Effective Field Theory of Majorana Dark Matter

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Abstract

We revisit thermal Majorana dark matter from the viewpoint of minimal effective field theory. In this framework, analytic results for dark matter annihilation into standard model background are derived. The dark matter parameter space subject to the latest LUX, PandaX-II and Xenon-1T limits is presented in a model-independent way. Applications to singlet-doublet, MSSM, and MSSM with vector-like lepton model are worked out.
I. INTRODUCTION

In the viewpoint of standard cosmology, dark matter (DM) is a neutral particle beyond the standard model (SM). Nevertheless, neither in particle astrophysical or collider experiments it has been observed to date. According to its electrically neutral property, in this paper we focus on DM being a Majorana fermion. This case covers a lot of well known models such as neutralino [1], singlet-doublet [2–9], Higgs-portal [10–13] and Z-portal [14–19] DM.

For a Majorana DM, the effective Lagrangian at the weak scale is described by,

\[ \mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{dark}}(\chi, \cdots), \] (1)

where the SM Lagrangian \( \mathcal{L}_{\text{SM}} \) contains interactions between DM mediators \( h \) and \( Z \) and SM particles,

\[ \mathcal{L}_{\text{SM}} \supset \frac{h}{v_{\text{EW}}} \left( \sum_f m_f \bar{\psi}_f \psi_f + 2m_w W^+ W^- + m_Z^2 Z^\mu Z_\mu \right) + Z_\mu \sum_f \bar{\psi}_f \gamma^\mu (g_V - g_A \gamma_5) \psi_f + \cdots \] (2)

with

\[ g_V = \frac{g}{\cos \theta_W} \left( \frac{3f}{2} - Q_f \sin^2 \theta_W \right) \quad \text{and} \quad g_A = \frac{g}{\cos \theta_W} \left( \frac{3f}{2} \right) . \] (3)

Here, the weak scale \( v_{\text{EW}} = 246 \text{ GeV} \), \( g \simeq 0.65 \) is the gauge coupling of \( SU(2)_L \) group, and \( \theta_W \) denotes the weak mixing angle. Note, \( Q_f \) is the electric charge, and \( T_{3f} = \pm \frac{1}{2} \) for up-type and down-type fermion, respectively.

Furthermore, \( \mathcal{L}_{\text{dark}} \) generally contains the following interaction terms between (4-component) Majorana DM and SM mediators,

\[ \mathcal{L}_{\text{dark}}(\chi, \cdots) \supset c_h h \bar{\chi} \chi + c_z Z^\mu \bar{\chi} \gamma_\mu \gamma_5 \chi + \cdots \] (4)

Note that a vector coupling between DM and \( Z \) boson is forbidden. The \( c_h \) and \( c_z \) interaction terms constitute the minimal framework from the perspective of effective field theory where higher-dimensional operators [20] responsible for obvious gauge invariance of Eq.(4) should be included. We refer it as the “minimal” effective field theory.

New physical particles beyond the minimal effective field theory impose diverse effects. If they are decoupled, their net effects are recorded in parameter \( c_h \) and \( c_z \) in Eq.(4). Conversely, if not, they should be included in Eq.(2) or Eq.(4), which either play the role of new mediator between DM and SM sector or contribute to new DM annihilation final states as long as kinetically allowed. In the former case, Lagrangians in Eq.(2) and Eq.(4) serve as the SM background of DM annihilation and DM scattering cross sections. In the later one, new particles with mass of order the weak scale yield a few new Feynman diagrams for these cross sections. When the number of new particles is large, the numerical treatment is more viable than the analytic one.

However, it is only the analytic treatment which can clearly show us the ingredients as required to fit future signatures of DM direct detection, which is the main starting point for this study. The paper is organized as follows. Sec.II is devoted to analytically derive DM annihilation into the SM background in the minimal effective field theory. We will compare...
FIG. 1: Feynman diagrams for DM annihilation into SM final states.

| channel | $a$ | $b$ |
|---------|-----|-----|
| $f \bar{f}$ | $a_{ff}$ | $b_{ff}$ |
| $ZZ$ | $a_{zz}$ | $b_{zz}$ |
| $W^+W^-$ | $b_{ww}$ |
| $hh$ | $b_{hh}$ |
| $Zh$ | $a_{zh}$ | $b_{zh}$ |

TABLE I: Coefficients of $\sigma v_\chi$ expansion in individual SM final state.

our results with those arising from numerical calculations. In Sec.III we show the parameter space subject to the latest DM direct detection limits in a model-independent way. In Sec.IV we apply our method to study the singlet-doublet, minimal supersymmetric standard model (MSSM) and MSSM with vector-like leptons with decoupling mass spectrum. Finally, we conclude in Sec.V.

II. RELIC DENSITY

According to the effective Lagrangian in Eq. (1) DM annihilates into SM final states $f \bar{f}$, $ZZ$, $WW$, $Zh$ and $hh$ through SM mediator $h$ and $Z$. In order to calculate DM relic density we firstly derive the thermally averaged cross section $\langle \sigma v_\chi \rangle$, responsible for which Feynman diagrams are shown in Fig. 1. Feynman diagrams similar to Fig. 1 have already been addressed in more complicated context such as neutralino DM [21]. However, completely analytic expression for annihilation cross section is viable only in simplified situations such as the minimal framework discussed above.

The DM annihilation cross section times DM relative velocity $v_\chi$ can be expanded in the standard way,

$$\sigma v_\chi = a + bv_\chi^2 + \mathcal{O}(v_\chi^4). \quad (5)$$
In Table I we introduce contributions to coefficients $a$ and $b$ for individual SM final state in Fig. 1. Direct evaluation of them yields

$$a_{ff} = \frac{2N_c e^2 r_{\mu\mu} g_A^2 m_f^2}{\pi m_Z^4},$$

$$a_{zz} = \frac{4e^4 r_{\mu\mu} (m^2 - m_Z^2)}{\pi (m_Z^2 - 2m^2)^2},$$

$$a_{zh} = \frac{c^2 r_{\tau\tau} m_Z^2}{64\pi^2 v_{EW}^4 m_Z^2},$$

and

$$b_{ff} = \frac{N_c e^2 m_Z^2 r_{\mu\mu} (m^2 - m_Z^2)}{2\pi v_{EW}^2 (m_h^2 - 4m^2)^2} + \frac{N_c e^2 g_A^2 m_f^2 (5m_f^2 - 4m^2)}{4\pi v_{EW} m_Z^2 m^2} + \frac{N_c e^2 r_{\mu\mu} (m^2 (g_V^2 - 2g_A^2) + 2m^2 (g_V^2 + g_A^2))}{3\pi (m_Z^2 - 4m^2)^2},$$

$$b_{ww} = \frac{c^2 r_{\mu\mu} g^2 \cos^2 \theta_W (17m_W^2 m^4 + 16m_W^4 m^2 - 3m_W^6 + 4m^6)}{6\pi v_{EW} (m_h^2 - 4m^2)^2} + \frac{c^2 r_{\mu\mu} g^2 \cos^2 \theta_W (17m_W^2 m^4 + 16m_W^4 m^2 - 3m_W^6 + 4m^6)}{6\pi v_{EW} (m_h^2 - 4m^2)^2} + \frac{c^2 r_{\tau\tau} (m^2 - 2m_Z m^2)}{3\pi v_{EW} (m_h^2 - 4m^2)^2} (m_Z^2 - 2m_Z m^2)^2,$$

$$b_{zz} = \frac{c^2 r_{\mu\mu} (-9m^2 m_Z^4 + 12m_Z^2 m^4 - 8m^6 + 2m_Z^6)}{8\pi v_{EW} (m_h^2 - 4m^2)^2} + \frac{c^2 r_{\mu\mu} (-9m^2 m_Z^4 + 12m_Z^2 m^4 - 8m^6 + 2m_Z^6)}{8\pi v_{EW} (m_h^2 - 4m^2)^2} + \frac{c^2 r_{\mu\mu} (m^2 - 2m_Z m^2)}{3\pi v_{EW} (m_h^2 - 4m^2)^2} (m_Z^2 - 2m_Z m^2)^2,$$

$$b_{hh} = \frac{c^2 r_{\tau\tau} (m_h^2 - 4m^2)^2}{32\pi v_{EW}^2 (m_h^2 - 4m^2)^2} + \frac{c^2 r_{\tau\tau} (m_h^2 - 4m^2)^2 (2m^2 - 5m_Z^2)}{2\pi v_{EW} (m_h^2 - 4m^2)^2} + \frac{c^2 r_{\tau\tau} (m_h^2 - 4m^2)^2 (2m^2 - 5m_Z^2)}{2\pi v_{EW} (m_h^2 - 4m^2)^2} (m_Z^2 - 2m^2)^2 + \frac{c^2 r_{\tau\tau} (m_h^2 - 4m^2)^2 (2m^2 - 5m_Z^2)}{2\pi v_{EW} (m_h^2 - 4m^2)^2} (m_Z^2 - 2m^2)^2,$$

$$b_{zh} = \frac{c^2 r_{\tau\tau} (m_h^2 - 4m^2)^2}{768\pi v_{EW}^2 (m_h^2 - 4m^2)^2} \left[ (4m_Z^2 (5m_h^2 + 9m_h^2) - 2m_Z^2 (5m_h^2 + 74m_h^2 m_h^2 + 14m_h^2)) \right. + \left. 6m_Z^2 (m_h^4 - 5m_h^2 m_h^2 + 4m_h^4) - 192m_h^4 (m_h^4 - 5m_h^2 m_h^2 + 4m_h^4) - 10m_Z^8 \right] + \frac{c^2 r_{\tau\tau} (m_h^2 - 4m^2)^2 (m_h^2 - 4m^2 + 3m_h^2)}{128\pi v_{EW} (m_h^2 - 4m_Z^2)^2} \left[ (2m_Z^2 (m_h^2 - 9m_h^2) - 2m_Z^2 (m_h^2 - 4m_h^2)^2) \right. - \left. m_Z^4 (m_h^4 + 14m_h^2 m_h^2 - 104m_h^4) + 2m_Z^2 (m_h^4 + 8m_h^2 m_h^2 - 48m_h^4) - 10m_Z^8 \right] + \frac{c^2 r_{\tau\tau} (m_h^2 - 4m^2)^2 (m_h^2 - 4m_Z^2)}{768\pi v_{EW}^2 (m_h^2 - 4m^2 + 3m_h^2)} \left[ m_Z^{10} + 2m_Z^6 (3m_h^4 + 16m_h^4) + 4m_Z^8 (m_h^2 - 2m_h^2) - 4m_Z^4 (m_h^2 - 4m_h^2)^2 (m_h^2 + 10m_h^2) + 4m_Z^2 (m_h^4 - 4m_h^2)^4 + m_Z^2 (m_h^2 - 4m_h^2)^2 (m_h^2 + 8m_h^2 m_h^2 + 80m_h^4) \right].$$

where $N_c = 1(3)$ for SM lepton (quark) and $m_h$ refers to DM mass. Functional $r_{ij}$ and $r_{ij}$
is defined as
\[ r_{ij} = \sqrt{1 - m_i^2/m_j^2}, \]
\[ r_{\chi ij} = \sqrt{m_i^4 - 2m_i^2(m_j^2 + 4m_\chi^2) + (m_j^2 - 4m_\chi^2)^2/m_\chi^2}, \]
respectively.

A few comments are in order regarding our results. At first, in the case \( c_h \to 0 \) all as in Eq.(6)-Eq.(7) coincide with results of Z portal \[15, 19\], but \( b_{ff} \) and \( b_{zz} \) in \[19\] are both two times of that in Eq.(9) and Eq.(11), respectively. Secondly, in the case \( c_z \to 0 \) all \( a_s \) in Eq.(6)-Eq.(8) disappear as the same as in the Higgs portal, and our \( b_{ff} \) in Eq.(9) and \( b_{zz} \) and \( b_{hh} \) (the \( c_h^2 \)-term) is in agreement with the results of \[13\] and \[10\], respectively. Thirdly, when both \( c_z \) and \( c_h \) are non-zero, interference effects appear in \( b_{zz} \) and \( b_{zh} \), which are explicitly shown for the first time. These interference effects can be ignored except in some specific DM mass range between \( m_z \) and \( m_h \), where it is not small relative to the other contributions. Finally, we have also included the SM Higgs self interaction contribution to \( b_{hh} \) in Eq.(12). We verified that our results agree with the numerical calculation in terms of MicrOMEGAs \[24\] with at most 10\% – 15\% difference on numerical value of DM relic density.

III. DIRECT DETECTION

The interactions in Eq.(4) yield both spin-dependent (SD) and spin-independent (SI) effective couplings between DM and SM nucleons. In particular, Yukawa coupling constant \( c_h \) and \( c_z \) controls SI and SD scattering cross section, respectively, which are given by \[1, 22\],
\[ \sigma_{SI} \simeq c_h^2 \times (2.11 \times 10^3 \text{zb}), \]
\[ \sigma_{SD}^p \simeq c_z^2 \times (1.17 \times 10^9 \text{zb}), \]
\[ \sigma_{SD}^n \simeq c_z^2 \times (8.97 \times 10^8 \text{zb}). \]

Here nuclear form factors have been chosen as in \[23\]. The approximations to \( \sigma_{SI} \) and \( \sigma_{SD} \) are always valid for DM mass range of most interest \( m_\chi > a \) few times \( m_{p,n} \).

In Fig.2 we show the parameter space of DM relic density \( \Omega_{DM} h^2 = 0.1199 \pm 0.0027 \) \[25\] in the two-parameter plane of \( c_h \) and \( c_z \), with contours referring to representative DM masses in unit of GeV. We also draw contours of the latest PandaX-II \[26\], Xenon-1T \[27\] and LUX 2016 \[28\] limit simultaneously. Parameter regions above the solid lines or on the right hand of the dashed line are excluded, from which we find that model-independent exclusion limit for DM mass is about \( \sim 155 \) GeV. Only a small region
\[ 0 \leq |c_z| \leq 0.018, \quad 0 \leq |c_h| \leq 0.06 \]
is left for future tests. Indirect detection, which will not be discussed here, may slightly improve parameter ranges in Eq.(15). When this region is further excluded by updated experimental limits, one can draw the conclusion that either new physical particle(s) is required to appear at the weak scale, or simplified Majorana DM models are disfavored. In what follows we will discuss implication of our results in a few simplified models.
FIG. 2: Parameter space of DM relic density in the two-parameter plane of $c_h$ and $c_z$ subject to the latest PandaX-II [26] (black), Xenon-1T [27] (red), and LUX 2016 [28] (green and yellow) limit. Contours referring to representative DM masses (in unit of GeV) are drawn for clarity, which implies that the model-independent exclusion limit for DM mass is about $\sim 155$ GeV.

IV. APPLICATIONS

A. Singlet-Doublet Dark Matter

This model contains two fermion doublets $L' = (l^0, l^-)^T$, $L = (l^+, l^0)^T$ and a fermion singlet $\psi_s$. The dark sector Lagrangian $\mathcal{L}_{\text{dark}}$ reads as [2–4],

$$\mathcal{L}_{\text{dark}} = \frac{i}{2} \left( \bar{\psi}_s \sigma^\mu \partial_\mu \psi_s + \bar{L}' \sigma^\mu \partial_\mu L' + \bar{L} \sigma^\mu \partial_\mu L \right) + \left( -y_1 L'H \psi_s - y_2 \bar{L} \psi_s H + \text{H.c} \right) - \frac{m_{\psi_s}}{2} \bar{\psi}_s \psi_s - m_D L'L \right) \tag{16}$$

where $m_s$, $m_D$ and $y_{1,2}$ are mass and Yukawa coupling parameters, respectively. $H$ denotes the SM Higgs doublet. In the basis $(\psi_s, l^0, l^0)$ the symmetric mass matrix for neutral fermions is given by,

$$M_{\chi} = \begin{pmatrix} m_s \frac{y_1 v_{\text{EW}}}{\sqrt{2}} & \frac{y_2 v_{\text{EW}}}{\sqrt{2}} & 0 \\ * & 0 & m_D \\ * & * & 0 \end{pmatrix}. \tag{17}$$

This model is similar to the neutralino sector of the next-to-minimal supersymmetric model (NMSSM) when bino and wino component are both decoupled. Imposing the decoupling limit $m_D >> m_s, v_{\text{EW}}$ on the dark sector yields only a light singlet-like DM with mass $m_{\chi} \simeq m_s$. Under this limit, the effective coupling $c_{h\bar{\chi}\chi}$ and $c_{z\bar{\chi}\chi}$ reduces to, respectively [7],

$$c_{h\bar{\chi}\chi} \simeq -\frac{v_{\text{EW}}}{m_D} \left( 2 y_1 y_2 + (y_1^2 + y_2^2) \frac{m_{\chi}}{m_D} \right),$$

$$c_{z\bar{\chi}\chi} \simeq \frac{1}{2} \frac{v_{\text{EW}}}{m_D} \frac{m_Z}{m_D} \left( y_1^2 - y_2^2 \right) \left( 1 - \frac{m_{\chi}^2}{m_D^2} \right). \tag{18}$$
Note that $|c_z|$ and $|c_h|$ are both unchanged under the exchange of $y_1 \leftrightarrow y_2$. Since the parameter ranges in Eq. (15) favor larger value of $|c_h|$ relative to $|c_z|$, it implies that the product $y_1y_2$ in Eq. (18) should be at most of order $m_\chi/m_D$. Otherwise, $|c_h|$ at the crossing points with contours of DM relic density would be too large to exceed the direct detection limits as shown in Fig. 2.

In Fig. 3 we show the contours of DM mass projected to the plane of $c_h - c_z$ for $y_1 = -3$ and $y_2 = 0.1$, where the condition that DM mass $m_\chi$ should be at least an order of magnitude smaller than $m_D$ has been imposed. The crossing points with the contours of DM relic density are indeed beneath the DM direct detection limits for DM mass range between 200 GeV and 600 GeV. When the magnitude of $y_1$ is tuned to be smaller than 2, these viable crossing points disappear. Actually, small Yukawa couplings are more favored by collider constraints on DM direct production at the LHC (see, e.g. [29–32]). In this sense, combing the DM direct detection and collider limits will result in more robust prediction on the parameter space.

### B. MSSM

Now, we discuss application to MSSM with decoupling mass spectrum, in which all supersymmetric particles except the lightest neutralino are decoupled from the weak scale. The symmetric neutralino mass matrix $M_\chi$ under the gauge eigenstates $(\tilde{B}^0, \tilde{W}^0, \tilde{H}_d^0, \tilde{H}_u^0)$ is given by,

$$
M_\chi = \begin{pmatrix}
M_1 & 0 & -m_Z s_W \cos \beta & m_Z s_W \sin \beta \\
* & M_2 & m_Z c_W \cos \beta & -m_Z c_W \sin \beta \\
* & * & 0 & -\mu \\
* & * & * & 0
\end{pmatrix}.
$$

(19)

Imposing the decoupling limit on the Higgs sector and the neutralino sector by $|\mu|$, $M_1 >> M_2$, $m_Z$ simultaneously leads to a wino-like DM with mass $m_{\chi^0} \simeq M_2$ and reduced effective
coupling coefficient \( c_h \) \( c_z \) [33–35],
\[ c_h \bar{\chi} \bar{\chi} \simeq \frac{g}{4} \cos \theta_W \frac{m_Z}{\mu} \left( \frac{\mu}{m_\chi_0} + \sin 2\beta \right), \]
\[ c_z \bar{\chi} \bar{\chi} \simeq -\frac{g}{4} \cos \theta_W \frac{m_Z^2}{\mu^2} \left( 1 - \frac{m_\chi_0^2}{\mu^2} \right). \] (20)
, respectively. Instead, imposing a different decoupling limit \( |\mu|, M_2 >> M_1, m_Z \) we obtain a bino-like DM with mass \( m_{\chi_1^0} \simeq M_1 \) and
\[ c_h \bar{\chi} \bar{\chi} \simeq \frac{g}{4} \sin \theta_W \frac{m_Z}{\mu} \left( \frac{\mu}{m_\chi_0} + \sin 2\beta \right), \]
\[ c_z \bar{\chi} \bar{\chi} \simeq -\frac{g}{4} \sin \theta_W \frac{m_Z^2}{\mu^2} \left( 1 - \frac{m_\chi_0^2}{\mu^2} \right). \] (21)
Both decoupling limits yield a light chargino \( \tilde{\chi}^\pm \) with mass slightly larger than DM mass. The modification to cross sections can be ignored in the case of bino-like DM, and only a tiny correction in the case of wino-like DM arises due to annihilation to \( W^+W^− \) final state. Under the decoupling limit we have \( |c_z| < 1.0 \times 10^{-3} \) both in Eq.(20) and Eq.(21) given \( |m_{\chi_0^0}/\mu| \leq 0.1 \) and \( m_{Z}/ |\mu| \leq 0.1 \). It reveals that light bino-like DM mass is excluded, as shown by the contours of DM relic density in Fig.2. This result is consistent with that bino-like DM should have mass of order \( \sim 2.7 - 3.0 \) TeV [36].

C. MSSM with Vector-like Doublets

Analysis above can be also applied to relatively complicated Majorana DM models such as MSSM with vector-like doublets [37]. Such vector-like doublets \( (L, \bar{L}) \), together with singlet superfield \( (N) \), constitute a \( 5 \) or \( \bar{5} \) representation of \( SU(5) \), which is unique because they are allowed to appear at the TeV scale without violating the perturbative \( SU(5) \) [38] or \( SO(10) \) [39] grand unification. In this model, the neutralino sector is extended by,
\[ W = W_{\text{MSSM}} + kH_u \bar{\bar{L}}N - hH_dLN + M_L \bar{\bar{L}}L + \frac{1}{2} M_N N^2 + M_D \bar{D}D \] (22)
where \( M_{L,D,N} \) refer to vector-like mass parameters, \( k \) and \( h \) denote the Yukawa coupling constants. Similarly, a few new soft mass parameters appear in \( L_{\text{soft}} \):
\[ -L_{\text{soft}} = -L_{\text{soft}}^{\text{MSSM}} + m_L |L|^2 + m_L |\bar{L}|^2 + m_D |D|^2 + m_D |\bar{D}|^2 + m_N |N|^2 + (kA_kH_u \bar{L}N - hA_hH_dLN + H.c.) \] (23)
where \( m_{L,N,D} \) and \( A_{k,h} \) represent supersymmetry-breaking mass parameters.

From Eq.(22) and Eq.(23) one can derive the modified neutralino mass matrix \( M_\chi \) in this model. Following the definitions \( L = (E^+, \langle v \rangle + \frac{v}{\sqrt{2}})^T, \bar{L} = (\langle v^c \rangle + \frac{v^c}{\sqrt{2}}, E^-)^T \) and
FIG. 4: Contours of DM mass in green dotted lines projected to the plane of \( c_h - c_z \) for \( 0.1 \leq |k| \leq 3 \). Contours of DM relic density are in solid color lines the same as in Fig. 2. The condition that DM mass \( m_\chi \) should be at least an order of magnitude smaller than \( M_L \) has been imposed.

\[
N = \langle N \rangle + \frac{N}{\sqrt{2}},
\]

we find that in the basis \((\tilde{B}^0, \tilde{W}^0, \tilde{H}_d^0, \tilde{H}_u^0, \tilde{c}, \tilde{v}, \tilde{N})\),

\[
M_\chi = \begin{pmatrix}
M_1 & 0 & -(1-x)m_Z s_W c_\beta & (1-x)m_Z s_W s_\beta & -x m_Z s_W c_\gamma & x m_Z s_W s_\gamma & 0 \\
* & M_2 & (1-x)m_Z c_W c_\beta & -(1-x)m_Z c_W s_\beta & -x m_Z c_W c_\gamma & x m_Z c_W s_\gamma & 0 \\
* & * & 0 & -\mu & 0 & h \langle N \rangle & h \langle v \rangle \\
* & * & * & 0 & -k \langle N \rangle & 0 & -k \langle \nu^c \rangle \\
* & * & * & * & 0 & -M_L & -k v_u \\
* & * & * & * & * & 0 & h v_d \\
* & * & * & * & * & * & -M_N
\end{pmatrix},
\]

where \( s_\beta = \sin \beta, c_\beta = \cos \beta, s_\gamma = \sin \gamma, c_\gamma = \cos \gamma; xv^2_{EW}/2 = \langle v \rangle^2 + \langle \nu^c \rangle^2, (1-x)v^2_{EW}/2 = \langle v_u \rangle^2 + \langle v_u \rangle^2; \) and \( \tan \beta = \langle v_u \rangle / \langle v_u \rangle, \quad \tan \gamma = \langle v \rangle / \langle \nu^c \rangle \). In Eq. (24) there are totally twelve model parameters, which dramatically complex the situation.

It is thus meaningful to consider various decoupling limits before a numerical treatment. However, it does not mean these decoupling limits automatically agree with other experimental constraints, especially the precision tests on SM Higgs couplings. LHC data suggests that the deviations from SM expectations for Higgs couplings are small, which implies that the mixing effects \([40, 41]\) between Higgs doublets and vector-like doublets/singlet should be small as well. It can be satisfied by imposing decoupling vector-like leptons. Note, \( M_L \) below mass about \( \sim 280 \text{ GeV} \) \([42]\) has been excluded by the 8 TeV LHC data.

With decoupling vector-like leptons that impose \( v = \nu^c = 0 \), the simplified DM can be divided into following cases. Firstly, imposing \( M_{2(1)}, \mu, M_L, M_N \gg M_{1(2)}, m_Z \) results in a bino (wino)-like DM, which are similar to discussions in Sec.IVB. Secondly, imposing \( M_1, M_2, \mu, M_L \gg M_N, m_Z \) contributes to a singlet-like DM. It differs from the singlet-like DM both in the singlet-doublet model and NMSSM. Because its effective couplings differ from Eq. (18) in the sense that \( y_1 \) and \( y_2 \) therein are not longer arbitrary but related to each other by \( \tan \beta \), and unlike in the NMSSM \([43]\), where small \( \tan \beta \) is favored by the Higgs mass, the model with superpotential in Eq. (22) favors a large \( \tan \beta \) value similar to MSSM.
Under the decoupling limit, we obtain from Eq. (24)
\begin{align}
  c_h \bar{\chi} \chi &\simeq -k^2 v_{EW} \frac{m_{\chi}}{M_L M_L}, \\
  c_z \bar{\chi} \chi &\simeq -\frac{k^2 v_{EW} m_Z}{2 M_L M_L} \left(1 - \frac{m_{\chi}^2}{M_L^2}\right),
\end{align}
(25)

where $m_{\chi} \simeq M_N$.

In Fig. 4 we show the contours of DM mass in green dotted lines projected to the plane of $c_h - c_z$ with Yukawa coupling $0.1 \leq |k| \leq 3.0$. Comparing them with the contours of DM relic density shown therein in solid lines, we arrive at the conclusion that a simplified singlet-like DM with mass near the weak scale in this model is excluded.

V. CONCLUSION

In this paper, we have revisited the Majorana DM, a weakly interacting massive particle, from the viewpoint of minimal effective field theory. Unlike the Dirac-type analogy, there is no vector coupling between a Majorana DM and $Z$ boson. In this framework, there are only three parameters - DM mass, Yukawa coupling constant $c_h$ and $c_Z$. Accordingly, one can sufficiently constrain the parameter space in a model-independent way. In order to achieve this, we have analytically derived annihilation of DM into the SM background, with interference effects and SM Higgs self-interaction contributions included. The fit to the latest LUX, PandaX-II and Xenon-1T limits points to DM mass lower bound about 155 GeV.

We also address applications to the singlet-doublet, MSSM and MSSM with vector-like leptons model. In singlet-doublet model, we find that singlet-like DM with mass between 200 GeV and 600 GeV still survives in the latest direct detection limits. In the MSSM, we recover the exclusion limit on bino DM mass. Finally, in the MSSM with vector-like leptons, we show that singlet-like DM is excluded, which is different from the case of singlet-doublet model and NMSSM.

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