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Gauge-singlet dark matter in a left-right symmetric model with spontaneous *CP* violation

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We propose a dark matter (DM) scenario in an extension of a left-right symmetric model with a gauge-singlet scalar field. The gauge-singlet scalar can automatically become a DM candidate, provided that both *P* and *CP* symmetries are only broken spontaneously. Thus no extra discrete symmetries are needed to make the DM candidate stable. After constraining the model parameters from the observed relic DM density we make predictions for direct detection experiments. We show that for some parameter range, the predicted weakly interacting massive particle-nucleon elastic scattering cross section can reach the current experimental upper bound, which can be tested by the experiments in the near future.

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I. INTRODUCTION

The standard model (SM) of particle physics, although greatly successful in phenomenology, gives no explanations for parity (*P*) and *CP* violation. The observed neutrino oscillations, the large baryonumber asymmetry and large energy density from nonbaryonic dark matter (DM) in the universe are clear indications for new physics beyond the SM. In the left-right (LR) symmetric models for weak interactions [1–4], the left- and right-handed fermions are treated equally, and *P* symmetry is restored prior to the spontaneous symmetry breaking (SSB). The LR models have other advantages such as the gauge group $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ can be elegantly embedded into grand unification theories, and the right-handed neutrinos are naturally required, etc.

The LR models may also contain DM candidates. In one of the minimal versions of the LR model, which contains one Higgs bidoublet ϕ and two Higgs triplets $\Delta_{L,R}$, a Z_2 symmetry on the left triplet $\Delta_L \leftrightarrow -\Delta_L$ can be used to resolve the so-called vacuum expectation value (VEV) seesaw problem [1]. A direct consequence of this discrete symmetry is that the neutral component δ_L^0 of Δ_L can only annihilate or be produced by pairs, which makes δ_L^0 a potential DM candidate. Unfortunately, due to the fact that δ_L^0 participates $SU(2)_L$ gauge interactions which is rather strong, the weakly interacting massive particle (WIMP)-nucleon elastic scattering experiments lead to a severe constraint on the dark matter relic density which is a few order of magnitudes below the observed value [5]. Thus it cannot be a main source of DM in the universe.

Besides the phenomenological problems, from a theoretical point of view, taking the neutral components of the $SU(2)_{L,R}$ triplets as DM candidates may lead to difficulties in neutrino mass generation. The neutral components are often stabilized by some discrete symmetries such as Z_2. In the minimal LR model, although the aforementioned Z_2 symmetry $\Delta_L \leftrightarrow -\Delta_L$ is essential for stabilizing the neutral scalar δ_{L0}, it eliminates the Majorana mass terms for both left- and right-handed neutrinos due to the *P* symmetry. The same problem appears in the $SU(2)_L$ triplet extension of the SM motivated by the type-II seesaw mechanism. Namely, the Majorana mass term for the left-handed neutrino is forbidden after a Z_2 symmetry is imposed on the triplet.

Other potential DM candidates in the minimal LR model may involve the heavy right-handed neutrinos. However, they are unlikely to be stable, as the right-handed neutrinos participate in right-handed gauge interactions. Furthermore, they can decay into Higgs bosons and light leptons through Yukawa interactions. Thus in the minimal LR model there is no suitable bosonic or fermionic DM candidate. For having a realistic DM candidate, one needs to consider extending the model with more particle contents. The most economical solution could be adding one or a few gauge-singlet scalars.

Another disadvantage of the minimal LR model is that although *P* can be broken spontaneously, the *CP* symmetry has to be broken explicitly, which looks quite unnatural. The reason is that without large fine-tuning on the Higgs potential, the condition for spontaneous *CP* violation (SCPV) cannot be satisfied [1,6,7]. Furthermore, in the minimal LR model with SCPV the predicted *CP* phase angle $\sin 2\beta \sim 0.1$ in the unitarity triangle of the Cabibbo-Kobayashi-Maskawa (CKM) matrix is far below the experimentally measured value of $\sin 2\beta = 0.671 \pm 0.024$ from the two B-factories [8]. The minimal LR model also suffers from strong phenomenological constraints from low energy flavor-changing-neutral-current (FCNC) processes, especially the neutral kaon mixings which push the masses of the right-handed gauge bosons and some neutral Higgs bosons above the TeV scale [9–15].

Motivated by the requirement of both spontaneous *P* and *CP* violation, we have recently discussed an extension of

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the minimal LR model with two Higgs bidoublets [16–18]. In this two-Higgs-bidoublet LR model (2HBLR), the additional Higgs bidoublet may modify the Higgs potential such that the fine-tuning problem can be avoided. The extra Higgs bidoublet can also change the interferences among different contributions in the box-diagrams in the neutral meson mixings, and lower the bounds for right-handed gauge boson masses to be below the TeV scale [16,17].

Note that the spontaneous P and CP violation in the LR models can also be useful for DM model-building. Before the SSB, the Lagrangian for the particle interactions prohibits the P-odd and CP-odd interactions, which may prevent the decays of the particles with odd CP parity. These particles can remain stable even after the SSB, provided that they do not develop VEVs and do not couple to the symmetry breaking sector. The simplest case would be that there is a gauge-singlet scalar field with odd CP parity, and has a vanishing VEV.

In this work we discuss this possibility by considering an extension of the 2HBLR with a gauge-singlet complex field S which plays the role of DM candidate, and the stability of DM is purely protected by the discrete P and CP symmetries. This model distinguishes itself from the previous gauge-singlet models (see, e.g. [19–26]) in that no *ad hoc* discrete symmetry of Z_2 type is introduced. This possibility has not been emphasized before in the literature, simply because most of the popular models such as SM and MSSM violate P and C maximally. This simple model shows that the DM may be connected to the fundamental symmetries of the quantum field theory. Recently, it is also noticed that the custodial symmetry of the gauge interaction can also be used to stabilize the DM candidate [27]. We calculate in this model the cross sections for the DM annihilation and the elastic scattering with the nucleus. The results show that for a large parameter space the DM relic density can be reproduced. The correlation between the DM annihilation and the elastic scattering off the nucleus depends on the Higgs and Yukawa sector of the model, and can be quite different from the ordinary gauge-singlet model which, in some limits, has a simple one-to-one correspondence. In general, the predictions for the direct detection experiments can be significantly larger and can even reach the current experimental upper bound for large Yukawa couplings.

This paper is organized as follows: In Sec. II, we outline the main features of the model. In Sec. III, we discuss the parameter space, and give the formulas for the main processes relevant to the DM annihilation and the elastic scattering off the nucleus in a simplified case where one Higgs bidoublet decouples from the theory. The case in which both Higgs bidoublets are active is discussed in Sec. IV. We finally conclude in Sec. V.

II. THE LR SYMMETRIC MODEL WITH A GAUGE-SINGLET

We begin with a LR model in the gauge group $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ which contains two Higgs bidoublets $\phi(2, 2^*, 0)$, $\chi(2, 2^*, 0)$, a left (right)-handed Higgs triplet $\Delta_{L(R)}$ $(3(1), 1(3), 2)$, and a gauge-singlet $S(0, 0, 0)$ with the following flavor contents

$$\phi = \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix}, \qquad \chi = \begin{pmatrix} \chi_1^0 & \chi_1^+ \\ \chi_2^- & \chi_2^0 \end{pmatrix}, \tag{1}$$
$$\Delta_{L,R} = \begin{pmatrix} \delta_{L,R}^+/\sqrt{2} & \delta_{L,R}^{++} \\ \delta_{L,R}^0 & -\delta_{L,R}^+/\sqrt{2} \end{pmatrix},$$
$$S = \frac{1}{\sqrt{2}}(S_\sigma + iS_D).$$

The introduction of Higgs bidoublet χ is to overcome the problem of fine-tuning in generating the SCPV in the minimal LR model and to relax the severe low energy phenomenological constraints [16–18]. Under the P- and CP-transformation, these fields transform as

$$\begin{array}{ccc} & P & CP \\ \phi & \phi^\dagger & \phi^* \\ \chi & \chi^\dagger & \chi^* \\ \Delta_{L(R)} & \Delta_{R(L)} & \Delta^*_{L(R)} \\ S & S & S^* \end{array} \tag{2}$$

We shall require P and CP invariance of the Lagrangian, which strongly restricts the structure of the Higgs potential. For instance, for the terms involving the ϕ and $\Delta_{L,R}$ fields the most general potential is given by

$$\begin{aligned} -\mathcal{V}_{\phi\Delta} = & -\mu_1^2 \operatorname{Tr}(\phi^\dagger\phi) - \mu_2^2[\operatorname{Tr}(\tilde{\phi}^\dagger\phi) + \operatorname{Tr}(\tilde{\phi}\phi^\dagger)] - \mu_3^2[\operatorname{Tr}(\Delta_L\Delta_L^\dagger) + \operatorname{Tr}(\Delta_R\Delta_R^\dagger)] + \lambda_1[\operatorname{Tr}(\phi^\dagger\phi)]^2 \\ & + \lambda_2^2\{[\operatorname{Tr}(\tilde{\phi}^\dagger\phi)]^2 + [\operatorname{Tr}(\tilde{\phi}\phi^\dagger)]^2\} + \lambda_3[\operatorname{Tr}(\tilde{\phi}^\dagger\phi)\operatorname{Tr}(\tilde{\phi}\phi^\dagger)] + \lambda_4\{\operatorname{Tr}(\phi^\dagger\phi)[\operatorname{Tr}(\tilde{\phi}^\dagger\phi) + \operatorname{Tr}(\tilde{\phi}\phi^\dagger)]\} \\ & + \rho_1\{[\operatorname{Tr}(\Delta_L\Delta_L^\dagger)]^2 + [\operatorname{Tr}(\Delta_R\Delta_R^\dagger)]^2\} + \rho_2[\operatorname{Tr}(\Delta_L\Delta_L)\operatorname{Tr}(\Delta_L^\dagger\Delta_L^\dagger) + \operatorname{Tr}(\Delta_R\Delta_R)\operatorname{Tr}(\Delta_R^\dagger\Delta_R^\dagger)] \\ & + \rho_3[\operatorname{Tr}(\Delta_L\Delta_L^\dagger)\operatorname{Tr}(\Delta_R\Delta_R^\dagger)] + \rho_4[\operatorname{Tr}(\Delta_L\Delta_L)\operatorname{Tr}(\Delta_R^\dagger\Delta_R^\dagger) + \operatorname{Tr}(\Delta_L^\dagger\Delta_L^\dagger)\operatorname{Tr}(\Delta_R\Delta_R)] + \alpha_1 \operatorname{Tr}(\phi^\dagger\phi)\operatorname{Tr}(\Delta_L\Delta_L^\dagger) \\ & + \operatorname{Tr}(\Delta_R\Delta_R^\dagger)] + \alpha_2 \operatorname{Tr}[(\tilde{\phi}^\dagger\phi) + (\tilde{\phi}\phi^\dagger)]\operatorname{Tr}[(\Delta_L\Delta_L^\dagger) + (\Delta_R\Delta_R^\dagger)] + \alpha_3[\operatorname{Tr}(\phi\phi^\dagger\Delta_L\Delta_L^\dagger) + \operatorname{Tr}(\phi^\dagger\phi\Delta_R\Delta_R^\dagger)] \\ & + \beta_1[\operatorname{Tr}(\phi\Delta_R\phi^\dagger\Delta_L^\dagger) + \operatorname{Tr}(\phi^\dagger\Delta_L\phi\Delta_R^\dagger)] + \beta_2[\operatorname{Tr}(\tilde{\phi}\Delta_R\phi^\dagger\Delta_L^\dagger) + \operatorname{Tr}(\tilde{\phi}^\dagger\Delta_L\phi\Delta_R^\dagger)] \\ & + \beta_3[\operatorname{Tr}(\phi\Delta_R\tilde{\phi}^\dagger\Delta_L^\dagger) + \operatorname{Tr}(\phi^\dagger\Delta_L\tilde{\phi}\Delta_R^\dagger)], \end{aligned} \tag{3}$$

where the coefficients μ_i, λ_i, ρ_i, α_i, and β_i in the potential are all real as all the terms are self-Hermitian. The Higgs potential $\mathcal{V}_{\chi\Delta}$ involving χ field can be obtained by the replacement $\chi \leftrightarrow \phi$ in Eq. (3). The mixing term $\mathcal{V}_{\chi\phi\Delta}$ can be obtained by replacing one of ϕ by χ in all the possible ways in Eq. (3).

In order to simplify the discussion, in this section we shall first consider a simple case in which the bidoublet χ does not mix significantly with other fields. In this case the model is reduced to the minimal LR model plus a gauge-singlet, which already contains the main features of the complete model. We postpone the discussions on the χ contributions into Sec. IV. The most general Higgs potential involving the singlet field S is given by

$$-\mathcal{V}_S = \frac{1}{\sqrt{2}}\tilde{\mu}_0^3(S+S^*) - \tilde{\mu}_S^2 SS^* - \frac{1}{4}\tilde{\mu}_\sigma^2(S+S^*)^2 + \sqrt{2}\tilde{\mu}_{\sigma S}(S+S^*)SS^* + \frac{1}{6\sqrt{2}}\tilde{\mu}_{3\sigma}(S+S^*)^3 + \tilde{\lambda}_S(SS^*)^2 - \frac{1}{4}\tilde{\lambda}_{\sigma S}(S+S^*)^2 SS^* - \frac{1}{16}\tilde{\lambda}_\sigma(S+S^*)^4 + \sum_{i=1}^{3}\left[-\frac{1}{\sqrt{2}}\tilde{\mu}_{i,\sigma}(S+S^*) + \tilde{\lambda}_{i,S}SS^* - \frac{1}{4}\tilde{\lambda}_{i,\sigma}(S+S^*)^2\right]O_i, \quad (4)$$

where

$$O_1 = \mathrm{Tr}(\phi^\dagger\phi),$$
$$O_2 = \mathrm{Tr}(\phi^\dagger\tilde{\phi} + \tilde{\phi}^\dagger\phi) \quad \text{and} \quad O_3 = \mathrm{Tr}(\Delta_L^\dagger\Delta_L + \Delta_R^\dagger\Delta_R). \quad (5)$$

Note that it only involves combinations of $(S+S^*)$ and SS^*. The terms proportional to odd powers of $(S-S^*)$ are absent in the singlet self-interactions as they are P-even but C-odd. Furthermore, they cannot mix with the Higgs multiplets in O_i because the three independent gauge-invariant combinations $O_i (i = 1, \ldots, 3)$ in Eq. (5) are both P- and C-even. Other possible Higgs multiplet combinations such as $\mathrm{Tr}(\phi^\dagger\tilde{\phi} - \tilde{\phi}^\dagger\phi)$ and $\mathrm{Tr}(\Delta_L^\dagger\Delta_L - \Delta_R^\dagger\Delta_R)$ are P-odd, thus cannot couple to S. The terms proportional to even powers of $(S-S^*)$ can be rewritten in terms of $(S+S^*)^2$ and SS^*. We have checked that the P- and CP-transformation rules for S defined Eq. (2) are actually the only possible way for the implementation of the DM candidate. For future convenience, we rewrite $\mathcal{V}_S$ in terms of the component fields S_σ and S_D.

$$-\mathcal{V}_S \equiv \mu_0^3 S_\sigma - \frac{1}{2}\mu_\sigma^2 S_\sigma^2 - \frac{1}{2}\mu_D^2 S_D^2 + \frac{1}{3}\mu_{3\sigma}S_\sigma^3 + \mu_{\sigma D}S_\sigma S_D^2 + \frac{1}{4}\lambda_\sigma S_\sigma^4 + \frac{1}{4}\lambda_D S_D^4 + \frac{1}{2}\lambda_{\sigma D}S_\sigma^2 S_D^2 + \sum_{i=1}^{3}\left(-\mu_{i,\sigma}S_\sigma + \frac{\lambda_{i,\sigma}}{2}S_\sigma^2 + \frac{\lambda_{i,D}}{2}S_D^2\right)O_i, \quad (6)$$

with the redefined coefficients

$$\mu_0 = \tilde{\mu}_0, \quad \mu_\sigma^2 = \tilde{\mu}_S^2 + \tilde{\mu}_\sigma^2, \quad \mu_D^2 = \tilde{\mu}_S^2, \quad \mu_{3\sigma} = \tilde{\mu}_{3\sigma} + 3\tilde{\mu}_{\sigma S}, \quad \mu_{\sigma D} = \tilde{\mu}_{\sigma S}, \quad \lambda_\sigma = \tilde{\lambda}_S - \tilde{\lambda}_{\sigma S} - \tilde{\lambda}_\sigma,$$
$$\lambda_D = \tilde{\lambda}_S, \quad \lambda_{\sigma D} = \tilde{\lambda}_S - \frac{1}{2}\tilde{\lambda}_{\sigma S}, \quad \mu_{i,\sigma} = \tilde{\mu}_{i,\sigma}, \quad \lambda_{i,\sigma} = \tilde{\lambda}_{i,S} - \tilde{\lambda}_{i,\sigma}, \quad \lambda_{i,D} = \tilde{\lambda}_{i,S}. \quad (7)$$

It follows from Eq. (6) that S_D can only be produced by pairs, thus is a potential dark matter candidate. After the SSB, the Higgs multiplets obtain nonzero VEVs

$$\langle\phi_{1,2}^0\rangle = \frac{\kappa_{1,2}}{\sqrt{2}} \quad \text{and} \quad \langle\delta_{L,R}^0\rangle = \frac{v_{L,R}}{\sqrt{2}}, \quad (8)$$

where κ_1, κ_2, v_L, and v_R are in general complex, and $\kappa \equiv \sqrt{|\kappa_1|^2 + |\kappa_2|^2} \approx 246$ GeV represents the electroweak symmetry breaking (EWSB) scale. The value of v_R sets the scale of LR symmetry breaking, which is directly linked to the right-handed gauge boson masses.

With the extra contributions from $\mathcal{V}_S$ to the whole Higgs potential $\mathcal{V} \equiv \mathcal{V}_{\phi\Delta} + \mathcal{V}_S$, one needs to redo the minimization with respect to ϕ and $\Delta_{L,R}$. However, from Eqs. (3) and (4) it follows that the minimization conditions for ϕ and $\Delta_{L(R)}$ remain to have the same form as that in the minimal LR model. This is because the mixing introduced by the singlet S only changes the overall coefficients μ_1, μ_2 and μ_3 of the ϕ and $\Delta_{L,R}$ potential term in Eq. (3). Hence the mass matrix of the Higgs multiplet ϕ and $\Delta_{L,R}$ remains the same as that in the minimal LR model in Refs. [1,28], which also indicates that the additional potential term $\mathcal{V}_S$ in Eq. (6) does not help in resolving the fine-tuning problem. The fine-tuning can only be relaxed by introducing another bidoublet χ. From the minimization condition for the singlet S_σ, one can eliminate one parameter μ_σ

$$\mu_\sigma^2 = \lambda_\sigma v_\sigma^2 - \frac{\mu_0^3}{v_\sigma} + \mu_{3\sigma}v_\sigma - \sum_i\left(\frac{\mu_{i,\sigma}}{v_\sigma}\langle O_i\rangle - \lambda_{i,\sigma}\langle O_i\rangle\right), \quad (9)$$

where $v_\sigma \equiv \langle S_\sigma \rangle$ is the VEV of S_σ. In order to ensure the stability of the dark matter candidate S_D, we require that S_D does not obtain a nonzero VEV, $\langle S_D \rangle = 0$, namely CP is not broken by the singlet fields. It follows that after the SSB, although P and CP are both broken, there is a Z_2 type of discrete symmetry on S_D remaining in the gauge-singlet sector. The discrete symmetry is induced from the original CP symmetry.

In the limit that $v_L \simeq 0$ and $\kappa_2 \ll \kappa_1$, which comes from the phenomenology of neutrino masses and neutral meson mixings, the mass eigenstates for the Higgs bidoublet and triplets approximately coincide with the corresponding flavor eigenstates. The mass terms for the Higgs bosons and gauge bosons are listed in Table I. There is only one light SM-like Higgs h^0 from the real part of ϕ_1^0, the mass of all the other scalars are set by v_R which can be very heavy.

The mass terms for S_D and S_σ are given by

$$
\begin{aligned}
M_D^2 &= \left(\tilde\lambda_\sigma + \frac{1}{2}\tilde\lambda_{\sigma S}\right) v_\sigma^2 - (\tilde\mu_{\sigma S} + \tilde\mu_{3\sigma}) v_\sigma + \tilde\mu_\sigma^2 \\
&\quad + \frac{\tilde\mu_0^3}{v_\sigma} + \frac{\tilde\mu_{i,\sigma}}{v_\sigma}\langle O_i\rangle, \\
M_\sigma^2 &= 2\lambda_\sigma v_\sigma^2 + (\tilde\mu_{3\sigma} + 3\tilde\mu_{\sigma S}) v_\sigma + \frac{\tilde\mu_0^3}{v_\sigma} + \frac{\tilde\mu_{i,\sigma}}{v_\sigma}\langle O_i\rangle, \\
M^2_{\sigma\phi_1^{0r}} &= \kappa(-\tilde\mu_{1,\sigma} + \lambda_{1,\sigma} v_\sigma), \\
M^2_{\sigma\phi_2^{0r}} &= 2\kappa(-\tilde\mu_{2,\sigma} + \lambda_{2,\sigma} v_\sigma), \\
M^2_{\sigma\delta_R^{0r}} &= v_R(-\tilde\mu_{3,\sigma} + \lambda_{3,\sigma} v_\sigma),
\end{aligned}
\tag{10}
$$

where $M^2_{\sigma\phi_1^{0r}}$, $M^2_{\sigma\phi_2^{0r}}$, and $M^2_{\sigma\delta_R^{0r}}$ denote the mixing between singlet S_σ and the other three neutral Higgs bosons. From the Lagrangian in Eq. (6) one can easily obtain the interaction terms among the scalars. Some of the relevant cubic and quartic scalar interactions are listed in Table II.

III. DM IN THE LR SYMMETRIC MODEL

There are a number of free parameters in the model such as the coefficients in the potentials and the VEV for S_σ. As shown in Eq. (10), the mass of S_D is related to two energy scales v_σ and v_R since $\langle O_3 \rangle \sim v_R^2$. In the minimal LR model with the spontaneous CP violation, the VEV v_R of the right-handed Higgs triplet Δ_R is subject to strong constraints from the K, B meson mixings [8,9,11–15,29] as well as low energy electroweak interactions [30–32]. The kaon mass difference and the indirect CP violation quantity ϵ_K set a bound for v_R around 10 TeV [30,33,34]. For a successful cold DM candidate S_D, its mass should be roughly between 10 GeV and a few TeV with annihilation cross section of approximately weak strength. For simplicity here we consider a case in which v_σ is heavy $v_\sigma \sim v_R \sim 10$ TeV, and before the SSB the Lagrangian has an approximate global $U(1)$ symmetry on S, i.e. under $S \to e^{i\delta} S$, which suppresses some of the parameters, namely

$$
\tilde\mu_0, \tilde\mu_\sigma, \tilde\mu_{\sigma S}, \tilde\mu_{3\sigma}, \tilde\mu_{i,\sigma} \ll v_\sigma, \qquad \tilde\lambda_{\sigma S}, \tilde\lambda_\sigma, \tilde\lambda_{i,\sigma} \ll 1,
\tag{11}
$$

which leads to a relatively light S_D in comparison with v_R, as it is the would-be Goldstone boson in the limit of exact global $U(1)$ symmetry. For a light S_D with $M_D \lesssim \mathcal{O}(1\ \text{TeV})$, without significant fine-tuning, one needs $\tilde\mu_{3,\sigma}/v_\sigma \lesssim 0.01$ from Eq. (10). It follows from Eq. (11)

TABLE I. The mass spectrum for the Higgs and gauge bosons in the LR symmetric model in the limit $v_L \simeq 0$ and $\kappa_2 \ll \kappa_1$. ϕ_i^{0r} and ϕ_i^{0i} stand for real and imaginary component of ϕ_i^0. The gauge boson $Z_1(W_1)$ corresponds to the $Z(W)$ boson in the SM.

Particles	Mass^2	Particles	Mass^2
$h^0 = \phi_1^{0r}$	$m_{h^0}^2 = 2\lambda_1\kappa^2$	$H_1^\pm = \phi_1^\pm$	$m^2_{H_1^\pm} = \frac{1}{2}\alpha_3 v_R^2$
$H_1^0 = \phi_2^{0r}$	$m^2_{H_1^0} = \frac{1}{2}\alpha_3 v_R^2$	$H_R^{\pm\pm} = \delta_R^{\pm\pm}$	$m^2_{H_R^{\pm\pm}} = 2\rho_2 v_R^2$
$A_1^0 = -\phi_2^{0i}$	$m^2_{A_1^0} = \frac{1}{2}\alpha_3 v_R^2$	$H_L^\pm = \delta_L^\pm$	$m^2_{H_L^\pm} = \frac{1}{2}(\rho_3 - 2\rho_1) v_R^2$
$H_2^0 = \delta_R^{0r}$	$m^2_{H_2^0} = 2\rho_1 v_R^2$	$H_L^{\pm\pm} = \delta_L^{\pm\pm}$	$m^2_{H_L^{\pm\pm}} = \frac{1}{2}(\rho_3 - 2\rho_1) v_R^2$
$H_3^0 = \delta_L^{0r}$	$m^2_{H_3^0} = \frac{1}{2}(\rho_3 - 2\rho_1) v_R^2$	$A_L^0 = \delta_L^{0i}$	$m^2_{A_L^0} = \frac{1}{2}(\rho_3 - 2\rho_1) v_R^2$
Z_1	$m^2_{Z_1} = m^2_{W_1}\sec^2\theta_W$	$W_1^\pm = W_L^\pm$	$m^2_{W_1} = g^2\kappa^2/4$
Z_2	$m^2_{Z_2} = \frac{g^2 v_R^2 \cos^2\theta_W}{\cos 2\theta_W}$	$W_2^\pm = W_R^\pm$	$m^2_{W_2} = g^2 v_R^2/2$

TABLE II. The cubic and quartic scalar couplings between Higgs singlets and multiplets, where HH^* stands for any states of $\{h^0h^0, H_1^0H_1^0, A_1^0A_1^0, H_1^+H_1^-\}$ and $\Delta\Delta^*$ stands for any states of $\{H_L^0H_L^0, A_L^0A_L^0, H_L^+H_L^-, H_L^{++}H_L^{--}, H_2^0H_2^0, H_R^{++}H_R^{--}\}$.

Interaction	Coupling	Interaction	Coupling	Interaction	Coupling	Interaction	Coupling
$S_DS_DS_\sigma S_\sigma$	$2\lambda_{\sigma D}$	$S_DS_Dh^0$	$\lambda_{1,D}\kappa$	$S_DS_DS_\sigma$	$2(\mu_{\sigma D} + \lambda_{\sigma D}v_\sigma)$	$S_DS_DH_2^0$	$\lambda_{3,D}v_R$
$S_DS_DHH^*$	$\lambda_{1,D}$	$S_\sigma S_\sigma h^0$	$-\mu_{1,\sigma} + \lambda_{1,\sigma}\kappa$	HH^*S_σ	$-\mu_{1,\sigma} + \lambda_{1,\sigma}v_\sigma$	$S_\sigma S_\sigma H_2^0$	$\lambda_{3,\sigma}v_R$
$S_DS_Dh^0H_1^0$	$2\lambda_{2,D}$	$S_DS_DH_1^0$	$2\lambda_{2,D}\kappa$	$h^0H_1^0S_\sigma$	$2(-\mu_{2,\sigma} + \lambda_{2,\sigma}v_\sigma)$	$S_\sigma S_\sigma S_\sigma$	$6\lambda_\sigma v_\sigma$
$S_DS_D\Delta\Delta^*$	$\lambda_{3,D}$	$S_\sigma S_\sigma H_1^0$	$2(-\mu_{2,\sigma} + \lambda_{2,\sigma}\kappa)$	$\Delta\Delta^*S_\sigma$	$-\mu_{3,\sigma} + \lambda_{3,\sigma}v_\sigma$	$h^0h^0H_2^0$	$\alpha_1 v_R$

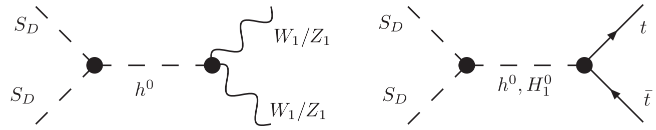

FIG. 1. Feynman diagrams for two DM candidate S_D annihilating into W_1W_1/Z_1Z_1 and $t\bar{t}$ final states.

that the coefficients for the quartic couplings and mixing terms in the potential $\mathcal{V}_S$ are roughly at the same order of magnitude

$$\lambda_\sigma \simeq \lambda_D \simeq \lambda_{\sigma D}, \qquad \lambda_{i,\sigma} \simeq \lambda_{i,D}. \tag{12}$$

One of the implications of this parameter region is that the mixing between S_σ and the SM-like Higgs h^0 will be small. This is because the mixing angle θ is proportional to

$$\tan 2\theta \simeq \frac{\kappa(-\tilde{\mu}_{1,\sigma} + \tilde{\lambda}_{1,\sigma} v_\sigma)}{2\lambda_\sigma v_\sigma^2 - 2\lambda_1 \kappa^2} \sim \mathcal{O}\left(\frac{\kappa}{v_\sigma}\right). \tag{13}$$

Thus the constraints from the precision electroweak data from LEP experiments become weak. There are of course other possible parameter regions. However, one will see in the next section that the parameter space corresponding to the approximate global $U(1)$ symmetry leads to the correct magnitude of the relic dark matter density.

A. Annihilation cross section

The relic density of the gauge-singlet DM S_D can be calculated from the annihilation cross sections which depend largely on mass spectrum of the particles in the model, especially the mass of the DM candidate. S_D can be very light. For 3 GeV $\lesssim m_D \lesssim$ 8 GeV which is consistent with the recent DAMA results [35]. S_D pairs can only annihilate to light fermion pairs through intermediate SM-like Higgs boson h^0. In this case, there is one-to-one correspondence between the DM relic density and the WIMP-nucleon elastic scattering cross section. The ratio between the two only depends on the mass of DM [22,36]. In order to satisfy both the DM relic density $0.105 \leq \Omega_{\rm DM} h^2 \leq 0.117$ [37] and the WIMP-nucleon elastic scattering cross section in the range 3×10^{-41} cm$^2 \lesssim \sigma_n^{\rm SI} \lesssim 5 \times 10^{-39}$ cm^2 reported by DAMA [35], a large $h^0 S_D S_D$ coupling is inevitable, which may cause the invisible decay of h^0 produced at LHC [36].

Here we consider a different parameter rang in which S_D is heavier than the SM-like Higgs and in a mass range $200 \leq m_D \leq 500$ GeV which can be covered by the CDMS and other experiments. Since we assume $v_\sigma \sim v_R \sim 10$ TeV, most of the scalars are heavy except for the SM-like one. In this case, the possible annihilation products are h^0h^0, W_1W_1/Z_1Z_1, and fermion pairs $q\bar{q}$, as shown in Figs. 1 and 2. The $q\bar{q}$ final states are dominated by the heavy t- quarks since the Yukawa coupling is the largest. For W_1W_1/Z_1Z_1 finale states the only possible intermediate state is h^0. For $q\bar{q}$ final states, the intermediate particles can be h^0 and H_1^0. But H_1^0 contribution is negligible as $m_{H_1^0} \gg m_{h^0}$. Since $H_{2,3}^0$ have nonzero $B-L$ charge they can only couple to Majorana neutrinos. For a high v_R around a few TeV, the right-handed neutrinos are also heavy, which cannot appear in the final states. Thus in our model the dominant annihilation processes in Fig. 1 are the same as in the minimal extension of SM with a gauge-singlet [22,26]. For h^0h^0 final states, the s-channel involves h^0, $H_{1,2}^0$ and S_σ while the t-channel involves h^0 only.

The relevant annihilation cross sections for Fig. 1 are given by

$$\begin{aligned}
(4E_1E_2\sigma v)_{W_1W_1} &= \frac{\lambda_{1,D}^2}{8\pi}\left(1 - \frac{m_{h^0}^2}{s}\right)^{-2}\left(1 - 4\frac{m_{W_1}^2}{s} + 12\frac{m_{W_1}^4}{s^2}\right)\left(1 - \frac{4m_{W_1}^2}{s}\right)^{1/2},\\
(4E_1E_2\sigma v)_{Z_1Z_1} &= \frac{\lambda_{1,D}^2}{16\pi}\left(1 - \frac{m_{h^0}^2}{s}\right)^{-2}\left(1 - 4\frac{m_{Z_1}^2}{s} + 12\frac{m_{Z_1}^4}{s^2}\right)\left(1 - \frac{4m_{Z_1}^2}{s}\right)^{1/2},\\
(4E_1E_2\sigma v)_{tt} &= \frac{3\lambda_{1,D}^2}{4\pi}\frac{m_t^2}{s}\left(1 - \frac{4m_t^2}{s}\right)^{3/2}\left(1 - \frac{m_{h^0}^2}{s}\right)^{-2},
\end{aligned} \tag{14}$$

and that for Fig. 2 is

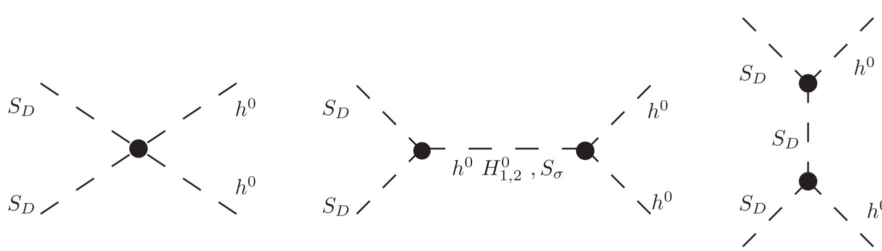

FIG. 2. Feynman diagrams for two DM candidate S_D annihilating into two SM-like Higgs bosons.

$$(4E_1E_2\sigma v)_{h^0h^0} = \frac{\lambda_{1,D}^2}{16\pi}\left(1 - \frac{4m_{h^0}^2}{s}\right)^{1/2} \times \left[\left(\frac{s-4m_{h^0}^2}{s-m_{h^0}^2} - \frac{m_\sigma^2}{s-m_\sigma^2} - \frac{\alpha_1\lambda_{3,D}}{\lambda_{1,D}}\frac{v_R^2}{s-m_{H_2^0}^2}\right)^2 + 4\lambda_{1,D}\frac{\kappa^2}{s-2m_{h^0}^2-2m_D^2} \times \left(\frac{s-4m_{h^0}^2}{s-m_{h^0}^2} - \frac{m_\sigma^2}{s-m_\sigma^2} - \frac{\alpha_1\lambda_{3,D}}{\lambda_{1,D}}\frac{v_R^2}{s-m_{H_2^0}^2}\right)Y(\xi) + 2\lambda_{1,D}^2\left(\frac{\kappa^4}{(m_D^2-m_{h^0}^2)^2} + \frac{4\kappa^4}{(s-2m_{h^0}^2-2m_D^2)^2}Y(\xi)\right)\right] \quad (15)$$

where s is the squared center-of-mass energy. E_1 and E_2 are the energies of the incidental particles. The quantity Y is defined as $Y(\xi) \equiv \mathrm{arctanh}(\xi)/\xi$ with $\xi \equiv \sqrt{(s-4m_{h^0}^2)(s-4m_D^2)}/(s-2m_{h^0}^2-2m_D^2)$. For the cross section $(4E_1E_2\sigma v)_{h^0h^0}$ in Eq. (15), the $H_2^0h^0h^0$ scalar coupling $\alpha_1 v_R$ has been used [5].

B. Constraints from the DM relic density

The thermal-average of the annihilation cross section times the relative velocity $\langle\sigma v\rangle$ is a key quantity in the determination of the cosmic relic abundance of S_D. For nonrelativistic gases, $\langle\sigma v\rangle$ can be expanded in powers of relative velocity and x^{-1} ($x \equiv m_D/T$). To the first order $\langle\sigma v\rangle \simeq \sigma_0 x^{-n}$, where $n = 0(1)$ for $s(p)$-wave annihilation process [38]. The general formula for $\langle\sigma v\rangle$ is given by [39]

$$\langle\sigma v\rangle = \sigma_0 x^{-n} = \frac{1}{m_D^2}\left[\omega - \frac{3}{2}(2\omega-\omega')x^{-1} + \ldots\right]_{s/4m_D^2=1}, \quad (16)$$

where $\omega \equiv E_1E_2\sigma v$, and the prime denotes the derivative with respect to $s/4m_D^2$. ω and its derivative are all to be evaluated at $s/4m_D^2 = 1$. The final DM density $\Omega_{\rm DM}h^2$ is given by [38]

$$\Omega_{\rm DM}h^2 = 1.07\times 10^9 \frac{(n+1)x_f^{n+1}}{g_*^{1/2}M_{\rm Pl}\sigma_0}\ \mathrm{GeV}^{-1} \quad (17)$$

with

$$x_f = \ln[0.038(n+1)(g_D/g_*^{1/2})M_{\rm Pl}m_D\sigma_0] - (n+1/2) \times \ln\{\ln[0.038(n+1)(g_D/g_*^{1/2})M_{\rm Pl}m_D\sigma_0]\}, \quad (18)$$

where $M_{\rm Pl} = 1.22\times 10^{19}$ GeV and $g_D = 1$ is the internal degree of freedom of S_D. g_* is the total number of effectively relativistic degrees of freedom at the time of freeze-out. For particles playing the role of cold DM, the relevant freeze-out temperature is $x_f = m_D/T_f \sim 25$. Since we consider the range 200 GeV $\le m_D \le$ 500 GeV in our analysis, we obtain $g_* = 345/4$.

The total annihilation cross section ω is

$$\omega = (E_1E_2\sigma v)_{h^0h^0} + (E_1E_2\sigma v)_{W_1W_1} + (E_1E_2\sigma v)_{Z_1Z_1} + (E_1E_2\sigma v)_{tt}. \quad (19)$$

From Eq. (15) there are seven unknown parameters enter the expression of total annihilation cross section, namely, m_{h^0}, m_D, $\lambda_{1,D}$, $\alpha_1\lambda_{3,D}$, m_σ^2, $m_{H_2^0}^2$, and v_R. But ω is highly insensitive to m_σ and $m_{H_2^0}^2$ when s_σ and H_2^0 masses are

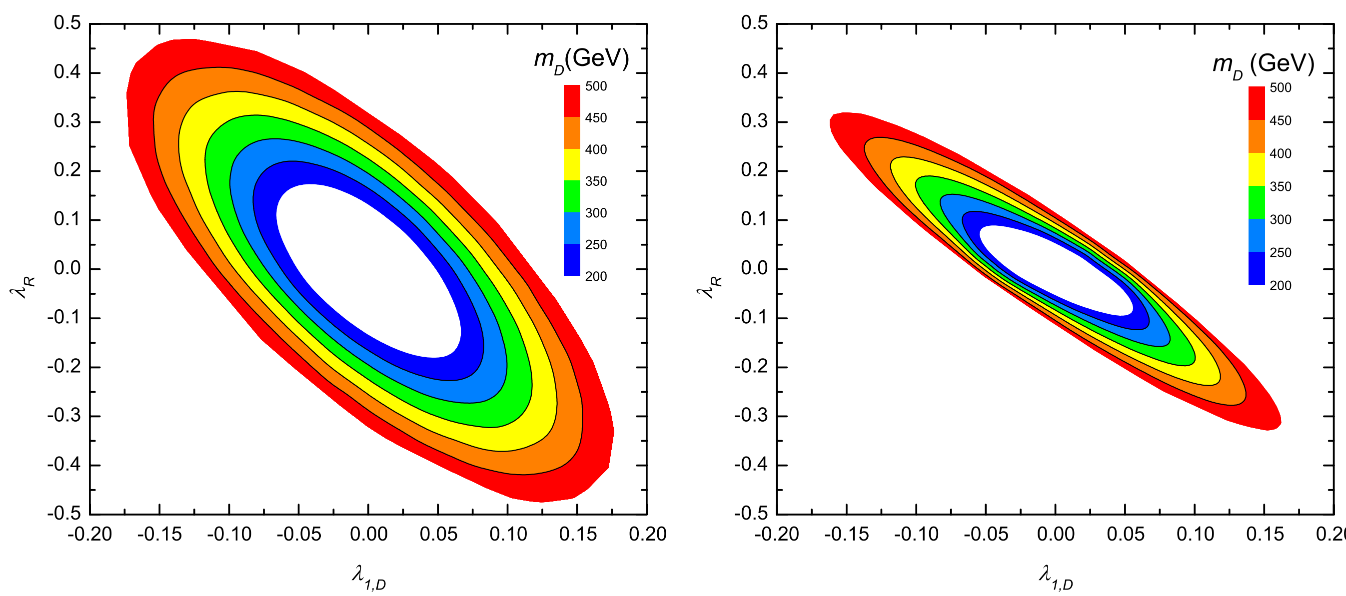

FIG. 3 (color online). The allowed region of $\lambda_{1,D}$ and λ_R for different m_D from DM relic density. The left panel corresponds to the annihilation involving only one Higgs bidoublet. The right panel corresponds to the annihilation involving two Higgs bidoublet. See the text for detailed explanation.

around v_R. Furthermore, the mass of H_2^0 can be related to v_R through $m_{H_2^0}^2 \approx 2\rho_1 v_R^2$ as it is shown in Table I. Thus only four parameters

$$m_{h^0}, m_D, \lambda_{1,D} \quad \text{and} \quad \lambda_R \equiv \alpha_1 \lambda_{3,D}/(2\rho_1)$$

are relevant to our numerical analysis. In numerical calculations, we fix the mass of the SM-like Higgs to $m_{h^0} = 120$ GeV, and perform a numerical scan over the parameters λ_R and $\lambda_{1,D}$ for the mass range 200 GeV $\leq m_D \leq$ 500 GeV. The allowed parameter space is shown in Fig. 3 (left panel), which gives an allowed range

$$-0.18 \lesssim \lambda_{1,D} \lesssim 0.18 \quad \text{and} \quad -0.48 \lesssim \lambda_R \lesssim 0.47.$$

The central region of this figure is excluded since these points cannot provide large enough annihilation cross section to give the desired relic abundance. For such a mass range of S_D, without significant fine-tuning, one needs $\tilde{\lambda} \lesssim 0.01$ and $\tilde{\mu}/v_\sigma \lesssim 0.01$ from Eq. (10), where $\tilde{\lambda}$ and $\tilde{\mu}$ denote the corresponding parameters in Eq. (11). Since the approximate global symmetry $U(1)$ requires $\tilde{\lambda} \ll \lambda_{1,D}$, the region near $\lambda_{1,D} = 0$ in Fig. 3 is disfavored.

C. Predictions for the DM direct detection experiments

The current DM direct detection experiments, such as the CDMS [40] and XENON [41], have imposed strong constraints on the WIMP-nucleus elastic scattering cross section for a wide range of DM mass. In our model, the DM candidate S_D interacts with nucleus $\mathcal{N}$ through Yukawa couplings interactions. For scalar interactions, the spin-independent elastic scattering cross section on a nucleus $\mathcal{N}$ is given by [42,43]

$$\sigma_{\mathcal{N}} = \frac{4M^2(\mathcal{N})}{\pi}[Zf_p + (A-Z)f_n]^2, \tag{20}$$

where $M(\mathcal{N}) = m_D M_{\mathcal{N}}/(m_D + M_{\mathcal{N}})$ and $M_{\mathcal{N}}$ is the target nucleus mass. Z and $A-Z$ are the numbers of protons and neutrons in the nucleus. $f_{p,n}$ is the coupling between WIMP and protons or neutrons, given by

$$f_{p,n} = \sum_{q=u,d,s} f_{Tq}^{(p,n)} a_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{TG}^{(p,n)} \sum_{q=c,b,t} a_q \frac{m_{p,n}}{m_q}, \tag{21}$$

where $f_{Tu}^{(p)} = 0.020 \pm 0.004$, $f_{Td}^{(p)} = 0.026 \pm 0.005$, $f_{Ts}^{(p)} = 0.118 \pm 0.062$, $f_{Tu}^{(n)} = 0.014 \pm 0.003$, $f_{Td}^{(n)} = 0.036 \pm 0.008$, and $f_{Ts}^{(n)} = 0.118 \pm 0.062$ [44]. The coupling $f_{TG}^{(p,n)}$ between WIMP and gluons from heavy quark loops is obtained from

$$f_{TG}^{(p,n)} = 1 - \sum_{q=u,d,s} f_{Tq}^{(p,n)}. \tag{22}$$

Traditionally, the results of WIMP-nucleus elastic scattering experiments are presented in the form of a normalized WIMP-nucleon scattering cross section σ_n^{SI} in the spin-independent case, which is straightforward

$$\sigma_n^{\mathrm{SI}} = \frac{1}{A^2}\frac{M^2(n)}{M^2(\mathcal{N})}\sigma_{\mathcal{N}}, \tag{23}$$

where $M(n) = m_D M_n/(m_D + M_n)$ is the reduced mass of the nucleon, and $M_n = m_{p,n}$ denotes the nucleon mass.

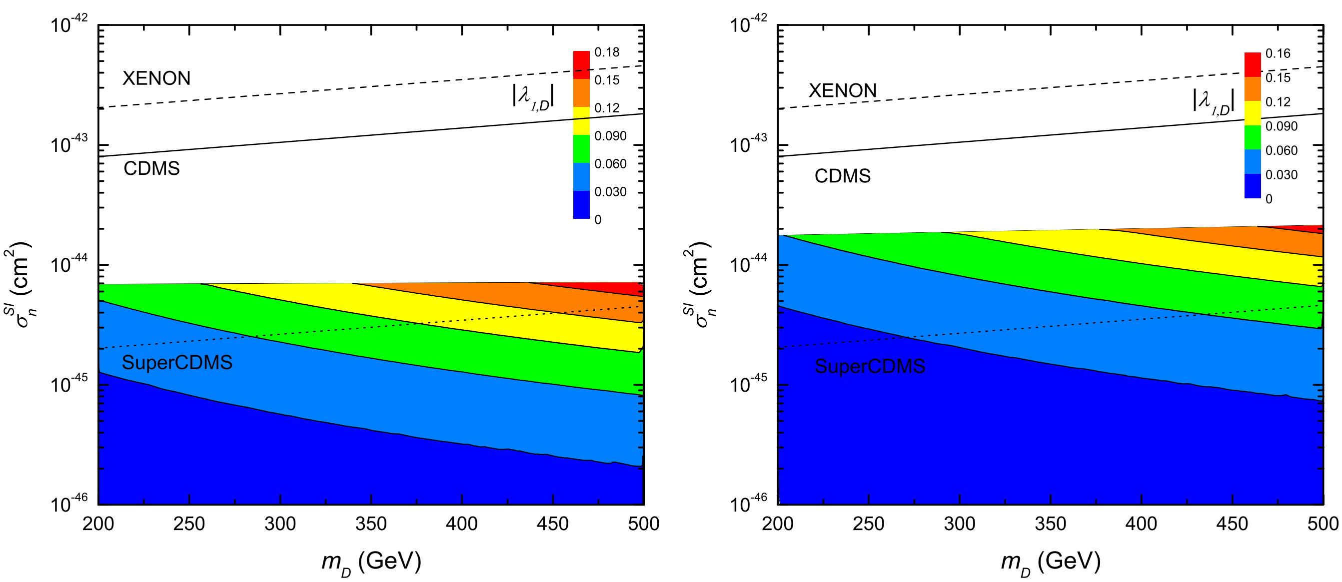

FIG. 4 (color online). Predicted region of the spin-independent WIMP-nucleon elastic scattering cross section σ_n^{SI} as functions of m_D and $\lambda_{1,D}$. The dashed line and solid line denote the present experimental upper bounds from the XENON and CDMS, respectively. The dotted line indicates the sensitivity of the future SuperCDMS [45]. The left panel corresponds to the annihilation involving only one Higss bidoublet. The right panel corresponds to the annihilation involving two Higgs bidoublet with the assumption of $y_q^{\eta_1} \simeq y_q^{\eta_2} \simeq y_q^h$. See the text for detailed explanation.

Because of $f_p \approx f_n$ in our model

$$\sigma_n^{\rm SI} \approx \frac{4M_n^2}{\pi} f_n^2. \quad (24)$$

The present bounds on the WIMP-nucleon elastic scattering cross section are $\sigma_n^{\rm exp} \leq 8 \times 10^{-44}$ cm^2–2×10^{-43} cm^2 from the CDMS [40] and $\sigma_n^{\rm exp} \leq 2 \times 10^{-44}$ cm^2–4.3×10^{-43} cm^2 for from the XENON [41] for the DM mass range 200–500 GeV.

The DM candidate S_D interacts with nucleus $\mathcal{N}$ through their couplings with quarks by exchanging Higgs bosons h^0 and H_1^0. Because H_1^0 is much heavier than h^0, the main contribution comes from the h^0-exchange only. In this case, the WIMP-quark coupling a_q in Eq. (21) is given by

$$a_q = \frac{\lambda_{1,D} y_q^h \kappa}{2\sqrt{2} m_D m_{h^0}^2} \quad (25)$$

where y_q^h $(q = u, d, s, c, b, t)$ denotes the Yukawa coupling of the SM-like Higgs to the quarks with $y_q^h \kappa/\sqrt{2} = m_q$. Using the allowed ranges for $\lambda_{1,D}$ we make predictions for the spin-independent WIMP-nucleon elastic scattering cross section $\sigma_n^{\rm SI}$. The numerical results are shown in Fig. 4 (left panel). One finds $\sigma_n^{\rm SI} \lesssim 7 \times 10^{-45}$ cm^2 for 200 GeV $\leq m_D \leq$ 500 GeV, which is far below the current experimental upper bounds. Nevertheless the future SuperCDMS (Phase A) experiment [45] is able to cover part of the allowed parameter space, especially in the small m_D region.

IV. CONTRIBUTIONS FROM THE OTHER HIGGS BIDOUBLET

In this section, we generalize the previous discussions to the case in which the other bidoublet χ mixes significantly with ϕ and $\Delta_{L,R}$. In this case the SCPV can be easily realized [16–18]. Comparing with the previous case, the main differences are that there could be more scalar particles entering the DM annihilation and scattering processes. Furthermore, the new contributions from these particles may modify the correlation between the DM annihilation and WIMP-nucleus elastic scattering cross sections, which leads to significantly different predictions from the other gauge-singlet scalar DM models and the previous discussions.

As shown in Eq. (1), the second Higgs bidoublet χ contains two neutral Higgs particles $\chi_{1,2}^0$. After the SSB, $\chi_{1,2}^0$ may obtain VEVs $w_{1,2}/\sqrt{2}$. The squared sum of all the VEVs including $\kappa_{1,2}$ should still lead to $v = (|\kappa_1|^2 + |\kappa_2|^2 + |w_1|^2 + |w_2|^2)^{1/2} \approx 246$ GeV. In the physical basis, some of the Higgs bosons from χ could be light around electroweak scale. The number of the light Higgs depends on the Higgs potential. In most cases, there are two more light neutral Higgs $\eta_{1,2}^0$ and one pair of light charged Higgs $\eta^\pm$ [18]. This feature can be easily understood in the limit $\kappa_2 \sim w_2 \sim 0$. In this case, one can determine h^0, η_1^0, and η_2^0 from ϕ_1^0 and χ_1^0, and $\eta_\pm$ from the mixing of $\phi_1^\pm$ and $\chi_1^\pm$. The number of kinematically allowed DM annihilation processes depends on the masses of the relevant particles. Here we consider a case in which S_D is heavier than all the light scalars and the SM-like h^0 remains the lightest scalar, i.e. $m_{h^0} \leq m_{\eta_{1,2}^0}$, $m_{\eta^+} \leq m_D$. The quartic interaction and the s-channel annihilation in Fig. 2 now have seven possible final states. They are combinations of any two of the three neutral states (h^0, η_1^0, η_2^0) and charged final states $\eta^+\eta^-$. Note that each s-channel diagram in Fig. 2 may have h^0, η_1^0, and η_2^0 as intermediate states besides S_σ and H_2^0. The t-channel diagram has six possible final states, due to the absence of the cubic scalar vertexes $S_D S_D \eta^\pm$.

The cubic coupling $S_D S_D \eta_{1,2}^0$ although can be different from that for $S_D S_D h^0$, may not modify the correlation between the DM annihilation and WIMP-nucleon elastic scattering cross sections in a significant way. As it is pointed out in Ref. [36], the ratio $R \equiv \langle \sigma v \rangle / \sigma_n^{\rm SI}$ is highly insensitive to these couplings because they cancel out largely. R is only sensitive to the mass of DM candidate and the Yukawa couplings. In the minimal scalar DM model the Yukawa couplings are the same as that in the SM. It is shown that the value R scales as m_D^2 [27]. For small m_D around a few GeV, the value of R is in agreement with the DAMA results. A large m_D around a few hundred GeV corresponds to a large R, which indicates that the WIMP-nucleon elastic scattering cross section may be far below the current direct detection bounds.

In 2HBLR the Yukawa couplings can be significantly different from those in the SM and the minimal LR model. Similar to the general two-Higgs-doublet model [46–53] the Yukawa couplings are not simply determined by the quark masses. This is because with the introduction of the additional bidoublet, the fermion mass matrices and Yukawa matrices are not proportional to each other. In general the Yukawa couplings can be parametrized as $y_q = \sqrt{2}\xi_q m_q/v$, the factor ξ_q depends on fermion flavor q and can be different from unity. For the DM annihilation processes, the heavy quark contribution dominates, while for WIMP-nucleon scattering processes the light quarks are more important as the quark mass dependence are reduced in Eq. (21). For a large ξ_q for light quark sector it is possible that the prediction for WIMP-nucleon elastic scattering cross section can be enhanced and the cross section for DM annihilation still coincides with the observed DM relic density.

For a concrete numerical illustration, we choose all the masses $m_{\eta_1^0}$, $m_{\eta_2^0}$, $m_{\eta^\pm}$ at 180 GeV and keep $m_{h^0} =$ 120 GeV. For cubic and quartic scalar couplings, we assume they are the same as that for the SM-like Higgs. Namely, the couplings of $S_D S_D \eta_{1,2}^0$ and $S_D S_D \eta_{1,2}^0 \eta_{1,2}^0$ are set equal to $\lambda_{1,D} v$ and $\lambda_{1,D}$, respectively. Similarly, the cubic scalar couplings among the light Higgs particles h^0,

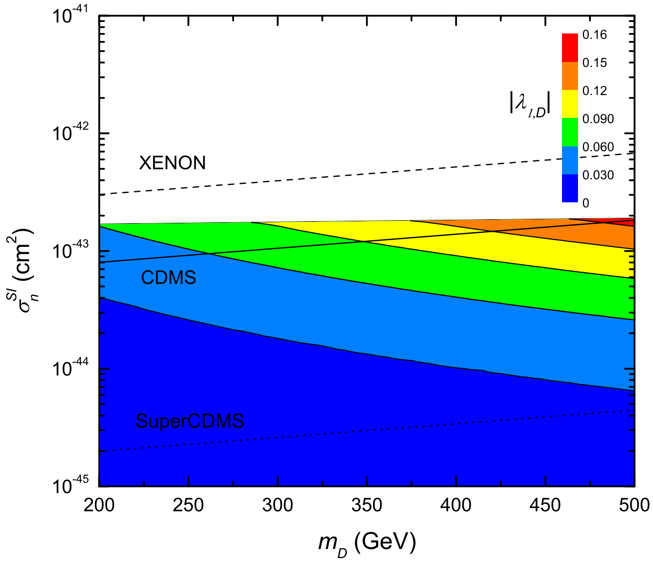

FIG. 5 (color online). Predicted region of the spin-independent WIMP-nucleon elastic scattering cross section $\sigma_n^{\rm SI}$ as functions of m_D and $\lambda_{1,D}$ in two Higgs bidoublet model with the assumption of $y_q^{\eta_1} \simeq y_q^{\eta_2} \simeq 10 y_q^h$ for three light quarks.

$\eta_{1,2}^0$, and $\eta^\pm$ are set equal to $3m_{h^0}^2/v$. For a comparison we consider two sets of Yukawa couplings:

(i) All the couplings y_q^h, $y_q^{\eta_1}$, $y_q^{\eta_2}$ for $h^0 q\bar{q}$, $\eta_1^0 q\bar{q}$, $\eta_2^0 q\bar{q}$, respectively are nearly the same:

$$y_q^{\eta_1} \simeq y_q^{\eta_2} \simeq y_q^h. \tag{26}$$

with $y_q^h = \sqrt{2} m_q/v$. In this case the annihilation cross section can be obtained simply by counting the number of new channels. In Fig. 3 (right panel), we give the constraints on $\lambda_{1,D}$, λ_R for different m_D. It is clear that there is a stronger constraint on the allowed parameter space, due to the increased number of intermediate and final states. For the WIMP-nucleon elastic scattering process, the WIMP-quark coupling a_q in Eq. (25) is given by

$$a_q = \frac{\lambda_{1,D} v}{2\sqrt{2} m_D}\left(\frac{y_q^h}{m_{h^0}^2} + \frac{y_q^{\eta_1}}{m_{\eta_1^0}^2} + \frac{y_q^{\eta_2}}{m_{\eta_2^0}^2}\right). \tag{27}$$

Using the allowed $\lambda_{1,D}$ and m_D from Fig. 3, we calculate the spin-independent WIMP-nucleon elastic scattering cross section $\sigma_n^{\rm SI}$. The numerical results are shown in Fig. 4 (right panel). We find that the $\sigma_n^{\rm SI}$ is enlarged roughly by a factor of 3, which is however still below the current experiment upper bounds.

(ii) The Yukawa couplings for $y_q^{\eta_1}$ and $y_q^{\eta_2}$ are significantly larger in the light quark sector (for $q = u, d, s$)

$$y_q^{\eta_1} \simeq y_q^{\eta_2} \simeq 10 y_q^h. \tag{28}$$

Since the annihilation process $S_D S_D \to q\bar{q}$ is dominated by heavy t-quarks, the enhanced Yukawa couplings $y_q^{\eta_1}$ and $y_q^{\eta_2}$ do not affect the total annihilation cross section. Thus the DM relic density remains unchanged. However, the predicted WIMP-nucleon scattering cross section $\sigma_n^{\rm SI}$ will be enhanced. The corresponding results have been shown in Fig. 5. We find that in this case $\sigma_n^{\rm SI}$ is enhanced by an order of magnitude compared with the one-Higgs bidoublet case. The future DM direct detection experiment SuperCDMS can cover most of the allowed parameter space.

V. CONCLUSIONS

In summary, we have discussed the possibility that the stability of DM can be protected by the fundamental symmetries P and CP of quantum field theory. It can be realized in the framework of a generalized LR symmetric model which allows SCPV. The DM candidate in our model is a gauge-singlet which transforms under CP as an ordinary complex scalar. In this model no extra discrete symmetry is required. We have scanned the parameter space allowed by the relic DM density and made predictions for direct detection experiments. We found that the model was in agreement with the current measurement in a large parameter space. Based on the constrained parameter space, we have made predictions for the WIMP-nucleon spin-independent cross sections, and further studied the correlations with the DM annihilation. We have found that in this model the correlation could be significantly different from other gauge-singlet DM models. The DM-nucleon elastic scattering cross section could reach the current experimental upper bound for large Yukawa couplings for light quarks, which could be tested by future experiments.

ACKNOWLEDGMENTS

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