Direct and Indirect Detection of Dark Matter in $D_6$ Flavor Symmetric Model

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Abstract

We study a fermionic dark matter in a non-supersymmetric extension of the standard model with a family symmetry based on $D_6 \times \hat{Z}_2 \times Z_2$. In our model, the final state of the dark matter annihilation is determined to be $e^+e^-$ by the flavor symmetry, which is consistent with the PAMELA result. At first, we show that our dark matter mass should be within the range of 230 GeV − 750 GeV in the WMAP analysis combined with $\mu \to e\gamma$ constraint. Moreover we simultaneously explain the experiments of direct and indirect detection, by simply adding a gauge and $D_6$ singlet real scalar field. In the direct detection experiments, we show that the lighter dark matter mass $\simeq 230$ GeV and the lighter standard model Higgs boson $\simeq 115$ GeV is in favor of the observed bounds reported by CDMS II and XENON100. In the indirect detection experiments, we explain the positron excess reported by PAMELA through the Breit-Wigner enhancement mechanism. We also show that our model is consistent with no antiproton excess suggested by PAMELA.
1 Introduction

The existence of the dark matter (DM) in the Universe has been established by measurements. The WMAP experiment tells us that the amount of the DM is considered about 23% of energy density of the Universe [1]. As indirect detection experiments of the DM, PAMELA reported excess of positron fraction in the cosmic ray [2]. This observation can be explained by annihilation and/or decay of DM particles with mass of $\mathcal{O}(10^{2-3})$ GeV. The PAMELA experiment searches antiproton as well in the cosmic ray, and it is consistent with the background [3]. Therefore, if these signals are from annihilation and/or decay processes of DM particles, this indicates that the leptophilic DM is preferable. However, even if the DM is leptophilic, the resultant positron fraction depends on the generation of final state leptons. For instance, if the final state of annihilation and/or decay of the DM is $\tau^+\tau^-$, it will overproduce gamma-rays as final state radiation [4]. Therefore it is considerable that the leptophilic DM can reflect flavor structure of elementary particles. In this point of view, several works discussing the DM and flavor structure have been done so far [5, 6, 7, 8, 9, 10, 11, 12, 13].

Flavor structure of elementary particles is thought to be determined by symmetry, so called flavor symmetry [14]. In our previous work [13], we have discussed fermionic DM model with the standard model (SM) extension with the $D_6$ flavor symmetry [15]. In this model, three generations of matter fields including right-handed neutrinos are embedded into doublet and singlet representations of $D_6$ group in particular way. The light neutrino masses are induced by radiative correction through inert $SU(2)_L$ doublet Higgs bosons $\eta$ which do not have vacuum expectation values(VEVs) [16, 17]. We identify a heavy Majorana neutrino of $D_6$ singlet $n_S$ with the DM candidate. The DM $n_S$ is stable because of the additional $Z_2$ symmetry. Since the $D_6$ symmetry completely determines flavor structure of the model, the final states of annihilation of the DM via Dirac Yukawa interaction $\eta^\dagger L n_S$ are fixed to be electron-positron pair and $\tau$ neutrino pair. In that paper, we have found that the DM mass is constrained to be in the range $230$ GeV $- 750$ GeV from the condition of the relic abundance and $\mu \rightarrow e\gamma$ process. However, this annihilation of the DM via $\eta$-mediated t- and u-channel processes does not give enough s-wave contribution to the cross section because it is proportional to mass of the final state $m_{e,\nu}$. Therefore, this model requires very large enhancement ($\sim 10^6$) of the cross section at the present Universe compared with that at the early Universe to explain the PAMELA data, which is not realistic.

In this paper, we extend our $D_6$ model of Ref.[13] by adding gauge and $D_6$ singlet scalar field $\varphi$, which couples with the DM $n_S$ as $\varphi n_S n_S$. While the final states of the DM annihilation are the same as those of the previous model, there exist s-channel annihilation processes mediated by $\varphi$. 
In this case, the Breit-Wigner enhancement mechanism works which can give enough boost factor [18, 19, 20]. Moreover, the new field $\varphi$ mixes with the $D_6$ singlet Higgs doublet $\phi_S$, which is responsible for mass of the quark sector. This mixing can simultaneously induce antiproton production by DM annihilation and interaction with quarks in atoms. We find that the spin-independent cross section of the DM and quarks via the mixing between $\varphi$ and $\phi_S$ can be close to sensitivities of direct DM detection experiments such as CDMS II [21] and XENON100 [22], suppressing antiproton flux in the cosmic ray.

This paper is organized as follows. In section 2, we review our model briefly and summarize the predictions for lepton sector coming from the flavor symmetry. In section 3, we analyze the Higgs potential and mixing between the SM Higgs and new singlet scalar $\varphi$. In section 4, we show constraints of DM mass from WMAP and $\mu \rightarrow e\gamma$ process. We discuss direct and indirect detection of DM in section 5 and 6, respectively. Section 7 is devoted to conclusions and discussions.

2 The Model

In this section, we briefly review a SM extension with $D_6 \times \hat{Z}_2 \times Z_2$ family symmetry [13].

2.1 Yukawa couplings

We introduce three “generations” of Higgs doublets $\phi_{I,S}$, inert doublets $\eta_{I,S}$, and one generation of inert singlet $\varphi$. Where $I = 1, 2$ and $S$ denote $D_6$ doublet and singlet, respectively, and assume that each field is charged in specific way under the family symmetry shown in Table 1.

| $SU(2)_L \times U(1)_Y$ | $L_S$ | $n_S$ | $e^c_S$ | $L_I$ | $n_I$ | $e^c_I$ |
|------------------------|------|------|--------|------|------|--------|
| $D_6$                  | $1$  | $1'$ | $1$    | $2'$ | $2'$ | $2'$   |
| $\hat{Z}_2$            | +    | +    | -      | +    | +    | -      |
| $Z_2$                  | +    | -    | +      | -    | +    | +      |

Table 1: The $D_6 \times \hat{Z}_2 \times Z_2$ assignment for the leptons. The subscript $S$ indicates a $D_6$ singlet, and the subscript $I$ running from 1 to 2 stands for a $D_6$ doublet. $L_I$ and $L_S$ denote the $SU(2)_L$-doublet leptons, while $e^c_I$, $e^c_S$, $n_I$ and $n_S$ are the $SU(2)_L$-singlet leptons.
Under the $Z_2$ symmetry (which plays the role of $R$ parity in the MSSM), only the right-handed neutrinos $n_I, n_S$ and the inert Higgs doublets $\eta_I, \eta_S$ are odd. All quarks are assumed to be singlet under the family symmetry so that the quark sector is basically the same as the SM, where the $D_6$ singlet Higgs doublet $\phi_S$ with $(+, +)$ of $\hat{Z}_2 \times Z_2$ plays a role in the SM Higgs in the quark sector. No other Higgs bosons can couple to the quark sector at the tree-level. In this way we can avoid tree-level flavor changing neutral currents (FCNCs) in the quark sector. The $\hat{Z}_2$ symmetry is introduced to forbid tree-level couplings of the $D_6$ singlet Higgs $\phi_S$ with $L_I, L_S, n_I$ and $n_S$, simultaneously to forbid tree-level couplings of $\phi_I, \eta_I$ and $\eta_S$ with quarks. As shall be discussed later, the gauge singlet $\varphi$ plays an important role in explaining an indirect detection reported by PAMELA. Furthermore, it is expected to explain the direct detection as CDMS II, because our dark matter $n_S, D_6$ singlet right-handed neutrino, couples to the quark sector by small mixing between $\varphi$ and $\phi_S$, which should be estimated to satisfy the experimental results. We will show the numerical analysis of the mixing for both experiments later.

The most general renormalizable $D_6 \times \hat{Z}_2 \times Z_2$ invariant Yukawa interactions in the lepton sector are found to be

$$
L_Y = \sum_{a,b,d=1,2,S} \left[ Y_{ab}^{ed}(L_a \sigma_2 \phi_d)e^c_b + Y_{ab}^{vd}(\eta_d^t L_a)n_b \right]
- \sum_{l=1,2} \left( \frac{M_1}{2} n_l n_I - \frac{M_S}{2} n_S n_S - \sum_{l=1,2} \frac{\mathcal{G}_1}{2} \varphi n_l n_I - \frac{\mathcal{G}_S}{2} \varphi n_S n_S + h.c. \right),
$$

where the coupling constants $\mathcal{G}_{1,S}$ are complex in general. The electroweak symmetry is broken by the VEVs $\langle \phi_1 \rangle = \langle \phi_2 \rangle \equiv v_D/2, \langle \phi_S \rangle = v_S/\sqrt{2}, V^2 \equiv v_D^2 + v_S^2 = (246 \text{ GeV})^2, \langle \eta_{I,S} \rangle = \langle \varphi \rangle = 0$ [23], and we obtain the following mass matrix $M_e$ and diagonalization matrix $U_{eL}$ of $M_e M_e^\dagger$ in the

| SU(2)_L \times U(1)_Y | $\phi_S$ | $\phi_I$ | $\eta_S$ | $\eta_I$ | $\varphi$ |
|------------------------|-------|-------|-------|-------|-------|
| $(2, -1/2)$ | $(2, -1/2)$ | $(2, -1/2)$ | $(2, -1/2)$ | $(1, 0)$ |
| $D_6$ | $1$ | $2'$ | $1''$ | $2'$ | $1$ |
| $\hat{Z}_2$ | $+$ | $-$ | $+$ | $+$ | $+$ |
| $Z_2$ | $+$ | $+$ | $-$ | $-$ | $+$ |

Table 2: The $D_6 \times \hat{Z}_2 \times Z_2$ assignment for the Higgs bosons.
As a result of this discussion, we can assume that DM, it mainly couples with electron (and positron) with large coupling as can been seen from Eq. (2.5), if one identifies the Higgs potential. When the seesaw mechanism at one-loop level [16]. In this mechanism, Majorana mass is proportional to origin of the maximal mixing of atmospheric neutrino mixing. Only an inverted mass spectrum does not work. Light Majorana neutrino masses are generated by radiative seesaw mechanism at one-loop level [16]. In this mechanism, Majorana mass is proportional to \( h_i^2 \kappa V^2 M/(16 \pi^2 (M^2 - m^2_\eta)) \), where \( \kappa \) denotes typical coupling constant of non self-adjoint terms in the Higgs potential. When \( \kappa = 0 \), an exact lepton number \( U(1)_L \) invariance is recovered, where the right-handed neutrinos \( n_{L,S} \) are neutral under \( U(1)_L \) in contrast to the conventional seesaw models. This \( U(1)_L \) forbids the neutrino masses, so that the smallness of the neutrino masses has a natural meaning. Now we can derive some predictions of our model based on the family symmetry:

1. If \( \epsilon_{e,\mu} = 0 \), the mixing matrix \( U_{eL} \) has the maximal mixing in its right-upper block which is the origin of the maximal mixing of atmospheric neutrino mixing. Only an inverted mass spectrum \( m_{\nu 3} < m_{\nu 1,2} \) is allowed.

2. Non-zero \( \theta_{13} \) is predicted as \( \sin^2 \theta_{13} \simeq \epsilon^2 = 1.2 \times 10^{-5} \). This small value of \( \theta_{13} \) is consistent with the best fit value \( 0.020^{+0.008}_{-0.009} \) with 1\( \sigma \) error [24].

3. The effective Majorana mass \( \langle m \rangle_{ee} \) is bounded from below as \( \langle m \rangle_{ee} \gtrsim 0.02 \) eV.

As a result of this discussion, we can assume that \( M_{1,S} = \mathcal{O}(\text{TeV}), \kappa \ll 1 \) and \( h_i = \mathcal{O}(1) \). Moreover, as can been seen from Eq. (2.3), if one identifies the \( D_6 \) singlet right-handed neutrino \( n_S \) to be the DM, it mainly couples with electron (and positron) with large coupling \( h_3 \sim 1 \). This selection rule is
remarkably determined by the family symmetry. These facts play a crucial role in the study of cold DM (CDM) as discussed below.

## 3 Higgs Potential

In this section, we analyze the Higgs potential. As discussed in Refs. [13, 23], the Higgs potential consists of $D_6$ symmetric and breaking terms. Since $D_6$ invariant Higgs potential has an accidental global $O(2)$ symmetry, the latter must be introduced in order to forbid massless Nambu-Goldstone (NG) bosons. Essentially, such soft $D_6$ breaking terms are mass terms of the Higgs bosons. For the potential of $(\phi_I, \phi_S)$, the soft $D_6$ breaking mass terms [23] are given by

$$V(\phi)_{\text{soft}} = \mu^2_2(\phi_2^\dagger \phi_1 + \phi_1^\dagger \phi_2) + \left( \mu^2_4 \phi_S^\dagger (\phi_1 + \phi_2) + h.c. \right),$$  

(3.1)

where $\mu^2_2$ is real while $\mu^2_4$ is complex in general. The mass term of $(\phi_I, \phi_S)$ is dominated by Eq. (3.1), and subdominantly given by $D_6$ invariant terms of order $V^2$. One finds that the $D_6$ breaking terms Eq. (3.1) preserve the minimum symmetry $S_2$ under $\phi_1 \leftrightarrow \phi_2$. The key point is that the $S_2$ invariance is required not only to ensure the vacuum alignment $\langle \phi_1 \rangle = \langle \phi_2 \rangle \neq 0$ but also to forbid NG bosons which violate the electroweak precision test of the SM.

Since the Higgs potential of $\phi_{I,S}$ and $\eta_{I,S}$ are analyzed in Ref. [13], we do not explicitly show that here again. In the present model, the new field $\varphi$ is introduced and it plays an important role in our analysis. Therefore we explicitly show the potential including $\varphi$. The most general renormalizable $D_6 \times \hat{Z}_2 \times Z_2$ invariant Higgs potential of $\varphi$ is given by

$$V(\varphi) = m_I^2 \varphi + m_2^2 \varphi^2 + m_3 \varphi^3 + \lambda_1 \varphi^4,$$

(3.2)

$$V(\phi, \varphi) = m_4 (\phi_S^\dagger \phi_S) \varphi + m_5 (\phi_I^\dagger \phi_I) \varphi + \lambda_2 (\phi_S^\dagger \phi_S) \varphi^2 + \lambda_3 (\phi_I^\dagger \phi_I) \varphi^2,$$

(3.3)

$$V(\eta, \varphi) = V(\phi, \varphi)(\phi \rightarrow \eta),$$

(3.4)

where all parameters are considered to be real without loss of generality. By using the decomposition of $SU(2)_L$ doublets $\phi_{I,S}$,

$$\phi_I = \frac{1}{\sqrt{2}} \begin{pmatrix} v_D / \sqrt{2} + \rho_I + i \sigma_I \\ \sqrt{2} \phi_I^- \end{pmatrix}, \quad \phi_S = \frac{1}{\sqrt{2}} \begin{pmatrix} v_S + \rho_S + i \sigma_S \\ \sqrt{2} \phi_S^- \end{pmatrix},$$

(3.5)

we find the mass matrix of neutral Higgs bosons as

$$H^T M_h^2 H = \frac{1}{2} \begin{pmatrix} \rho & \sigma & \varphi \\ \rho \sigma & \rho & \rho \sigma \\ \rho \sigma & \rho \sigma & \rho \end{pmatrix} \begin{pmatrix} M_{\rho,\rho} & M_{\rho,\sigma} & M_{\rho,\varphi} \\ M_{\rho,\sigma} & M_{\sigma,\sigma} & M_{\sigma,\varphi} \\ M_{\rho,\varphi} & M_{\sigma,\varphi} & M_{\varphi,\varphi} \end{pmatrix} \begin{pmatrix} \rho \\ \sigma \\ \varphi \end{pmatrix},$$

(3.6)
where \( \rho = (\rho_I, \rho_S) \), \( \sigma = (\sigma_I, \sigma_S) \). Each 3 \times 3 element \( M^2_{\rho,\rho} \)’s are given by [13]

\[
M^2_{\rho,\rho} \approx \begin{pmatrix}
0 & 2\mu_2^2 & \sqrt{2}\text{Re}(\mu_4^2) \\
2\mu_2^2 & 0 & \sqrt{2}\text{Re}(\mu_4^2) \\
\sqrt{2}\text{Re}(\mu_4^2) & \sqrt{2}\text{Re}(\mu_4^2) & 0
\end{pmatrix} + \begin{pmatrix}
a_{\rho,\rho} v^2_D & a_{\rho,\rho} v^2_D & b_{\rho,\rho} v_D v_S \\
a_{\rho,\rho} v^2_D & a_{\rho,\rho} v^2_D & b_{\rho,\rho} v_D v_S \\
b_{\rho,\rho} v_D v_S & b_{\rho,\rho} v_D v_S & c_{\rho,\rho} v^2_S
\end{pmatrix}, \quad (3.7)
\]

\[
M^2_{\sigma,\sigma} \approx \begin{pmatrix}
0 & 2\mu_2^2 & \sqrt{2}\text{Re}(\mu_4^2) \\
2\mu_2^2 & 0 & \sqrt{2}\text{Re}(\mu_4^2) \\
\sqrt{2}\text{Re}(\mu_4^2) & \sqrt{2}\text{Re}(\mu_4^2) & 0
\end{pmatrix} + \begin{pmatrix}
a_{\sigma,\sigma} v^2_D + a'_{\sigma,\sigma} v^2_S & b_{\sigma,\sigma} v^2_D & c_{\sigma,\sigma} v_D v_S \\
b_{\sigma,\sigma} v^2_D & a_{\sigma,\sigma} v^2_D + a'_{\sigma,\sigma} v^2_S & c_{\sigma,\sigma} v_D v_S \\
c_{\sigma,\sigma} v_D v_S & c_{\sigma,\sigma} v_D v_S & d_{\sigma,\sigma} v^2_D
\end{pmatrix}, \quad (3.8)
\]

\[
M^2_{\rho,\sigma} \approx \begin{pmatrix}
0 & 0 & \sqrt{2}\text{Im}(\mu_4^2) \\
0 & 0 & \sqrt{2}\text{Im}(\mu_4^2) \\
\sqrt{2}\text{Im}(\mu_4^2) & \sqrt{2}\text{Im}(\mu_4^2) & 0
\end{pmatrix} + \begin{pmatrix}
a_{\rho,\sigma} v^2_S & 0 & -b_{\rho,\sigma} v_D v_S \\
0 & a_{\rho,\sigma} v^2_S & -b_{\rho,\sigma} v_D v_S \\
b_{\rho,\sigma} v_D v_S & b_{\rho,\sigma} v_D v_S & c v^2_D
\end{pmatrix}, \quad (3.9)
\]

where the coefficients \( a_{\rho,\rho} \)’s are of \( O(1) \). The \( \varphi \)-dependent terms are given by

\[
\rho M^2_{\rho,\varphi} = (\rho_1, \rho_2, \rho_S) \begin{pmatrix}
v_D m_5/\sqrt{2} \\
v_D m_5/\sqrt{2} \\
v_S m_4/\sqrt{2}
\end{pmatrix} \varphi, \quad (3.10)
\]

\[
M^2_{\varphi,\varphi} = 2m_2^2 + v^2_S \lambda_2 + v^2_D \lambda_3. \quad (3.11)
\]

The stable minimum conditions are found by partially differentiating the potential by \( \varphi \) as

\[
\frac{\partial V}{\partial \varphi} \bigg|_{\varphi \to 0} = m_1^3 + \frac{1}{2} \left(v^2_S m_4 + v^2_D m_5\right) = 0, \quad (3.12)
\]

and

\[
\frac{\partial^2 V}{\partial \varphi^2} \bigg|_{\varphi \to 0} = M^2_{\varphi,\varphi}, \quad \frac{\partial^2 V}{\partial \varphi \partial v_{S(D)}} \bigg|_{\varphi \to 0} = \frac{1}{\sqrt{2}} v_{S(D)} m_{4(5)}. \quad (3.13)
\]

Therefore, we obtain the vacuum conditions for \( \langle \phi_I, S \rangle \neq 0 \) and \( \langle \varphi \rangle = 0 \) as

\[
m_1^3 + \frac{1}{2} \left(v^2_S m_4 + v^2_D m_5\right) = 0, \quad M^2_{\varphi,\varphi} > 0, \quad v_{S(D)} m_{4(5)} > 0. \quad (3.14)
\]

The mass matrix \( M^2_{\rho} \) is diagonalized by the 7 \times 7 orthogonal matrix \( O \), as \( O M^2_{\rho} O^T \). Notice that quarks couple only with \( \phi_S \) via Yukawa interactions, and \( \phi_S(\rho_S) \) mixes with \( \varphi \) via \( m_4 \). This mixing
parameter $m_4$ will induce both interaction of the DM with atoms (direct detection) and antiproton flux in the cosmic ray. We will discuss these DM phenomenology below. Note also that there is no mixing between $\phi$ and $\eta$ because $\eta_{I,S}$ do not get VEVs.

The SM Higgs is described in terms of the linear combination of flavor eigenstate fields as

$$SM - Higgs = O_{11}\rho_1 + O_{12}\rho_2 + O_{13}\rho_S + O_{14}\sigma_1 + O_{15}\sigma_2 + O_{16}\sigma_S + O_{17}\phi,$$

and the other combinations correspond to heavy neutral Higgs bosons with mass of several hundred GeV. Therefore the $\rho_S - \phi$ mixing is proportional to $O_{31}^T O_{71}^T$. In the following analysis, we give numerical values of the matrix $O$.

### 4 WMAP and $\mu \rightarrow e\gamma$ Constraint

In this section, we derive conditions for mass of the DM $M_S$ and charged component of $\eta$ boson $M_\eta$, following the result of Ref. [13].

#### 4.1 $\mu \rightarrow e\gamma$ Constraint

The DM mass $M_S$ is constrained from the $\mu \rightarrow e\gamma$ process. The branching fraction of $\mu \rightarrow e\gamma$ from Fig. 1 is approximately given by

$$B(\mu \rightarrow e\gamma) = \frac{3\alpha}{64\pi G_F} X^4 \approx |X^2 900 \text{ GeV}^2|^2, \quad X^2 \approx h_3^2 \frac{m_e}{m_\mu} \frac{F_2(M_S^2/M_\eta^2)}{M_\eta^2},$$

and

$$F_2(x) = \frac{1 - 6x + 3x^2 + 2x^3 - 6x^2 \ln x}{6(1 - x)^4},$$

where $x = M_S^2/M_\eta^2$. As we can see from the Yukawa matrix of Eq. (2.5), only $\eta_S$ couples to $n_S$ with $e_L$ and with $\mu_L$, where the coupling with $\mu_L$ is suppressed by $m_e/m_\mu \approx 0.005$. In the next subsection, we will obtain the constraints of the DM mass $M_S$ which is consistent with the observed DM relic density $\Omega_d h^2 \approx 0.12 \ [1]$ and $\mu \rightarrow e\gamma$, assuming $n_S$ to be the DM.

#### 4.2 WMAP

In the analysis of Ref. [13], we have found that it is more natural and promising that only $n_S$ of three right-handed neutrinos remains as a fermionic CDM candidate. Furthermore since charged
component of $\eta_S$ boson couples to $e_L$ and $n_S$ due to our original matrix in Eq.(2.5), it remarkably leads to be a clean signal if the charged extra Higgs boson $\eta_S$ is produced at LHC.

We simply find the thermally averaged cross section $\langle \sigma_1 v \rangle$ for the annihilation of two $n_S$’s [25] from Fig.2 in the limit of the vanishing final state lepton masses:

$$ \langle \sigma_1 v \rangle = a_1 + b_1 \frac{6}{x} + \cdots, \quad a_1 = 0, \quad b_1 = \frac{|h_3|^2 r^2 (1 - 2r + 2r^2)}{24\pi M_\eta^2}, $$

(4.3)

$$ r = \frac{M_S^2}{(M_\eta^2 + M_S^2)}, \quad x = \frac{M_S}{T} \quad (4.4) $$

where $M_\eta$ is $\eta_S$ mass, $M_S$ is $n_S$ mass which is our DM candidate and $T$ is temperature of the Universe. The thermally averaged cross section Eq.(4.3) does not contain s-wave contribution as a consequence of massless limit of the final state particles, and we find that the allowed region for the DM mass is around $\mathcal{O}(10^2)$ GeV from the constraints of WMAP results [1] and $\mu \rightarrow e\gamma$ decay.

In Fig.2 we present the allowed region in the $M_\eta - M_S$ plane, in which $\Omega_n h^2 = 0.12$ and $B(\mu \rightarrow e\gamma) < 1.2 \times 10^{-11}$ [26] are satisfied, where we take $|h_3| < 1.5$. From Eq.(4.3), retaining $h_3 = \mathcal{O}(1)$
is quite important to find the promising DM mass regions, as we mentioned before. Note that there is no allowed region even for $|h_3| \lesssim 0.8$. As can been seen from Fig.3 we find the mass range as follows:

$$230 \text{ GeV} < M_S < 750 \text{ GeV}, \quad 300 \text{ GeV} < M_\eta < 750 \text{ GeV}.$$ (4.5)

In this analysis, we have calculated the mass bound for Sunyaev and Zeldovich (SZ) effect [27]. In our model, $\eta^+_S$, which decays to high energy $e^+_L$, may affect the CMB by the inverse Compton scattering, if the lifetime is not between $10^{-(5-7)}$ sec. From the condition that the lifetime of $\eta^+_S$ comes into the allowed region, mass $M_\eta$ has the bound of $30 \text{ GeV} < M_\eta < 750 \text{ GeV}$. Where the Yukawa coupling nearly equals to 1, and $M_\eta \gg M_S$ are assumed. Hence, one finds that the SZ effect satisfies the both constraints of $\mu \rightarrow e\gamma$ and cosmological pair annihilation of CDMs sufficiently.

![Figure 3: The allowed region in the $M_\eta - M_S$ plane in which $\Omega_d h^2 = 0.12, B(\mu \rightarrow e\gamma) < 1.2 \times 10^{-11}$ and $|h_3| < 1.5$ are satisfied.](image)

## 5 Direct Detection

We analyze the direct detection search through the experiments of CDMS II [21] and XENON100 [22]. The main contribution to the spin-independent cross section is from the t-channel diagram with the mixing between $\varphi$ and $\phi_S$, as depicted in Fig.4. Then the resultant cross section for a proton is given by

$$\sigma_{SI}^{(p)} = \frac{4}{\pi} \left( \frac{m_p M_S}{m_p + M_S} \right)^2 |f_p|^2,$$ (5.1)
with the hadronic matrix element

$$f_p \frac{m_p}{m_q} = \sum_{q=u,d,s} f^{(p)}_{Tq} \frac{\alpha_q}{m_q} + \frac{2}{27} \sum_{q=c,b,t} f^{(p)}_{Tq} \frac{\alpha_q}{m_q},$$

(5.2)

where $m_p$ is the proton mass. The effective vertex $\alpha_q$ in our case is given by

$$\alpha_q \approx \frac{O_{31}^T O_{71}^T S_q Y_q}{m_{SM-Higgs}^2}.$$  

(5.3)

Here $m_{SM-Higgs}$ is the SM Higgs mass and $Y_q \propto (m_q/v)$ is a Yukawa coupling constant of the quark sector. Notice that the quark sector couples only to $\phi_S$. In the numerical analysis, we set the Higgs masses to avoid the lepton flavor violation (LFV) process as follows:

$$115 \text{ GeV} \leq m_{SM-Higgs} \leq 200 \text{ GeV}, \quad 500 \text{ GeV} \lesssim \text{other six neutral Higgs boson masses}.$$  

(5.4)

Under this setup, the elastic cross section is shown in Fig.5. Where we set $|O_{31}^T O_{71}^T S_q| = 0.1$, which we call the “mixing”. We plot the DM mass $M_S$ in the region $10 - 1000$ GeV. Since the allowed region of the DM mass is $230 \text{ GeV} - 750 \text{ GeV}$ from the WMAP analysis combined with $\mu \rightarrow e\gamma$ constraint, rather smaller SM Higgs mass is favored if these experiments could detect the DM near the current bound.

6 Indirect Detection

The PAMELA experiment implies that there could be positron excess [2], but not be antiproton excess [3]. In order to describe the PAMELA results successfully through an annihilation process of the DM, we need enhancement of the cross section by using the Breit-Wigner mechanism [18].
Figure 5: The spin-independent cross section as a function of the DM mass for the direct detection [21, 22]. The “mixing” is defined by $|O^{T}_{31}O^{T}_{71}\mathcal{S}|$ and set to be 0.1. The longitudinal black line represents the SM Higgs boson mass range.

### 6.1 Positron Production from DM annihilation

The main channel of the DM annihilation in the present Universe is depicted in Fig.6. The $n_s$ annihilation cross section to leptons is given by

$$ (\sigma v) \simeq \frac{4}{\pi(4\pi)^4} \frac{m^2_M^4 |h_3|^4 (O_{R7}^2)^4 (\text{Im}\mathcal{S}_S)^2}{M^4_{\eta} (s - M^2_R)^2 + M^2_R \Gamma^2_R} \left[ (\text{Re}\mathcal{S}_S)^2 (I^2_1 + I^2_2) + 2(\text{Im}\mathcal{S}_S)^2 I^2_3 \right], $$

where the spin of initial states is averaged and $\alpha = M^2_{h} / M^2_{\eta}$ and $\beta = m^2_e / M^2_{\eta}$. Notice that Eq.(6.1) has the s-wave contributions because the coupling $\mathcal{S}_S$ is complex. The mass parameter $M_R$ is a mass eigenvalue of the Higgs mass matrix $M^2_h$ which is satisfied the resonance mass relation $M_R \simeq 2M^1_h$.

$^{1}$We take account of physical pole ($\Delta > 0$). Unphysical pole analysis is studied in detail in [18].
\[ \eta + S + S_0, \eta_S^0, \eta_S^+ \]  

\[ e^-, \nu_\tau \]

\[ e^+, \bar{\nu}_\tau \]

Figure 6: The main process for the positron excess from the DM annihilation. The s-channel diagram induces the Breit-Wigner enhancement.

\[ \Gamma_R \] is the decay width to \[ n_S n_S \] and \[ \Delta = 1 - 4 M_S^2 / M_R^2 \]. The resonance particle \[ R \] is described in terms of the linear combination of flavor eigenstate fields as

\[ R = O_{R1} \rho_1 + O_{R2} \rho_2 + O_{R3} \rho_S + O_{R4} \sigma_1 + O_{R5} \sigma_2 + O_{R6} \sigma_S + O_{R7} \phi. \]  

(6.7)

There are the other contributions to the \[ n_S \] annihilation cross section such as \( t, u \)-channel in Fig. 2 or the interference contributions between \( t, u \)-channel and \( s \)-channel. However all we have to consider is the contribution of Eq. (6.1) because this is dominant at the present Universe that the DM relative velocity is \( v \sim 10^{-3} \). One finds that the flavor symmetry remarkably fixes the final states to be positron/electron in our scenario.

The thermally averaged annihilation cross section \( \langle \sigma_2 v \rangle \) is defined as

\[ \langle \sigma_2 v \rangle = \frac{\int d^3 p_1 d^3 p_2 (\sigma_2 v) f_1^{eq} f_2^{eq}}{\int d^3 p_1 d^3 p_2 f_1^{eq} f_2^{eq}}, \]  

(6.8)

where \( p_i \) is the momentum of initial particle \( i \) and \( f_i^{eq} = e^{-E_i / T} \) is the Maxwell-Boltzmann distribution function. If we can expand the annihilation cross section in terms of \( v^2 \) as \( \sigma_2 v = a_2 + b_2 v^2 \), we can calculate it easily as \( \langle \sigma_2 v \rangle = a_2 + b_2 x / x \), where \( x = M_S / T \). Although such a naive treatment is not justified when the annihilation cross section has a resonance point, an approximate estimation is obtained as follows if the condition \( \gamma_R \ll \Delta \) is satisfied

\[ \langle \sigma_2 v \rangle \approx \frac{|h_3|^2 (O_{R7})^2 \text{Im} \mathbb{G}_S}{(4\pi)^4 \pi^{1/2}} \left[ (\text{Re} \mathbb{G}_S)^2 \left( \frac{I_1^2}{2} + \frac{I_2^2}{2} \right) + (\text{Im} \mathbb{G}_S)^2 I_3^2 \right] \frac{m_e \sqrt{\Delta \gamma_R}}{M_\eta^4} \frac{1}{x^{3/2}} e^{-x \Delta}, \]  

(6.9)

\footnote{Although there are other decay channels like \( \phi_S \phi_S \), we assume that decay to \( n_S n_S \) is dominant to lead to the Breit-Wigner enhancement.}

\footnote{We assume that \( n_I \) in the loop does not contribute to the positron production because \( n_I \) can produce the tauon final state with no suppression, which is now forbidden by the Fermi-LAT \( \gamma \)-ray experiment \cite{4}. Such a condition can be realized in our model by controlling the coupling \( \mathcal{G}_1 \) to be small.}
where \( \gamma_R = \Gamma_R/M_R \). Since \( \Gamma_R \) is proportional to \( \sqrt{\Delta} \), one might suspect that large annihilation cross section is obtained under the condition \( \gamma_R/\Delta < 0.1 \). We will discuss this point below.

We define the boost factor \( BF \) as

\[
BF \equiv \frac{\langle \sigma_2 v \rangle}{3.0 \times 10^{-9} \text{[GeV}^{-2}\text{]}},
\]

and contours of the boost factor are shown in Fig.7 where the red regions satisfy the condition \( \gamma_R/\Delta < 0.1 \) and \( M_\eta = 500 \text{ GeV} \) is taken as a typical example. The degree of the fine tuning is smaller(larger) if the smaller(larger) \( M_\eta \) value is taken because the thermally averaged cross section \( \langle \sigma_2 v \rangle \) is inversely proportional to \( M_\eta^4 \). One finds that a large boost factor is obtained through the Breit-Wigner enhancement from Fig.7 if the parameters satisfy \( \Delta \lesssim 10^{-13} \) and \( \text{Im}S \ll \text{Re}S \).

Under this condition, the thermally averaged annihilation cross section and the decay width of \( R \) are written as

\[
\langle \sigma_2 v \rangle \simeq \frac{16\sqrt{\pi}}{(4\pi)^3}|h_3|^4 O_{R7}^2 (\text{Re}S)^2 \left( \frac{I_1^2}{2} + \frac{I_2^2}{2} \right) \frac{m_e^2 x^{3/2} e^{-x\Delta}}{M_\eta^4},
\]

\[
\Gamma_R \simeq \frac{\text{Im}S^2}{16\pi} O_{R7}^2 M_R \sqrt{\Delta},
\]

thus one find that \( \text{Im}S \ll \text{Re}S \) is important to obtain large annihilation cross section and small decay width.

The flux of positron and electron from DM annihilation is given by \( \Phi_{e^\pm}(\epsilon) \) [28] and the positron fraction is given by

\[
\text{Positron Fraction} \equiv \frac{\Phi_{e^+}(\epsilon) + \Phi_{e^+}^{\text{sec}}(\epsilon)}{\Phi_{e^+}(\epsilon) + \Phi_{e^-}^{\text{sec}}(\epsilon) + \Phi_{e^-}(\epsilon) + \Phi_{e^-}^{\text{prim}}(\epsilon) + \Phi_{e^-}^{\text{sec}}(\epsilon)},
\]

\[
\gamma_\Delta < 0.1
\]

\[
BF = 1
\]

\[
BF = 10
\]

\[
BF = 100
\]

\[
BF = 1000
\]

\[
BF = 10000
\]
where $\Phi_{e^{\pm}}(\epsilon)$ are the contributions from DM annihilation, and the others are the background fluxes given by

$$\Phi_{e^{-}}^{\text{prim}}(\epsilon) = \frac{0.16 \epsilon^{-1.1}}{1 + 11 \epsilon^{0.9} + 3.2 \epsilon^{2.15}} \text{(GeV}^{-1}\text{cm}^{-2}\text{s}^{-1}\text{sr}^{-1}),$$

$$\Phi_{e^{-}}^{\text{sec.}}(\epsilon) = \frac{0.70 \epsilon^{0.7}}{1 + 110 \epsilon^{1.5} + 600 \epsilon^{2.9} + 580 \epsilon^{4.2}} \text{(GeV}^{-1}\text{cm}^{-2}\text{s}^{-1}\text{sr}^{-1}),$$

$$\Phi_{e^{+}}^{\text{sec.}}(\epsilon) = \frac{4.5 \epsilon^{0.7}}{1 + 650 \epsilon^{2.3} + 1500 \epsilon^{4.2}} \text{(GeV}^{-1}\text{cm}^{-2}\text{s}^{-1}\text{sr}^{-1}).$$

(6.14)

Figure 8: The positron fraction for DM annihilation($n_S$) into $e^+e^-$. The red, green and blue lines are the best fits for $M_S = 230$ GeV and $\langle \sigma v \rangle = 8.5 \times 10^{-8}$ GeV$^{-2}$, $M_S = 450$ GeV and $\langle \sigma v \rangle = 2.6 \times 10^{-7}$ GeV$^{-2}$ and $M_S = 750$ GeV and $\langle \sigma v \rangle = 6.8 \times 10^{-7}$ GeV$^{-2}$ respectively.

The direct positron fraction is plotted in Fig.8 for some fixed parameters. The $BF$ of order $10^2$ is required in all cases. This $BF$ is not large enough to fit the Fermi-LAT data [29]. Thus the constraints from diffuse gamma rays and neutrinos are not severe as long as isothermal dark matter profile is considered [4]. It might be worth mentioning that the DM mass less than $O(\text{TeV})$ is in favor of the experiment recently reported by HESS [30], if one considers the NFW profile [31].

### 6.2 Muon Flux Measurement from Super-Kamiokande

We briefly mention that the high energy neutrinos induced by DM annihilations in the earth, the sun, and the galactic center are an important signal for the indirect detection of the DM [32]. Such energetic neutrinos induce upward through-going muons from charged current interactions, which provide the most effective signatures in Super-Kamiokande (SK) [33]. Once the thermally averaged
cross section of the muon flux reaches the same order of the cross section required by the PAMELA results, it is natural to expect that such a value of cross section is close to the upper bound of the muon flux measured by SK. In fact, our model has the large cross section enhanced by the Breit-Wigner mechanism with \(\nu_\tau\) pair final state as can been seen in Fig[6]. However since the total cross section is proportional to the neutrino mass as in Eq.(6.1), the neutrino flux is extremely suppressed than positron and electron fluxes.

### 6.3 Antiproton Production from DM Annihilation

![Diagram](image)

Figure 9: The antiproton flux is mainly induced by the top-pair production through the Higgs mixing between \(\varphi\) and \(\phi_S\) in our model. Where \(T = E - M_S\) is the kinetic energy of antiproton, we set \(\langle \sigma_3 v \rangle = 3 \times 10^{-9} \text{ GeV}^{-2}\).

Finally, we briefly discuss the antiproton flux in the cosmic ray. Since our model has the quark-DM coupling through the Higgs mixing between \(\varphi\) and \(\phi_S\), as discussed in section[5], we have to verify that our antiproton flux is consistent with the antiproton experiment of PAMELA. The main source comes from the top quark pair production, and substantially bottom and charm pair production. The cross section of the \(n_S n_S \rightarrow q \bar{q}\) processes is given by

\[
\sigma_3 v(n_S n_S \rightarrow q \bar{q}) = \frac{1}{2\pi} \sqrt{1 - \frac{m_q^2}{M_S^2}} \frac{m_q^2}{v_S^2} \frac{(\text{Im} \mathcal{G}_S)^2 M_S^2}{(s - M_R^2)^2 + M_R^2 \Gamma_R^2} \times \left[ (\mathcal{O}_{R3} \mathcal{O}_{R7})^2 \left( 1 - \frac{m_q^2}{M_S^2} \right) + (\mathcal{O}_{R6} \mathcal{O}_{R7})^2 \right],
\]

(6.15)

where the index \(q\) is summed over top, bottom, and charm quark. The energy-squared \(s\) of the initial state is defined in Eq.(6.5).
The thermally averaged annihilation cross section is expressed in terms of $\Delta$, $\gamma_S$ and some couplings as

$$\langle \sigma_3 v \rangle \simeq \frac{(\text{Im} S S)}{16\pi^{1/2}} \sqrt{1 - \frac{m_q^2}{M_S^2}} \frac{m_q^2}{v_S^2 M_S^2} \left[(O_{R3} O_{R7})^2 \left(1 - \frac{m_q^2}{M_S^2}\right) + (O_{R6} O_{R7})^2\right] x^{3/2} \frac{\sqrt{\Delta}}{\gamma_R} e^{-\Delta x}. \quad (6.16)$$

The PAMELA experiment implies the positron excess, but no antiproton excess. Thus the ratio of the annihilation cross section to leptons and quarks constrains the mixing parameters between $\varphi$ and $\phi_S$. The ratio is given by

$$R \equiv \frac{\langle \sigma_3 v \rangle}{\langle \sigma_2 v \rangle} \sim \left(\frac{m_q}{m_e}\right)^2 \left(\frac{M_\eta^4}{v_S^2 M_S^2}\right) \frac{(4\pi)^4}{|h_{31}|^2} \frac{O_{R3}^2 + O_{R6}^2}{(\text{Re} S S)^2 + (\text{Im} S S)^2}, \quad (6.17)$$

where $I_1$, $I_2$ and $I_3$ are taken as $\mathcal{O}(10^{-2})$ which is evaluated by numerical analysis. If we require the boost factors for leptons and quarks to be 100 and 1 respectively, the constraint to the couplings becomes

$$\frac{O_{R3}^2 + O_{R6}^2}{(\text{Re} S S)^2 + (\text{Im} S S)^2} \lesssim \mathcal{O}(10^{-24}), \quad (6.18)$$

where we have taken the masses of $M_S = 450$ GeV and $M_\eta = 500$ GeV. We find that the mixing matrix elements $O_{R3}$ and $O_{R6}$ which appear in $\langle \sigma_3 v \rangle$ need to be suppressed by $\mathcal{O}(10^{-12})$ in order to have no antiproton excess if $S$ is $\mathcal{O}(1)$.

The flux of antiproton from DM annihilation is given in Ref.\[28\]. We plot the antiproton flux as a function of the kinetic energy of antiproton $T = E - M_S$ in Fig.9. Where we adopt $\langle \sigma_3 v \rangle = 3 \times 10^{-9}$ GeV$^{-2}$, i.e. $BF = 1$, which is required to explain the WMAP experiment, and the same set up as the positron case. The key parameters contributing to the direct detection are $O_{13}$ and $O_{17}$ which come from SM Higgs mediation, while those to the indirect detection of the antiproton are $O_{R3}$, $O_{R6}$, and $O_{R7}$ which come from resonant bosons. It suggests that both of them can be explained by independent way. Hence it is easy to find the allowed region avoiding such an enhancement as well by controlling many parameters in the Higgs sector.

### 7 Summary and Conclusions

In this paper, we have considered that two important issues of the dark matter in a non-supersymmetric extension of the radiative seesaw model with a family symmetry based on $D_6 \times \hat{Z}_2 \times Z_2$: direct detection recently reported by CDMS II and indirect detection reported by PAMELA. We suppose that the $D_6$ singlet right-handed neutrino is the promising candidate of the DM. Analyzing the
\( \mu \rightarrow e\gamma \) together with the WMAP result, we have shown the allowed region for the DM mass to be \( 230 \text{ GeV} < M_S < 750 \text{ GeV} \), within a perturbative regime. In the analysis of the direct detection experiment of CDMS II and XENON100, we have shown that the Higgs mixing between \( \varphi \) and \( \phi_S \) plays an important role in generating the quark effective couplings, and also there exist allowed region to be detected by those experiments in near future. As a result of the positron production analysis through PAMELA, a couple of remarks are in order. In the case of \( M_S = 230 \text{ GeV} \), \( M_S = 450 \text{ GeV} \) and \( M_S = 750 \text{ GeV} \), each of \( \langle \sigma v \rangle = 8.5 \times 10^{-8} \text{ GeV}^{-2} \), \( \langle \sigma v \rangle = 2.6 \times 10^{-7} \text{ GeV}^{-2} \) and \( \langle \sigma v \rangle = 6.8 \times 10^{-7} \text{ GeV}^{-2} \) is required, respectively. In all cases the required boost factor is at most \( \sim \mathcal{O}(10^2) \), which is realized by the Breit-Wigner enhancement mechanism if the suitable parameter regions are choosen. Also such boost factor is not large enough to fit the Fermi-LAT result. Thus constraints from diffuse gamma rays and neutrinos are not severe as long as isothermal dark matter profile is considered. Finally, we have investigated the antiproton flux in the cosmic ray to compare to the direct detection. We found that the constraint of the mixing from the direct detection can easily satisfy the allowed region for no antiproton excess by controlling many parameters in the Higgs sector.

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