Multi-component dark matter with magnetic moments for Fermi-LAT gamma-ray line

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We propose a model of multi-component dark matter with magnetic moments to explain the 130 GeV gamma-ray line hinted by the Fermi-LAT data. Specifically, we consider a $U(1)_{X}$ dark sector which contains two vector-like fermions besides the related gauge and Higgs fields. A very heavy messenger scalar is further introduced to construct the Yukawa couplings of the dark fermions to the heavy $[SU(2)]$-singlet leptons in the $SU(3)_{c} \times SU(2)_{L} \times SU(2)_{R} \times U(1)_{B-L}$ left-right symmetric models for universal seesaw. A heavier dark fermion with a very long lifetime can mostly decay into a lighter dark fermion and a photon at one-loop level. The dark fermions can serve as the dark matter particles benefited from their annihilations into the dark gauge and Higgs fields. In the presence of a $U(1)$ kinetic mixing, the dark matter fermions can be verified by the ongoing and forthcoming dark matter direct detection experiments.

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I. INTRODUCTION

Astronomical and cosmological observations indicate the existence of dark matter in the present universe. Many dark matter candidates have been suggested in various scenarios beyond the $SU(3)_{c} \times SU(2)_{L} \times U(1)_{Y}$ standard model (SM). The dark matter particles can be directly detected through their scatterings off the familiar nucleons and/or indirectly detected through their annihilations/decays into the SM species. Recently, the Fermi-LAT data on the cosmic gamma-ray spectrum from the Galactic center (GC) have revealed a tentative evidence for a line-like feature at an energy around 130 GeV [1–16]. Such monochromatic photon can be induced by a dark matter annihilation or decay. There have been a lot of models realizing the required dark matter annihilations or decays.

Usually, we need some particles heavier than the dark matter to mediate the significant annihilations or decays of the dark matter into the monochromatic photons at loop level since the dark matter is only allowed to have an extremely tiny electric charge [54]. On the other hand, in the $SU(3)_{c} \times SU(2)_{L} \times SU(2)_{R} \times U(1)_{B-L}$ left-right symmetric models [55] for universal seesaw [56], some heavy $[SU(2)]$-singlet leptons and quarks are allowed to have a kinetic mixing, the dark matter masses and solve the strong CP problem [59]. Such heavy fermions could take part in mediating the radiative dark matter annihilations or decays for the 130 GeV gamma-ray line. In this paper, we shall demonstrate this possibility in a multi-component dark matter scenario. Specifically, we shall introduce a $U(1)_{X}$ dark sector which contains two vector-like fermions besides the related gauge and Higgs fields. There is also a very heavy messenger scalar having the Yukawa couplings with the dark fermions and the heavy leptons. The heavier dark fermion can have a very long lifetime although it mostly decays into the lighter fermion with a photon. The dark fermions can obtain the desired dark matter relic density through their annihilations into the dark gauge and Higgs fields. As the $U(1)_{X}$ and $U(1)_{B-L}$ gauge fields are allowed to have a kinetic mixing, the dark matter fermions can be verified by the ongoing and forthcoming dark matter direct detection experiments.

II. THE MODEL

For simplicity, we will not give the full Lagrangian. Instead, we only show the kinetic mass and Yukawa terms relevant to our demonstration,

$$
\mathcal{L} = -\frac{1}{4} W_{\mu \nu}^{a} W_{\mu \nu}^{a} - \frac{1}{4} W_{\mu \nu}^{a} W_{\mu \nu}^{a} - \frac{1}{4} B_{\mu \nu} B_{\mu \nu} - \frac{1}{4} C_{\mu \nu} C_{\mu \nu}
$$

$$
- \delta \frac{1}{2} B_{\mu \nu} C_{\mu \nu} + (D_{\mu} \phi_{L})^{\dagger} D^{\mu} \phi_{L}
$$

$$
+ (D_{\mu} \phi_{R})^{\dagger} D^{\mu} \phi_{R} + (D_{\mu} \sigma) \dagger D^{\mu} \sigma + (D_{\mu} \delta) \dagger D^{\mu} \delta
$$

$$
+ i\bar{\nu}_{L} \not{D} \nu_{L} + i\bar{\nu}_{L} \not{D} \nu_{R} + i\bar{\nu}_{R} \not{D} \nu_{L} + i\bar{\nu}_{R} \not{D} \nu_{R} + i\bar{\nu}_{L} \not{D} \nu_{R} + i\bar{\nu}_{R} \not{D} \nu_{L}
$$

$$
+ 2\bar{\nu}_{L} \not{D} \nu_{R} + i\bar{\nu}_{L} \not{D} \nu_{L} + i\bar{\nu}_{R} \not{D} \nu_{R} - M_{\nu}^{2} \delta \dagger - \bar{\nu}_{L} M_{\nu} \nu_{R} - \bar{\nu}_{L} M_{\nu} \nu_{L} - \bar{\nu}_{R} M_{\nu} \nu_{R} - \bar{\nu}_{R} M_{\nu} \nu_{L}
$$

$$
- \bar{\nu}_{L} \not{D} \nu_{R} + \bar{\nu}_{R} \not{D} \nu_{L} - \bar{\nu}_{L} \not{D} \nu_{L} + \bar{\nu}_{R} \not{D} \nu_{R} + \bar{\nu}_{L} \not{D} \nu_{R} - \bar{\nu}_{R} \not{D} \nu_{L} + \bar{\nu}_{R} \not{D} \nu_{R} + \bar{\nu}_{R} \not{D} \nu_{L} + \text{H.c.}
$$

(1)

Here $W_{\mu}^{a}$, $W_{\nu}^{a}$, $B$ and $C$ are the gauge fields associated with the $SU(2)_{L}$, $SU(2)_{R}$, $U(1)_{B-L}$ and $U(1)_{X}$ gauge groups, respectively. The Higgs scalars

$$
\phi_{L}(\pm 1, 0) = \begin{bmatrix} \phi_{L}^{+} \\ \phi_{L}^{0} \end{bmatrix}, \quad \phi_{R}(\pm 1, 0) = \begin{bmatrix} \phi_{R}^{+} \\ \phi_{R}^{0} \end{bmatrix}, \quad \sigma(0, +2) (2)
$$

are an $SU(2)_{L}$ doublet, an $SU(2)_{R}$ doublet and an $SU(2)$ singlet, respectively. Here and thereafter the first and second numbers in parentheses are the $U(1)_{B-L}$ charges $B - L$ and the $U(1)_{X}$ charges $X$. The messenger scalar $\delta$
is an $SU(2)$ singlet and carries both of the $U(1)_{B-L}$ and $U(1)_Y$ charges,
\[
\delta(2, -\frac{2}{3}). \tag{3}
\]

Among the fermions, the $[SU(2)]$-doublet quarks $q_{L,R}$, the $[SU(2)]$-doublet leptons $l_{L,R}$, and the $[SU(2)]$-singlet quarks $D_{L,R}$ and $U_{L,R}$ as well as the $[SU(2)]$-singlet leptons $E_{L,R}$ carry the $U(1)_{B-L}$ charges,
\[
q_L(\frac{1}{3}, 0) = \begin{bmatrix} u_L \\ d_L \end{bmatrix}, \quad q_R(\frac{1}{3}, 0) = \begin{bmatrix} u_R \\ d_R \end{bmatrix},
\]
\[
l_L(-1, 0) = \begin{bmatrix} l'_L \\ e_L \end{bmatrix}, \quad l_R(-1, 0) = \begin{bmatrix} l'_R \\ e_R \end{bmatrix},
\]
\[
D_{L,R}(\frac{2}{3}, 0), \quad U_{L,R}(\frac{4}{3}, 0), \quad E_{L,R}(-2, 0), \tag{4}
\]

while the $[SU(2)]$-singlet dark fermions $\chi_{L,R}$ carry the $U(1)_Y$ charges,
\[
\chi_{L,R}(0, -\frac{2}{3}). \tag{5}
\]

The covariant derivatives are
\[
D_\mu = \partial_\mu - ig_X \frac{X}{2} C_\mu - ig_{B-L} \frac{B - L}{2} B_\mu - igI_3 LW^3_{\mu},
\]
\[
-igR I_{3R} W^3_{R\mu} + ..., \tag{6}
\]

where we have only written down the diagonal components of the $SU(2)$ gauge fields ($W^3_L, W^3_R$) which can mix with the other $U(1)$ gauge fields ($B, C$).

### A. Symmetry breaking

The $[SU(2)]$-doublet Higgs scalars $\phi_R$ and $\phi_L$ are responsible for breaking the left-right symmetry down to the electroweak symmetry and then the electromagnetic symmetry, i.e.
\[
SU(2)_L \times SU(2)_R \times U(1)_{B-L} \quad \downarrow \langle \phi_R \rangle
\]
\[
SU(2)_L \times U(1)_Y \quad \downarrow \langle \phi_L \rangle
\]
\[
U(1)_{em} \quad \downarrow \langle \phi_{em} \rangle. \tag{7}
\]

where $\langle \phi_R \rangle$ and $\langle \phi_L \rangle$ are the vacuum expectation values (VEVs),
\[
\langle \phi_L \rangle = \langle \phi^0_L \rangle = \frac{1}{\sqrt{2}} v_L \quad (v_L \simeq 246 \text{ GeV}),
\]
\[
\langle \phi_R \rangle = \langle \phi^0_R \rangle = \frac{1}{\sqrt{2}} v_R. \tag{8}
\]

As for the dark symmetry $U(1)_X$, it will be broken when the $[SU(2)]$-singlet Higgs scalar $\sigma$ develops its VEV,
\[
\langle \sigma \rangle = \frac{1}{\sqrt{2}} v_X. \tag{9}
\]

Roughly, the above symmetry breakings take place at the temperatures $T = O(\langle \phi_R \rangle), O(\langle \phi_L \rangle)$ and $O(\langle \sigma \rangle)$, respectively.

### B. Fermions

The charged fermions have the masses,
\[
\mathcal{L} \ni -[\bar{d}_L \bar{T}_L] \begin{bmatrix} 0 & \frac{1}{\sqrt{2}} y_D^D v_L \\ \frac{1}{\sqrt{2}} y_D^D v_L & M_D \end{bmatrix} \begin{bmatrix} d_R \\ D_R \end{bmatrix}
\]
\[
- [\bar{u}_L \bar{U}_L] \begin{bmatrix} 0 & \frac{1}{\sqrt{2}} y_U^U v_R \\ \frac{1}{\sqrt{2}} y_U^U v_R & M_U \end{bmatrix} \begin{bmatrix} u_R \\ U_R \end{bmatrix}
\]
\[
- [\bar{e}_L \bar{E}_L] \begin{bmatrix} 0 & \frac{1}{\sqrt{2}} y_E^E v_R \\ \frac{1}{\sqrt{2}} y_E^E v_R & M_E \end{bmatrix} \begin{bmatrix} e_R \\ E_R \end{bmatrix} + \text{H.c.}, \tag{10}
\]

which can be block diagonalized by
\[
\mathcal{L} \ni -\bar{d}_L \left( -\frac{g_D^L}{2 M_D} v_L y_D^D \right) d_R - \bar{T}_L M_D D_R
\]
\[
- \bar{u}_L \left( -\frac{g_U^L}{2 M_U} v_R y_U^U \right) u_R - \bar{U}_L M_U U_R
\]
\[
- \bar{e}_L \left( -\frac{g_E^L}{2 M_E} v_R y_E^E \right) e_R - \bar{E}_L M_E E_R + \text{H.c.}, \tag{11}
\]

Remarkably, we have the heavy charged fermions besides the SM charged fermions in this universal seesaw scenario.

In the following we will work in the base where the mass matrices of the dark fermions $\chi$ and the heavy charged fermions $F = D, U, E$ are diagonal and real, i.e.
\[
m_\chi = \text{diag}(m_{\chi_1}, m_{\chi_2}),
\]
\[
M_F = \text{diag}(M_{F_1}, M_{F_2}, M_{F_3}), \tag{12}
\]

and then define the vector-like fermions:
\[
\chi_i = \chi_{Li} + \chi_{Ri}, \quad F_a = F_{La} + F_{Ra}. \tag{13}
\]

Moreover, we will assume $m_{\chi_1} \leq m_{\chi_2}$ and $M_{F_1} \leq M_{F_2} \leq M_{F_3}$ without loss of generality. Note the left-right symmetry breaking scale and the heavy fermion masses should be large enough to escape from the experimental constraint. For example, the right-handed charged gauge boson should be heavier than a few TeV [70]. In the
In the following we will focus on the case that the orthogonal fields $Z_B$, $Z_L$ and $Z_X$ approximate to the mass eigenstates for $m_{Z_R}^2 \gg m_{Z_L}^2 \gg m_{Z_X}^2$ and $\epsilon \ll 1$. In this case, the quasi-mass-eigenstate $Z_L$ is identified to the SM $Z$ boson.

### III. Dark Gauge Boson Decay

The dark gauge filed $C$ which is mostly the quasi-mass-eigenstate $Z_X$ can couple to the SM fermions besides the dark fermions,

$$\mathcal{L} \supset -\frac{1}{6} \xi g_{B-L} \bar{d} \gamma^\nu d Z_{\chi \mu} - \frac{1}{6} \xi g_{B-L} \bar{u} \gamma^\nu u Z_{\chi \mu} + \frac{1}{2} \xi g_{B-L} \bar{e} \gamma^\nu e Z_{\chi \mu} + \frac{1}{2} \xi g_{B-L} \bar{\nu} \gamma^\nu \nu Z_{\chi \mu} - \frac{1}{3} \sqrt{1 - \epsilon^2} g_X \chi \gamma^\mu \chi Z_{\chi \mu}.$$ (19)

Therefore, if the dark gauge boson $Z_X$ is heavy enough, it can decay into the SM fermion pairs $ff$. For $2m_f < m_{Z_X} < 2m_\chi$, the decay width should be

$$\Gamma_{Z_X} = \sum_f \Gamma_{Z_X \rightarrow ff} \sim \sum_f \frac{N_c \epsilon^2 g_{B-L}^2}{12\pi} \left( \frac{B - L}{2} \right)^2 m_{Z_X} \left( 1 + \frac{m_f^2}{m_{Z_X}^2} \right) \times \sqrt{1 - \frac{4m_f^2}{m_{Z_X}^2}},$$ (20)

where $(N_c, B - L) = (3, \frac{1}{3})$ for a quark and $(N_c, B - L) = (1, -1)$ for a lepton. We then find

$$\tau_{Z_X} \simeq \left( \frac{1.34 \times 10^{-11}}{\epsilon} \right)^2 \left( \frac{0.428}{g_{B-L}} \right)^2 \left( \frac{500~\text{MeV}}{m_{Z_X}} \right) \sec$$ (21)

So, the dark gauge boson $Z_X$ with a mass $m_{Z_X} = 500$ MeV can have a lifetime shorter than 1 second if we take $\epsilon > 1.34 \times 10^{-11}$. Currently, the measurement on the muon magnetic moment constrains $\epsilon^2 c_{B-L}^2 < 2 \times 10^{-4}$ for $m_{Z_X} = 500$ MeV [74].

### IV. Dark Matter Relic Density

The dark fermions $\chi_{1,2}$ can annihilate into the dark gauge field $C$ and the dark Higgs field $\sigma$. The heavier dark fermion $\chi_2$ can also annihilate into the lighter dark fermion $\chi_1$. The relevant diagrams are shown in Fig. 1. If the annihilations freeze out before the $U(1)_X$ symmetry breaking at the temperature $T = O(\sigma)^2$, the present work, we will take $v_R \sim M_{F,1,2,3} = O(10^3 \text{TeV})$ to give a numerical example.

The present model can accommodate a discrete parity symmetry to solve the strong CP problem without an axion [59, 71]. Furthermore, the neutral neutrinos can have a two-loop induced Dirac mass matrix proportional to the SM charged lepton mass matrix [72]. One can introduce more fermions or scalars to generate the desired neutrino masses and mixing.

### C. Gauge fields

We can remove the kinetic mixing between the $U(1)_X$ and $U(1)_B-L$ gauge fields by making a non-unitary transformation [72],

$$B_\mu = \tilde{B}_\mu - \frac{\epsilon}{\sqrt{1 - \epsilon^2}} \tilde{C}_\mu = \tilde{B}_\mu - \xi \tilde{C}_\mu,$$

$$C_\mu = \frac{1}{\sqrt{1 - \epsilon^2}} \tilde{C}_\mu,$$ (14)

and then define the orthogonal fields,

$$A_\mu = W^3_{3\mu} s_W + (W^3_{R\mu} s_R + \tilde{B}_\mu c_R) c_W,$$ (15a)

$$Z_{L\mu} = W^3_{L\mu} c_W - (W^3_{R\mu} s_R + \tilde{B}_\mu c_R) s_W,$$ (15b)

$$Z_{R\mu} = W^3_{R\mu} c_R - \tilde{B}_\mu s_R,$$ (15c)

$$Z_{X\mu} = \tilde{C}_\mu,$$ (15d)

with

$$s_R = \sin \theta_R, \quad c_R = \cos \theta_R \quad \text{for} \quad t_R = \tan \theta_R = \frac{g_{B-L}}{g_R},$$

$$s_W = \sin \theta_W, \quad c_W = \cos \theta_W \quad \text{for} \quad t_W = \tan \theta_W = \frac{g_{B-L} g_R/g}{\sqrt{g_{B-L}^2 + g_R^2}} = \frac{g'}{g}. $$ (16)

Here $g$ and $g'$ are the SM gauge couplings with $g \simeq 0.653$ and $g' \simeq 0.358$ while $\theta_W$ is the Weinberg angle $s_W^2 \equiv 0.231$. If a parity symmetry is imposed, we can determine the unknown gauge couplings $g_R$ and $g_{B-L}$ by

$$g_R = g \simeq 0.653, \quad g_{B-L} = \frac{g' g}{\sqrt{g^2 - g'^2}} \simeq 0.428. $$ (17)

Among the orthogonal fields (15), $A$ is the massless photon $\gamma$, while $Z_L$, $Z_R$ and $Z_X$ have the mass terms as below,

$$\mathcal{L} \supset \frac{1}{2} m_{Z_R}^2 (Z_{R\mu} + \xi Z_{X\mu} s_R) (Z_{R\mu} + \xi Z_{X\mu} s_R) + \frac{1}{2} m_{Z_L}^2 (Z_{L\mu} + Z_{R\mu} t_R s_W + \xi Z_{X\mu} t_R s_W)$$

$$\times (Z_{L\mu} + Z_{R\mu} t_R s_W + \xi Z_{X\mu} t_R s_W) + \frac{1}{2} m_{Z_X}^2 Z_{X\mu} Z_{X\mu}$$

with

$$m_{Z_R} = \frac{1}{2} \cos \theta_R g_{B-L},$$

$$m_{Z_L} = \frac{1}{2} \cos \theta_W g_{B-L} \simeq 91 \text{GeV},$$ (18)

$$m_{Z_X} = \frac{1}{\sqrt{1 - \epsilon^2}} g_X \psi_X,$$
FIG. 1: The dark fermion annihilations. Here $\chi_{i,j}$ ($i \neq j$) denotes the dark matter fermions, $C$ is the $U(1)_X$ gauge field, while $\sigma$ is the Higgs scalar for breaking the $U(1)_X$ symmetry.

The freeze-out temperature determined by $\left[75\right]$ is the Higgs scalar for breaking the $U(1)_X$ symmetry.

Annihilation cross sections should be dominated by

$$\sigma_{\chi_1} = \langle \sigma_{\chi_1 \bar{\chi}_1 \to CC} v_{\text{rel}} \rangle + \langle \sigma_{\chi_1 \bar{\chi}_1 \to \sigma \bar{\sigma}} v_{\text{rel}} \rangle$$

and

$$\sigma_{\chi_2} = \langle \sigma_{\chi_2 \bar{\chi}_2 \to CC} v_{\text{rel}} \rangle + \langle \sigma_{\chi_2 \bar{\chi}_2 \to \sigma_1 \bar{\sigma}_1} v_{\text{rel}} \rangle$$

Here $v_{\text{rel}}$ is the relative velocity between the two annihilating particles in their center-of-mass frame.

As we will clarify later, the heavier dark fermion $\chi_2$ has a very long lifetime and hence contributes to the dark matter relic density together with the lighter and stable dark fermion $\chi_1$. The relic density of the dark fermion $\chi_1$ can be calculated by $\left[75\right]$ to obtain

$$\frac{\sigma_{\chi_1}}{m_{\chi_1}} \simeq 2.45 \times 10^{-7} \text{ GeV}^{-2}, \quad T_1 \simeq m_{\chi_1}/25.3,$$

and

$$\frac{\sigma_{\chi_2}}{m_{\chi_2}} \simeq 1.87 \times 10^{-9} \text{ GeV}^{-2}, \quad T_2 \simeq m_{\chi_2}/23.0,$$

The total dark matter relic density thus should be

$$\Omega_{\chi_1,\chi_2} \chi^2 \simeq 0.001,$$

which is well consistent with the observations $\left[76\right]$. In the above numerical estimation, the heavier dark fermion $\chi_2$ rather than the lighter dark fermion $\chi_1$ dominates the dark matter relic density, i.e.

$$\frac{\Omega_{\chi_1,\chi_2} \chi^2}{\Omega_{\text{DM}} \chi^2} \simeq 0.99 \%, \quad \frac{\Omega_{\chi_2} \chi^2}{\Omega_{\text{DM}} \chi^2} \simeq 99.1 \%,$$

as a result of their hierarchical masses $m_{\chi_2}^2 \gg m_{\chi_1}^2$. Actually, it is easy to see

$$\frac{\Omega_{\chi_1,\chi_2} \chi^2}{\Omega_{\text{DM}} \chi^2} \leq 50\% \leq \frac{\Omega_{\chi_2} \chi^2}{\Omega_{\text{DM}} \chi^2} \text{ for } m_{\chi_1} \leq m_{\chi_2}.$$

V. DARK MATTER DECAY

The dark matter fermions $\chi_{1,2}$ can couple to the $U(1)_{B-L}$ and $U(1)_X$ field strength tensors ($B_{\mu\nu}, C_{\mu\nu}$) at one-loop level. We show the relevant diagrams in Fig. 2. For $m_{\chi_{1,2}} < M_{E_6}, M_{\delta}$, the effective interactions should
where the couplings $\lambda_{ij}$ and $\kappa_{ij}$ are given by

$$\lambda_{ij} = \frac{g_{B-L}}{32\pi^2} \langle f_L^j \rangle_{ia} \langle f_R^j \rangle_{aj} \frac{1}{M_{E_a}} \text{F}_B \left( \frac{M_e^2}{M_{E_a}^2} \right) \quad \text{with}$$

$$\text{F}_B(x) = \frac{1}{1-x} + \frac{x}{(1-x)^2} \ln x,$$

and

$$\kappa_{ij} = \frac{g_X}{96\pi^2} \langle f_L^i \rangle_{ia} \langle f_R^j \rangle_{aj} \frac{1}{M_{E_a}} \text{F}_C \left( \frac{M_e^2}{M_{E_a}^2} \right) \quad \text{with}$$

$$\text{F}_C(x) = -\frac{1+x}{2(1-x)^2} - \frac{x}{(1-x)^3} \ln x.$$

Therefore, the heavier dark matter fermion $\chi_2$ can decay into the lighter dark matter fermion $\chi_1$ and a gauge boson $(\gamma, Z, Z_X$ or $Z_R)$ as long as the kinematics is allowed. Provided that

$$m_{Z_R} \gg m_{\chi_2} - m_{\chi_1} > m_{Z}, m_{Z_X},$$

we can obtain the decay widths,

$$\Gamma_{\chi_2 \to \chi_1 \gamma} = \Gamma_{\tilde{\chi}_2 \to \tilde{\chi}_1 \gamma}$$

$$\approx \frac{c_W^2 c_Y^2}{2\pi} \langle \lambda_{12} \rangle^2 m_{\chi_2}^3 F_2 \left( \frac{m_{\chi_1}^2}{m_{\chi_2}^2}, \frac{m_Z^2}{m_{\chi_2}^2} \right),$$

with

$$F_2(x, y) = \left[ (1-x)^2 - \frac{1}{2} y (1+x+y) \right] \times \sqrt{(1-x-y)^2 - 4xy}.$$
and then
\[ \Gamma_{\chi_2 \to \chi_1 \gamma} = \sum_{\alpha, \beta} \Gamma_{\chi_2 \to \chi_1 e^\alpha e^\beta} \approx \frac{1}{2g_{\pi}^2} (R_{11} R_{22} + R_{11} L_{22} + R_{22} L_{11}) + L_{11} L_{22} \frac{m_2^2}{M_\chi^2} F_3 \left( \frac{m_2^2}{m_{\chi_2}^2} \right), \] (38)

with
\[ R_{ii} = \left( f_L \frac{\nu_R}{\sqrt{2} M_E} y_E \frac{\nu_R}{\sqrt{2} M_E} f_L \right)_{ii}, \]
\[ L_{ii} = \left( f_R \frac{\nu_L}{\sqrt{2} M_E} y_E \frac{\nu_L}{\sqrt{2} M_E} f_R \right)_{ii}, \] (39)

and
\[ F_3(x) = \frac{1}{12} (1 - 8x + 8x^3 - x^4) + x^2 \ln x. \] (40)

The heavier dark matter fermion $\chi_2$ can have a very long lifetime. For example, by inputting
\[ g_{B-L} = 0.428, \quad g_X = 0.592, \quad \epsilon = 10^{-7}, \]
\[ f_L = f_R = \text{diag} \left\{ \sqrt{2} m_e, \sqrt{2} m_e, \sqrt{2} m_e, \sqrt{2} m_e \right\}, \]
\[ M_{E_{1,2,3}} = \frac{1}{\sqrt{2}} \nu_R = 900 \text{ TeV}, \quad M_\delta = 10^{16} \text{ GeV}, \]
\[ m_\chi_2 = 262 \text{ GeV}, \quad m_{\chi_1} = 20 \text{ GeV}, \] (41)
we can obtain
\[ \Gamma_{\chi_2 \to \chi_1 \gamma} = 5.03 \times 10^{-53} \text{ GeV}, \]
\[ \Gamma_{\chi_2 \to \chi_1 z} = 1.23 \times 10^{-53} \text{ GeV}, \]
\[ \Gamma_{\chi_2 \to \chi_1 z} = 2.22 \times 10^{-57} \text{ GeV}, \]
\[ \Gamma_{\chi_2 \to \chi_1 l^+ l^-} = 6.20 \times 10^{-66} \text{ GeV}, \] (42)

and then
\[ \text{Br}_{\chi_2 \to \chi_1 \gamma} = \frac{\Gamma_{\chi_2 \to \chi_1 \gamma}}{\Gamma_{\chi_2 \to \chi_1 \gamma} + \Gamma_{\chi_2 \to \chi_1 z}} \approx 80.3\%, \]
\[ \tau_{\chi_2} = \frac{1}{\Gamma_{\chi_2}} \approx 1.05 \times 10^{28} \text{ sec}, \] (43)

with $\Gamma_{\chi_2}$ and $\Gamma_{\bar{\chi}_2}$ being the total decay width,
\[ \Gamma_{\chi_2} = \Gamma_{\bar{\chi}_2} = \Gamma_{\chi_2 \to \chi_1 \gamma} + \Gamma_{\chi_2 \to \chi_1 z} + \Gamma_{\chi_2 \to \chi_1 z} + \Gamma_{\chi_2 \to \chi_1 l^+ l^-} \approx \Gamma_{\chi_2 \to \chi_1 \gamma} + \Gamma_{\chi_2 \to \chi_1 z} \]
\[ = 9.26 \times 10^{-53} \text{ GeV}, \] (44)

Furthermore, the photons from the decays $\chi_2 \to \chi_1 \gamma$ and $\bar{\chi}_2 \to \bar{\chi}_1 \gamma$ have the determined energy
\[ E_\gamma = \frac{m_2^2 - m_{\chi_1}^2}{2m_{\chi_2}} \approx 130 \text{ GeV}. \] (45)

The gamma-ray flux from the dark matter decays $\chi_2 \to \chi_1 \gamma$ and $\bar{\chi}_2 \to \bar{\chi}_1 \gamma$ can be written as
\[ \frac{d\Phi}{dE\Omega} = \frac{\Omega_{\chi_2+\bar{\chi}_2} h^2}{\Omega_{\Delta M} h^2} \frac{\text{Br}_{\chi_2 \to \chi_1 \gamma}}{4\pi m_{\chi_2}^2} \int_{l.o.s.} ds \rho(s, \psi) \frac{dN}{dE}, \] (46)

where $dN/dE$ is the differential gamma spectrum per dark matter decay with $E$ being the gamma-energy, $r(s, \psi) = (r_\odot^2 + s^2 - 2r_\odot s \cos \psi)^{1/2}$ is the coordinate centered on the GC with $s$ being the distance from the Sun along the line-of-sight (l.o.s.), $r_\odot$ being the distance from the Sun to the GC and $\psi$ being the angle between the direction of observation in the sky and the GC, $\rho(r(s, \psi))$ is the dark matter density profile. In Ref. [27], the authors have shown that the decay of a single-component dark matter fermion into a neutrino and a photon can explain the gamma-ray line in the Fermi-LAT data [27], i.e.
\[ \frac{d\Phi}{dE d\Omega} = \frac{\text{Br}_{\nu \gamma}}{4\pi m_{\chi_2}^2} \int_{l.o.s.} ds \rho(s, \psi) \frac{dN}{dE}. \] (47)

Here
\[ m_{\chi} \approx 2E_\gamma \approx 260 \text{ GeV}, \] (48)
is the dark matter mass, $\tau_\chi$ is the dark matter lifetime, $\text{Br}_{\nu \gamma}$ is the branching ratio of the decay modes $\chi \to \nu \gamma$. By comparing Eqs. [15-16] and [17-18], we can take
\[ \frac{m_2^2 - m_{\chi_1}^2}{2m_{\chi_2}} = m_{\chi_1}, \quad \text{Br}_{\chi_2 \to \chi_1 \gamma} = \text{Br}_{\nu \gamma}, \]
\[ \tau_{\chi_2} = \frac{\Omega_{\chi_2+\bar{\chi}_2} h^2}{\Omega_{\Delta M} h^2} m_{\chi_2} \tau_\chi, \] (49)
to account for the fitting results [27]. For example, we can obtain the Fermi-LAT gamma-ray line by inputting
\[ m_{\chi_2} = 262 \text{ GeV}, \quad m_{\chi_1} = 20 \text{ GeV}, \]
\[ \frac{\Omega_{\chi_2+\bar{\chi}_2} h^2}{\Omega_{\Delta M} h^2} = 99.1\%, \quad \text{Br}_{\chi_2 \to \chi_1 \gamma} = 80.3\%, \]
\[ \tau_{\chi_2} = 1.05 \times 10^{28} \text{ sec}. \] (50)

VI. DARK MATTER SCATTERING AND SELF-INTERACTION

As shown in Fig. 4, the dark matter fermions $\chi_{1,2}$ can scatter off the nucleon $N$ through the kinetic mixing between the $U(1)_X$ and $U(1)_{B-L}$ gauge fields. The elastic
scattering cross section is computed by
\[
\sigma_{\chi_i N} = \frac{\sigma_{\chi_i N \to \chi_i N}}{\sigma_{\chi_i N \to \chi_i N}} = \sigma_{\chi_i N \to \chi_i N} \simeq \frac{\sqrt{g_{\chi i} g_{B-L}}}{324 \pi} \frac{\mu_r^2}{m_{Z_X}^2} \sigma_{\chi_i N \to \chi_i N} = 3.9 \times 10^{-45} \text{cm}^2 \left( \frac{\epsilon}{10^{-7}} \right)^2 \frac{(g_X/0.592)^2}{(500 \text{ MeV})^4} \left( \frac{\mu_r}{1 \text{ GeV}} \right)^2.
\]

Here we have defined the reduced mass,
\[
\mu_r = \frac{m_{\chi_i}^2 m_N}{m_{\chi_i} + m_N} \simeq m_N \simeq 1 \text{ GeV} \text{ for } m_{\chi_i} \gg m_N. \tag{52}
\]

Such scattering can be measured by the dark matter direct detection experiments. The event rate per unit time per nucleon should be
\[
\frac{R_{\chi_i}}{m_{\chi_i}} \approx \frac{\Omega_{\chi_i + \chi_i}^2 h^2}{2 \pi} \rho_\odot \sigma_{\chi_i N} \tag{53}
\]

with \(\rho_\odot\) being the local dark matter density. Currently, the measured experimental rate is given in the single-component dark matter hypothesis,
\[
\frac{R_{\chi_i}}{m_{\chi_i}} \approx \frac{\rho_\odot}{m_{\chi_i}} \sigma_{\chi_i N} \tag{54}
\]

where \(\sigma_{\chi_i N}\) denotes the cross section of a single-component dark matter \(\chi\) with the mass \(m_{\chi_i}\) scattering off the nucleon \(N\). The XENON10 and XENON100 experiments \cite{Experiments_1, Experiments_2} have stringently put an upper bound \(\sigma_{\exp}\) on the cross section \(\sigma_{\chi_i N}\) for a given mass \(m_{\chi_i}\). Therefore, we should constrain
\[
\sigma_{\chi_i N} < \frac{\Omega_{\chi_i + \chi_i}^2 h^2}{2 \pi} \sigma_{\chi_i N} \exp. \tag{55}
\]

For a proper parameter choice such as
\[
m_{\chi_1} = 20 \text{ GeV}, \quad m_{\chi_2} = 262 \text{ GeV}, \quad \sigma_{\chi_{1} N} = \sigma_{\chi_{2} N} = 3.9 \times 10^{-45} \text{cm}^2, \quad \frac{\Omega_{\chi_1 + \chi_1}^2 h^2}{\Omega_{\chi_2 + \chi_2}^2 h^2} \simeq 3.9 \times 10^{-43} \text{cm}^3 \left( \frac{g_X}{0.592} \right)^4 \tag{56}
\]

the dark matter fermions \(\chi_{1,2}\) can be verified by the ongoing and forthcoming dark matter direct detection experiments.

Furthermore, the dark matter fermions \(\chi_{1,2}\) can have a self-interaction as shown in Fig. \ref{fig:kip}. For \(m_{\chi_i} \gg m_{Z_X}\), the self-interacting cross section should be
\[
\frac{\sigma_{\chi_i \chi_i \to \chi_i \chi_i}}{m_{\chi_i}} = \frac{\sigma_{\chi_i \chi_i \to \chi_i \chi_i}}{m_{\chi_i}} \simeq \frac{g_X^4 m_{\chi_i}}{628 \pi m_{Z_X}^2}, \tag{57}
\]

from which we read
\[
\frac{\sigma_{\chi_1 \chi_1 \to \chi_1 \chi_1}}{m_{\chi_1}} \simeq 3 \times 10^{-43} \text{cm}^3 \left( \frac{g_X}{0.592} \right)^4 \left( \frac{\chi_{\chi_1}}{2 \text{ GeV}} \right)^4 \left( \frac{\chi_{\chi_1}}{500 \text{ MeV}} \right)^4, \tag{58}
\]

For a single-component dark matter with the mass \(m\), its self-interacting cross section \(\sigma\) has an upper bound \(\sigma/m < 4.4 \times 10^{-42} \text{cm}^3\). \cite{Xenon_2019}. We hence should require
\[
\frac{\Omega_{\chi_i + \chi_i}^2 h^2 \sigma_{\chi_i \chi_i \to \chi_i \chi_i}}{\Omega_{\chi_i + \chi_i}^2 h^2 m_{\chi_i}} = \frac{\Omega_{\chi_1 + \chi_1}^2 h^2 \sigma_{\chi_1 \chi_1 \to \chi_1 \chi_1}}{\Omega_{\chi_1 + \chi_1}^2 h^2 m_{\chi_1}} < 4.4 \times 10^{-42} \text{cm}^3. \tag{59}
\]

Clearly, the parameter choice \cite{Xenon_2019} can satisfy the limit since we have \(\frac{\Omega_{\chi_i + \chi_i}^2 h^2}{\Omega_{\chi_i + \chi_i}^2 h^2} < 1.\)
FIG. 5: The dark matter self-interactions $\chi_i\chi_i \rightarrow \chi_i\chi_i$ and $\chi_i\bar{\chi}_i \rightarrow \chi_i\bar{\chi}_i$. Here, $\chi_i$ denotes the dark matter fermions while $C$ is the $U(1)_X$ gauge field. For simplicity, we don’t show the processes $\bar{\chi}_i\bar{\chi}_i \rightarrow \bar{\chi}_i\bar{\chi}_i$.

### VII. SUMMARY

It was suggested that the 130 GeV gamma-ray line hinted by the Fermi-LAT data could be understood by the dark matter annihilation or decay into monochromatic photons. We hence propose a multi-component dark matter model, where a heavier dark matter fermion mostly decays into a lighter dark matter fermion and a photon, to explain the Fermi-LAT gamma-ray line. In our model, the neutral dark matter fermions have a highly suppressed magnetic moment at one-loop level because of their Yukawa couplings to a charged scalar and three non-SM leptons. The new scalar besides the dark matter fermions is gauged by a $U(1)_X$ symmetry which will be spontaneously broken below the GeV scale. As for the non-SM leptons, they play an essential role for generating the SM lepton masses in the $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left-right symmetric models for the universal seesaw scenario where the strong CP problem can be solved without an axion. The dark matter fermions can obtain a thermally produced relic density through their annihilations into the $U(1)_X$ gauge and Higgs fields. The kinetic mixing between the $U(1)_X$ and $U(1)_{B-L}$ gauge fields can result in a testable dark matter scattering.

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