New interpretation of the recent result of AMS-02 and multi-component decaying dark matters with non-Abelian discrete flavor symmetry

Yuji Kajiyama\textsuperscript{1,}a, Hiroshi Okada\textsuperscript{2,}b, Takashi Toma\textsuperscript{3,}c

\textsuperscript{1} Akita Highschool, Tegata-Nakadai 1, Akita 010-0851, Japan
\textsuperscript{2} School of Physics, KIAS, Seoul 130-722, Korea
\textsuperscript{3} Institute for Particle Physics Phenomenology, University of Durham, Durham DH1 3LE, UK

Received: 26 September 2013 / Accepted: 23 December 2013 / Published online: 5 February 2014
© The Author(s) 2014. This article is published with open access at Springerlink.com

Abstract Recently the AMS-02 experiment has released the data of positron fraction with a very small statistical error. Because of the small error, it is no longer easy to fit the data with single dark matter for a fixed diffusion model and dark matter profile. In this paper, we propose a new interpretation of the data: that it originates from decay of two-component dark matter. This interpretation gives a rough threshold of the lighter DM component. When DM decays into leptons, the positron fraction in the cosmic rays depends on the flavor of the final states, and this is fixed by imposing a non-Abelian discrete symmetry on our model. By assuming two gauge-singlet fermionic decaying DM particles, we show that a model with non-Abelian discrete flavor symmetry, e.g. $T_{13}$, can give a much better fitting to the AMS-02 data compared with a single-component dark matter scenario. Few dimension-six operators of the universal leptonic decay of DM particles are allowed in our model, since its decay operators are constrained by the $T_{13}$ symmetry. We also show that the lepton masses and mixings are consistent with current experimental data, due to the flavor symmetry.

1 Introduction

The latest experiment of Planck [1] tells us that about 26.8% of the energy density of the universe consists of Dark Matter (DM). Many experiments are being performed to search DM signatures. The recent result of the indirect detection experiment of AMS-02 [2] is in favor of previous experiments such as PAMELA [3,4] and Fermi-LAT [5], which had reported an excess of positron fraction in the cosmic rays. Moreover, it smoothly extends the anomaly line of positron fraction with energy up to about 350 GeV with a small statistical error compared with the previous experiments. These observations can, in general, be explained by scattering and/or decay of the GeV/TeV-scale DM particles. In addition, leptophilic DM is preferable, since PAMELA observed no antiproton excess [6]. Along this line of thought, several papers have been released [7–15]. Due to the smallness of the statistical error of AMS-02, it became difficult to make a fit to the data, in the same way as previous experiments like PAMELA [14].

In this paper, we show that we can obtain a better fitting to the data with two-component decaying DM. We introduce two kinds of fermionic DM particles, which are gauge-singlet fermions $X$ and $X'$, with leptons by dimension-six operators $\bar{L}E_LX/(\Lambda^2)$ due to the $T_{13}$ symmetry. DM particles decay into leptons via these operators with the suppression factor $\Lambda^2 \sim 10^{16}$ GeV, giving the desired lifetime of DM particles, $\Gamma^{-1} \sim (\text{TeV})^5/\Lambda^2 \sim 10^{26}$ s [17,18].

(i) The flavor of the final states of DM decay is determined by the $T_{13}$ symmetry.

(ii) The flavor of the final states of DM decay is determined by the $T_{13}$ symmetry.

We give a concrete example of the universal final states $X/X' \rightarrow \nu_e e^+ e^-/\nu_{\mu} \mu^+ \mu^-/\nu_{\tau} \tau^+ \tau^-$. Due to a specific selection rule by the flavor symmetry mentioned above, we show that the two-component DM model is preferable for the explanation of the precise AMS-02 result. In addition to that, we find a set of parameters that is consistent with the observed lepton masses and their mixings, especially a somewhat large angle of $\theta_{13}$, as recently reported by several experiments [19–25].
Table 1 The $T_{13}$ and $Z_3$ charge assignment of the SM fields and the Majorana DM $X$ and the Dirac DM $X'$, where $\omega = e^{2i\pi/3}$

|     | $Q$       | $U$       | $D$       | $L$       | $E$       | $H$       | $H'$     | $X$   | $X'$   |
|-----|-----------|-----------|-----------|-----------|-----------|-----------|----------|-------|-------|
| $SU(2)_L \times U(1)_Y$ | $2_{1/6}$  | $1_{2/3}$ | $1_{-1/3}$| $2_{-1/2}$| $1_{-1}$  | $2_{1/2}$ | $2_{1/2}$| $1_0$ | $1_0$ |
| $T_{13}$ | $1_{0,1,2}$| $1_{0,1,2}$| $1_{0,1,2}$| $3_1$     | $3_2$     | $3_1,3_2$ | $1_{0,1,2}$| $1_0$ | $1_1$ |
| $Z_3$   | $1$       | $\omega$  | $\omega^2$| $1$       | $1$       | $1$       | $\omega$  | $1$   | $1$   |

This paper is organized as follows. In Sect. 2, we briefly review our model based on the non-Abelian discrete group $T_{13}$, which is isomorphic to $Z_3 \times Z_3$ [16,26–29]. The $T_{13}$ group is a subgroup of $SU(3)$, and it is known as the minimal non-Abelian discrete group having two complex triplets as the irreducible representations; see Ref. [16] for details.

Lepton masses and mixings are derived from the setup shown in Table 1. Here, $Q$, $U$, $D$, $L$, $E$, $H(H')$, and $X(X')$ denote left-handed quarks, right-handed up-type quarks, right-handed down-type quarks, left-handed leptons, right-handed charged leptons, Higgs bosons, and gauge-singlet fermions, respectively. Here one should notice that $X$ and $X'$ are Majorana- and Dirac-type DM, respectively, which directly comes from the charge assignment of $T_{13}$. Due to the $T_{13}$ flavor symmetry in addition to an appropriate choice of the additional $Z_3$ symmetry, triplet Higgs bosons $H(3_1)$ and $H(3_2)$ couple only to leptons, while $T_{13}$ singlet Higgs bosons $H'(1_{0,1,2})$ couple only to quarks. Hence a linear combination of $H'$ is the SM-like Higgs boson and is created at LHC by gluon fusion. Therefore, the mass matrices of the quark sector are not constrained, while those of the lepton sector are determined by the $T_{13}$ symmetry. For the neutrino sector, since the Yukawa couplings $LH X$ and $LH X'$ are forbidden by the $T_{13}$ symmetry, the left-handed Majorana neutrino mass terms are derived from dimension-five operators $LHLH$. Here notice that $X$ and $X'$ have dimension-six operators $LHLX'$, $\bar{L}EX'$, and mass terms $m_X X X$, $m_{X'} \bar{X'} X'$. For the matter content and the $T_{13}$ assignment given in Table 1, the charged-lepton and neutrino masses are generated from the $T_{13}$ invariant operators

$$\mathcal{L}_Y = \sqrt{2}a_e \bar{E}LH^c(3_2) + \sqrt{2}b_e \bar{E}LH^c(3_1)$$

$$+ \frac{a_e}{\Lambda}LH(3_2)H(3_2) + \frac{b_e}{\Lambda}(LH(3_2)^3_2)_{32}(LH(3_1)_{31})_{32}$$

$$+ \frac{c_e}{\Lambda}(LH(3_2)_{31}(LH(3_1)_{31})_{31} + h.c.,$$

(2.1)

where $H^c = e H^*$, and $LH(3_2)LH(3_1)$ is $T_{13}$ invariant in two different products, corresponding to $b_e$ and $c_e$. The fundamental scale $\Lambda = 10^{11} \text{ GeV}$ is needed for the certain neutrino mass scale ($\Lambda/\sqrt{2} \sim 10^{10} \text{ GeV}$ is required to obtain the desired lifetime of DM, where $\lambda$ is the coupling constant of the DM decay operators, as we will discuss later). After the electroweak symmetry breaking, the Lagrangian Eq. (2.1) gives rise to mass matrices of the charged leptons $M_e$ and neutrinos $M_\nu$ as follows:

$$M_e = \begin{pmatrix} 0 & b_e v_1 & a_e \bar{v}_2 \\ a_e \bar{v}_3 & 0 & b_e v_2 \\ b_e v_3 & a_e \bar{v}_1 & 0 \end{pmatrix},$$

(2.2)

$$M_\nu = \frac{1}{\Lambda} \begin{pmatrix} c_\nu \bar{v}_3 v_1 & a_\nu \bar{v}_1^2 + b_\nu \bar{v}_3 v_1 & a_\nu \bar{v}_3^2 + b_\nu \bar{v}_2 v_3 \\ a_\nu \bar{v}_1^2 + b_\nu \bar{v}_3 v_1 & c_\nu \bar{v}_1 v_3 & a_\nu \bar{v}_2^2 + b_\nu \bar{v}_1 v_2 \\ a_\nu \bar{v}_2^2 + b_\nu \bar{v}_2 v_3 & a_\nu \bar{v}_1 v_2 & c_\nu \bar{v}_2 v_1 \end{pmatrix},$$

(2.3)

where the vacuum expectation values (VEVs) of the Higgs bosons are defined as $\langle H(3_1)^i \rangle = v_i/\sqrt{2}$, $\langle H(3_2)^i \rangle = \tilde{v}_i/\sqrt{2}$, $\langle H'(1_{0,1,2})^i \rangle = v'_i/\sqrt{2}$, $\sum_{i=1}^3 (v_i^2 + \tilde{v}_i^2 + v'_i) = (246 \text{ GeV})^2$.

Now we give a numerical example. By the following choice of parameters:

$$v_1 = 0.4269 \text{ GeV}, \quad v_2 = 16.11 \text{ GeV}, \quad v_3 = 7.862 \text{ GeV},$$

$$\tilde{v}_1 = 1 \text{ GeV}, \quad \tilde{v}_2 = 16.82 \text{ GeV}, \quad \tilde{v}_3 = 0.004836 \text{ GeV},$$

$$a_e = 0.1057, \quad b_e = 0, \quad a_\nu = -8.220 \times 10^{-3},$$

$$b_\nu = 8.439 \times 10^{-3}, \quad c_\nu = 3.632 \times 10^{-1},$$

(2.4)

the mass matrices of Eqs. (2.2) and (2.3) give rise to mass eigenvalues and related observables as follows:

---

1 All the assignments and particle contents are the same as our previous work [16] except the DM sector.
$m_e = 0.511 \text{ MeV}$, $m_\mu = 105.7 \text{ MeV}$, $m_\tau = 1.777 \text{ MeV}$, $m_{\nu_1} = 6.324 \times 10^{-3} \text{ eV}$, $m_{\nu_2} = 1.078 \times 10^{-2} \text{ eV}$, $m_{\nu_3} = 5.046 \times 10^{-2} \text{ eV}$, 
$\Delta m_{21}^2 = m_{\nu_2}^2 - m_{\nu_1}^2 = 7.62 \times 10^{-5} \text{ eV}^2$, 
$\Delta m_{32}^2 = m_{\nu_3}^2 - m_{\nu_2}^2 = 2.43 \times 10^{-3} \text{ eV}^2$, 
$\langle m \rangle_{ee} = 2.83 \times 10^{-4} \text{ eV}$, $\sum_i m_{\nu_i} = 5.49 \times 10^{-2} \text{ eV}$,
(2.5)
and the mixing matrices are given by

$U_{eL} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad U_{eR} = \begin{pmatrix} 0 & 0 & 1 \\ 1 & 0 & 0 \\ 0 & 1 & 0 \end{pmatrix},$
(2.6)

$U_{\text{MNS}} = U_{eL} U_\nu = \begin{pmatrix} 0.819 & 0.552 & -0.156 \\ -0.304 & 0.648 & 0.698 \\ -0.487 & 0.524 & -0.698 \end{pmatrix}$,
(2.7)
which are all consistent with the present experimental data [30,31]. In particular in the case of $U_{eL} = 1$, the mass matrices in Eqs. (2.2) and (2.3) require a normal hierarchy $m_{\nu_1} < m_{\nu_2} < m_{\nu_3}$ of the neutrino masses and $(U_{\text{MNS}})_{33} \neq 0$. A comprehensive analysis of the $T_{13}$ symmetric models has been made by several authors [32–35]. Although one can sweep the whole range of parameters, we adopt those of Eq. (2.4), giving universal final states due to the mixing matrices Eq. (2.6), since such an analysis is out of scope of the present paper.

As the Higgs sector, since the present model contains nine Higgs doublets, it causes flavor changing neutral current processes such as $K^0 - \bar{K}^0$ mixings. Therefore, extra Higgs bosons must be heavy enough. Moreover, additional massless bosons appear because the $T_{13}$ symmetric Higgs potential has accidental $U(1)$ symmetry. Therefore one can introduce soft $T_{13}$ breaking terms such as $H_i^{h1} H_i^{h1} + H_i^{h1} H_i^{h2} + H_i^{h2} H_i^{h2}$ and $H_i^{h1} \sum_i H(\bar{3}_2)^i$ in order to avoid those problems.

### 3 Decaying dark matter in the $T_{13}$ model

It is well known that the cosmic-ray anomalies measured by PAMELA [3] and Fermi-LAT [5] can be explained by DM decay with lifetime of $\Gamma^{-1} \sim 10^{26} \text{ s}$. If the DM ($X$ and $X'$ in our case) decays into leptons by dimension-six operators $L_i E_i X^{(i)}/\Lambda^2$ with $\Lambda \sim 10^{16} \text{ GeV}$, such a long lifetime can be achieved. In general, however, there exist several gauge invariant decay operators of lower dimensions; dimension-four operators inducing too rapid DM decay, and dimension-six operators including quarks, Higgs, and gauge bosons in the final states, which must be forbidden in a successful model. By the field assignment of Table 1, most decay operators listed in Table 2 [36] are forbidden due to the $T_{13}$ symmetry, except for $L_i E_i X^{(i)}$. Therefore, one does not have to worry about production of antiprotons and secondary positrons by scattering with a nucleon and the interstellar medium. With the notation $L_i = (v_i, \ell_i) = (U_{eL})_{ia}(v_a, \ell_a)$ and $E_i = (U_{eR})_{ib} \bar{E}_b \quad (i = 1, 2, 3; \quad \alpha, \beta = e, \mu, \tau)$, the four-Fermi decay interaction is explicitly written as

$$L_{\text{decay}} = \frac{\lambda_X}{\Lambda^2} \sum_{i=1}^3 (\bar{L}_i E_i) \bar{L}_i X$$

$$+ \frac{\lambda_X}{\Lambda^2} \sum_{i=1}^3 (\omega^{2(i-1)})(\bar{L}_i E_i) \bar{L}_i X' + \text{h.c.}$$

$$= \frac{\lambda_X}{\Lambda^2} \left[ \sum_{i=1}^3 (U_{eL})_{ia}^\dagger (U_{eR})_{ib} (U_{eL})_{ib}^\dagger \sum_{i=1}^3 (U_{eL})_{ia}^\dagger (U_{eR})_{ib} (U_{eL})_{ib}^\dagger \right] X \rightarrow X' + \text{h.c.},$$

(3.1)

where the factor $(\omega^{2(i-1)})$ is only for the case of $X'$ decay because of the multiplication rule of the $T_{13}$ flavor symmetry. As seen from Eq. (3.1), the decay mode of the DM particles $X$ and $X'$ depends on the mixing matrices $U_{eL}$ and $U_{eR}$, which are given in Eq. (2.6).

Next, we consider the decay width of the decaying DM through the $T_{13}$ invariant interaction Eq. (3.1). Due to the particular generation structure, the DM particles $X$ and $X'$ decay into a final state with several tri-leptons with a mixing-dependent rate. The decay width of DM $X$ per each flavor is defined as $\Gamma_{\alpha\beta\gamma} \equiv \Gamma(X \rightarrow \nu_\alpha \bar{E}_\beta \bar{E}_\gamma) + \Gamma(X \rightarrow \bar{\nu}_\alpha \ell_\beta \ell_\gamma)$.

Table 2 The higher dimensional operators which cause decay of $X$ and $X'$ up to dimension six [36]

| Dimensions | DM decay operators |
|------------|--------------------|
| 4          | $\bar{L} H^c X^{(i)}$ |
| 5          | $L E_i X^{(i)}$ |
| 6          | $L E_i X^{(i)}$, $H^i H^c X^{(i)}$, $(H^i)^c D_\alpha H^c E_\beta X^{(i)}$, $\bar{Q} D \bar{L} X^{(i)}$, $\bar{L} D \bar{Q} X^{(i)}$, $\bar{Q} \gamma_{\mu} D \bar{E}_\nu X^{(i)}$, $D^\alpha H^c D_\alpha \bar{L} X^{(i)}$, $D^\alpha D_\alpha H^c X^{(i)}$, $W_{\mu \nu}^a L \sigma^{\mu \nu} e \bar{H}^c X^{(i)}$ |
| $B_{\mu \nu}$, $W_{\mu \nu}^a$, and $D_\alpha$ are the field strength tensors of hypercharge gauge boson, weak gauge boson, and the electroweak covariant derivative |

$\bar{Q} D \bar{L} X^{(i)}$, $\bar{L} D \bar{Q} X^{(i)}$, $\bar{Q} \gamma_{\mu} D \bar{E}_\nu X^{(i)}$, $D^\alpha H^c D_\alpha \bar{L} X^{(i)}$, $D^\alpha D_\alpha H^c X^{(i)}$, $W_{\mu \nu}^a L \sigma^{\mu \nu} e \bar{H}^c X^{(i)}$ |

$\Delta \text{invariant decay operators of lower dimensions; dimension-four operators inducing too rapid DM decay, and dimension-six operators including quarks, Higgs, and gauge bosons in the final states, which must be forbidden in a successful model. By the field assignment of Table 1, most decay opera-
and the decay width $\Gamma_{a\beta\gamma}$ is calculated as
\[
\Gamma_{a\beta\gamma} = \frac{|\lambda_X|^2 m_X^5}{32 (4\pi)^2 \Lambda^4} (U_{a\beta\gamma} + U_{a\gamma\beta}),
\] (3.2)
where
\[
U_{a\beta\gamma} = \left| \sum_{i=1}^{3} (U_{eL})_{ia}^\ast (U_{eR})_{ib} (U_{eL})_{i\gamma}^\ast \right|^2.
\] (3.3)

The decay width of $X'$, denoted $\Gamma_{a'β'γ}'$, is obtained by replacing $X \to X'$. The differential decay width is written as
\[
\frac{d\Gamma_{aβγ}}{dx} = \frac{|\lambda_X|^2 m_X^5}{48 (4\pi)^2 \Lambda^4} x^2 \times \left( (6 - 2x)U_{aβγ} + (15 - 14x) U_{aγβ} \right),
\] (3.4)
where $x = 2E_{pF}/m_X$. This is required to enable one to calculate the energy distribution function of the injected $e^\pm$ from DM decay, $dN_{e\pm}/dE$. Here we have neglected the masses of the charged leptons in the final states. In both the $X$ and the $X'$ DM cases, the flavor dependent factor $U_{a\beta\gamma}$ gives a factor 3 when one takes the sum of flavor indices $α, β$ and $γ$. That is, not by a particular choice of parameters Eq. (2.4), but by the $T_{13}$ symmetry. Moreover, the branching fraction of each decay mode is given by $\text{Br}(X \to \nu_\alpha \ell_\beta^\pm \ell_\gamma^\mp, \overline{\nu}_\alpha \ell_\beta^\mp \ell_\gamma^\pm) = (U_{aβγ} + U_{aγβ})/6$. The DM mass $m_X$ and the total decay width $Γ_X = \sum_{a,\beta,\gamma} Γ_{aβγ}$ are chosen to be free parameters in the following analysis, since it can be always tuned with the coupling $λ_X$ and the cut-off scale $Λ$.

Given the differential decay width and the branching ratios, the primary source term of the positron and electron coming from DM decay at the position $r$ of the halo associated with our galaxy is expressed as
\[
Q(E, r) = n_X(r) Γ_X \sum_f \text{Br}(X \to f) \left( \frac{dN_{e\pm}}{dE} \right)_f + (X \to X'),
\] (3.5)
where $(dN_{e\pm}/dE)_f$ is the energy distribution of $e^\pm$ coming from the DM decay with the final state $f$, and $E$ is the energy of the injected $e^\pm$. We use the PYTHIA 8 [39] to evaluate the energy distribution function. Although it is often assumed that the relic density of the DM is thermally determined, nonthermal production of the DM dark matter is also possible [41]. We thus do not specify the origin of the relic DM in the following analysis, and we assume that the number densities of $X$ and $X'$ are the same for the simplest cases. The non-relativistic DM number density $n_X(r)$ is rewritten by $n_X(r) = ρ_X(r)/m_X$ with the DM profile $ρ_X(r)$. In this work, we adopt the Navarro–Frank–White (NFW) profile [42],
\[
ρ_{\text{NFW}}(r) = ρ_{⊙}(r_{⊙} + r_c)^2 / r(r + r_c)^2,
\] (3.6)
where $ρ_{⊙} \approx 0.40 \text{ GeV/cm}^3$ is the local DM density at the solar system, $r$ is the distance from the galactic center whose special values $r_{⊙} \approx 8.5 \text{ kpc}$ and $r_c \approx 20 \text{ kpc}$ are the distance to the solar system and the core radius of the profile, respectively. The diffusion equation must be solved to evaluate the $e^\pm$ flux observed at the Earth, and it depends on the diffusion model. The observable $e^\pm$ flux at solar system, $dΦ_{e\pm}/dE$, which is produced by DM decay is given by
\[
\frac{dΦ_{e\pm}}{dE} = \sum_{X,X'} \frac{ν_{e\pm}}{4πb(E)} \frac{ρ_{⊙}}{m_X} Γ_X \sum_f \text{Br}(X \to f)
\times \int_E^{m_X} \left( \frac{dN_{e\pm}}{dE'} \right)_f I(Φ_{X, E'}) dE',
\] (3.7)
where $b(E)$ is a space-independent energy loss coefficient written as $b(E) = E^2/(t_{⊙} \cdot 1 \text{ GeV})$ with $t_{⊙} = 5.7 \times 10^{15} \text{ s}$, and $I(Φ_{X, E'})$ is the reduced halo function at the solar system which is expressed by a Fourier–Bessel expansion [43]. A fitting function for the reduced halo function $I(λ_D)$ is given in Ref. [43] as a function of the single parameter $λ_D$, which is called the diffusion length. It is given by
\[
λ_D^2 = \frac{4K_0 t_{⊙}}{1 - δ} \left[ E_0^{δ-1} - E^{δ-1} \right],
\] (3.8)
where we use the following diffusion parameters: $δ = 0.70$, $K_0 = 0.0112 \text{ kpc}^2/\text{Myr}$, which is called MED. In addition, the diffusion zone is considered as a cylinder that sandwiches the galactic plane with height of $2L$ and radius $R$ where $L = 4 \text{ kpc}$ and $R = 20 \text{ kpc}$.

As seen from Eqs. (2.6) and (3.2), the DM decays into $e^\pm$ as well as $μ^\pm$ and $τ^\pm$ in equal rates. As a result, pure leptonic decays give dominant contributions, and it is consistent with no antiproton excess of the PAMELA results [6]. We may take into account the gamma-ray constraint since a lot of gamma rays are produced by the hadronization of $τ^\pm$. As we shall see below, the obtained lifetimes of the DM particles $X$ and $X'$ are roughly $τ_X, τ_{X'} \gtrsim 5 \times 10^{26} \text{ s}$. Thus, we do not need to consider the gamma-ray constraint seriously as long as we are comparing with Ref. [44].

3.1 Result for AMS-02

We use 31 data points of AMS-02 which are higher than 20 GeV for our chi-square analysis. The only statistical error is taken into account as the experimental errors here [2]. The positron fraction for the scenario of the leptonically decaying DM with $T_{13}$ symmetry is depicted in Fig. 1 with the experimental data of AMS-02 and PAMELA. The flux coming from only one-component DM is also shown in the figure for comparison. The obtained best fit point for one-component DM is $m_X = 521 \text{ GeV}$, $Γ_X^{-1} = 5.1 \times 10^{20} \text{ s}$ with $χ^2_{min} = 172.2$ (29 d.o.f.). For the single DM, the positron fraction at high
energy cannot be fit well as one can see from the figure. This is because the experimental data in the low-energy region $E \sim 20$ GeV has much higher precision, and the energy spectrum $dN/e^+/dE$ is fixed by the imposed flavor symmetry. Thus the predicted flux in the higher energy region is almost determined by the flavor symmetry. One should note that fitting with one-component DM would be better for different diffusion models or different DM halo profiles, since the evaluated $e^\pm$ flux has a large dependence on them.

On the other hand, the fitting parameters for two-component DM are

\[ m_X = 208 \text{ GeV}, \quad \Gamma_X^{-1} = 1.9 \times 10^{27} \text{ s}, \quad (3.9) \]
\[ m_{X'} = 1112 \text{ GeV}, \quad \Gamma_{X'}^{-1} = 4.7 \times 10^{26} \text{ s}, \quad (3.10) \]

with $\chi^2_{\text{min}} = 22.62$ (27 d.o.f) at the best fit point. Therefore a much better fitting is obtained with the two-component case. This is the result of multi-component DM and the fixed flavor of final states by $T_{13}$ symmetry. That is, not by the particular choice of parameters Eq. (2.4), but by the $T_{13}$ symmetry as mentioned below Eq. (3.4). A sharper drop-off is expected if we have a larger branching ratio for directly produced positrons.

4 Conclusions

We revisited a decaying DM model with a non-Abelian discrete symmetry $T_{13}$, and we extended it to the two-component DM scenario by adding extra DM $X'$. We have shown that our model is consistent with all the observed masses and mixings in the lepton sector. Also due to the specific selection rule of $T_{13}$, we have found that both DM particles have a universal decay coming from dimension-six operators, which gives a promising model for current indirect detection searches of DM.

Fitting to the positron fraction with single-component DM under the assumption of the MED diffusion model and the NFW DM profile can no longer give a good interpretation of the positron excess by DM decay because of the precise measurement of AMS-02. However, taking into account two-component DM as our model gives a much better fitting to the AMS-02 observation. The obtained parameters are $m_X = 208 \text{ GeV}$ with $\Gamma_X^{-1} = 1.8 \times 10^{27} \text{ s}$ and $m_{X'} = 1112 \text{ GeV}$ with $\Gamma_{X'}^{-1} = 4.7 \times 10^{26} \text{ s}$, assuming that $X$ and $X'$ have an equal number density.

Acknowledgments We would like to thank Prof. Shigeri Matsumoto for crucial advice. T.T. acknowledges support from the European ITN project (FP7-PEOPLE-2011-ITN, PITN-GA-2011-289442-INVISIBLES). The numerical calculations were carried out on SR16000 at YITP in Kyoto University.

Open Access This article is distributed under the terms of the Creative Commons Attribution License which permits any use, distribution, and reproduction in any medium, provided the original author(s) and the source are credited.

Funded by SCOAP3 / License Version CC BY 4.0.

References

1. P.A.R. Ade et al. [Planck Collaboration]. [arXiv:1303.5076 [astro-ph.CO]]
2. M. Aguilar et al., Phys. Rev. Lett. 110, 141102 (2013)
3. O. Adriani et al., Nature 458, 607 (2009)
4. O. Adriani et al., [PAMELA Collaboration], Science 332, 69 (2011). [arXiv:1103.4055 [astro-ph.HE]]
5. M. Ackermann et al. [Fermi LAT Collaboration], Phys. Rev. Lett. 108, 011103 (2012). [arXiv:1109.0521 [astro-ph.HE]]
6. O. Adriani et al., Phys. Rev. Lett. 102, 051101 (2009)
7. P.-F. Yin, Z.-H. Yu, Q. Yuan, X.-J. Bi, Phys. Rev. D 88, 023001 (2013). [arXiv:1304.4128 [astro-ph.HE]]
8. Q. Yuan, X.-J. Bi, Phys. Lett. B 727, 1–7 (2013). [arXiv:1304.2687 [astro-ph.HE]]
9. J. Koop, Phys. Rev. D 88, 076013 (2013). [arXiv:1304.1184 [hep-ph]]
10. A. De Simone, A. Riotto, W. Xue, JCAP 1305, 003 (2013). [arXiv:1304.1336 [hep-ph]]
11. Q. Yuan, X.-J. Bi, G.-M. Chen, Y.-Q. Guo, S.-J. Lin, X. Zhang. [arXiv:1304.1482 [astro-ph.HE]]
12. M. Ibe, S. Iwamoto, S. Matsumoto, T. Moroi, N. Yokozaki, JHEP 1308, 029 (2013). [arXiv:1304.1483 [hep-ph]]
13. T. Linden, S. Profumo, Astophys. J. 772, 18 (2013). [arXiv:1304.1791 [astro-ph.HE]]
14. I. Cholis, D. Hooper, Phys. Rev. D 88, 023013 (2013). [arXiv:1304.1840 [astro-ph.HE]]
15. H.-B. Jin, Y.-L. Wu, Y.-F. Zhou, JCAP 1311, 026 (2013). [arXiv:1304.1997 [hep-ph]]
16. Y. Kajiyama, H. Okada, Nucl. Phys. B 848, 303 (2011). [arXiv:1011.5753 [hep-ph]]
17. N. Haba, Y. Kajiyama, S. Matsumoto, H. Okada, K. Yoshioka, Phys. Lett. B 695, 476 (2011). [arXiv:1008.4777 [hep-ph]]
18. K. Hashimoto, H. Okada. [arXiv:1110.3640 [hep-ph]]
19. T2K Collaboration, K. Abe et al., Phys. Rev. Lett. 107, 041801 (2011). 1106.2822
20. J.K. Ahn et al. [RENO Collaboration], Phys. Rev. Lett. 108, 191802 (2012). [arXiv:1204.0626 [hep-ex]]
21. DAYA-BAY Collaboration, F. An et al., Phys. Rev. Lett. 108, 171803 (2012). 1203.1669
22. DOUBLE-CHOOZ Collaboration, Y. Abe et al., Phys. Rev. Lett. 108, 131801 (2012). 1112.6353
23. CHOOZ Collaboration, M. Apollonio et al., Eur. Phys. J. C 27, 331 (2003). [arXiv:hep-ex/0301017]
24. Palo Verde Collaboration, F. Boehm et al., Phys. Rev. D 64, 112001 (2001). [arXiv:hep-ex/0107009]
25. MINOS Collaboration, P. Adamson et al., Phys. Rev. Lett. 107, 181802 (2011). 1108.0015
26. W.M. Fairbairn, T. Fulton, J. Math. Phys. 23, 1747 (1982)
27. S.F. King, C. Luhn, JHEP 0910, 093 (2009)
28. For a review of non-Abelian discrete symmetry, H. Ishimori, T. Kobayashi, H. Ohki, H. Okada, Y. Shimizu, M. Tanimoto, Lect. Notes Phys. 858, 1 (2012)
29. H. Ishimori, T. Kobayashi, H. Ohki, H. Okada, Y. Shimizu, M. Tanimoto, Prog. Theor. Phys. Suppl. 183, 1 (2010)
30. J. Beringer et al. [Particle Data Group Collaboration], Phys. Rev. D 86, 010001 (2012)
31. D.V. Forero, M. Tortola, J.W.F. Valle, Phys. Rev. D 86, 073012 (2012). [arXiv:1205.4018 [hep-ph]]
32. K.M. Parattu, A. Wingerter, Phys. Rev. D 84, 013011 (2011). [arXiv:1012.2842 [hep-ph]]
33. G.-J. Ding, Nucl. Phys. B 853, 635 (2011). [arXiv:1105.5879 [hep-ph]]
34. C. Hartmann, A. Zee, Nucl. Phys. B 853, 105 (2011). [arXiv:1106.0333 [hep-ph]]
35. C. Hartmann, Phys. Rev. D 85, 013012 (2012). [arXiv:1109.5143 [hep-ph]]
36. F. del Aguila, S. Bar-Shalom, A. Soni, J. Wudka, Phys. Lett. B 670, 399 (2009). [arXiv:0806.0876 [hep-ph]]
37. G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012)
38. S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012)
39. T. Sjostrand, S. Mrenna, P.Z. Skands, JHEP 0605, 026 (2006)
40. T. Sjostrand, S. Mrenna, P.Z. Skands, Comput. Phys. Commun. 178, 852 (2008)
41. T. Moroi, L. Randall, Nucl. Phys. B 570, 455 (2000). [arXiv:hep-ph/9906527]
42. J.F. Navarro, C.S. Frenk, S.D.M. White, Astrophys. J. 490, 493 (1997)
43. M. Cirelli, G. Corcella, A. Hektor, G. Hutsi, M. Kadastik, P. Panci, M. Raidal, F. Sala et al., JCAP 1103, 051 (2011). [arXiv:1012.4515 [hep-ph]. [Erratum-ibid. 1210, E01 (2012)]]
44. M. Ackermann et al. [LAT Collaboration], Phys. Rev. D 86, 022002 (2012). [arXiv:1205.2739 [astro-ph.HE]]