In view of ongoing measurements of the Higgs-like boson at the LHC and direct searches for dark matter, we explore the possibility of accommodating the potential results in a simple new-physics model with discrete gauge symmetry as well as light neutrino masses. Specifically, we study collider and relic-density constraints on the new gauge coupling, predict the cross section of the dark matter scattering off nucleons, and compare it with current direct search data. We also discuss some of the implications if the dark matter is light. The new gauge sector of the model allows it to be compatible with the latest LHC information on the Higgs-like particle and simultaneously satisfy the requirements in its dark matter sector.

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I. INTRODUCTION

The particle with mass around 125 GeV recently discovered at the LHC has characteristics which, based on the measurements reported so far [1], suggest that it may be the Higgs boson of the standard model (SM). Whether that is indeed the case will likely become clearer in the near future after sufficiently more data are accumulated at the LHC. Needless to say, the imminent potential confirmation on the Higgs’ existence will have far-reaching implications for attempts to identify the nature of new physics beyond the SM. For it is widely accepted that new physics is necessary at least to explain the astronomical evidence for dark matter (DM) and the numerous experimental indications of neutrino mass [2]. In anticipating the outcomes of the Higgs quest as well as upcoming results of direct searches for DM, which are also ongoing, it is of great interest to explore a simple framework of new physics that can accommodate the possibilities.

To do so, we adopt an economical model that offers not only a Higgs boson and DM of the popular weakly interacting massive particle (WIMP) type, but also a means to produce light neutrino masses. The parameter space of the model has ample room for a Higgs with SM characteristics or one that is less SM-like. The Higgs and DM sectors are linked in that the Higgs may decay substantially into DM particles if kinematically allowed. This makes the Higgs and DM searches complementary for probing the model. Another of its salient features is that the DM stability is realized under a discrete gauge symmetry. We collect the quantum number assignments for the fermion and scalar fields of our model in Table I. One can therefore see that U(1)$_\zeta$ is none other than the usual complex doublet. For quarks and leptons, $\zeta_f = 1/3$ and $-1$, respectively.

II. MODEL DESCRIPTION

We introduce a minimal number of new particles for our purposes: three right-handed neutrinos, $\nu_{\ell R}$, two complex scalar fields, $S$ and $D$, which are singlets under the SM gauge group, and a $Z'$ boson for the U(1)$_\zeta$ gauge symmetry. We collect the quantum number assignments for the fermion and scalar fields of our model in Table II. We assume that the SM fermions carry U(1)$_\zeta$ charges, and this allows us to introduce right-handed neutrinos for gauge-anomaly cancellation and generating light neutrino masses by means of the well-known seesaw mechanism [3], which is activated with the involvement of the same new scalar field.

This paper is organized as follows. In Sec. II, we describe our model in greater detail. We also briefly compare the important features of our model with a number of other scenarios in the literature. In Sec. III, we obtain constraints on the new gauge sector from collider experiments. In Sec. IV, we address further constraints on the model from the relic abundance data and predict the DM-nucleon scattering cross-section subject to direct detection measurements. We also discuss the invisible decay of the Higgs boson if the DM is sufficiently light and consider some of the implications in relation to the Higgs hunt at the LHC. Finally, Sec. V contains the summary of our work and some additional discussions.

| $f_{SM}$ | $\nu_{\ell R}$ | $H$ | $S$ | $D$ |
|---------|---------------|-----|-----|-----|
| SU(2), U(1)$_Y$ | $g_{SM}$ | 1, 0 | 2, 1/2 | 1, 0 | 1, 0 |
| U(1)$_\zeta$ | $\zeta_f$ | [-1] | 0 [+1] | 2 [+2] | 1 [-] |

TABLE I: Charge assignments of the fermions and scalars in the model. $f_{SM}$ ($g_{SM}$) denotes SM fermions (their assignments) and $H$ the usual complex doublet. For quarks and leptons, $\zeta_f = 1/3$ and $-1$, respectively.
than $U(1)_{B-L}$, with $B$ and $L$ referring to baryon and lepton numbers, respectively. The kinetic terms of $S$ and $D$ along with a renormalizable potential $V$ for them and $H$ are given by

$$\mathcal{L} = \langle D^\mu D^\nu \rangle \partial^\mu \partial_\nu D - \langle D^\mu D^\nu \rangle S \partial^\mu \partial_\nu S - V, \quad (1)$$

$$V = \mu_D^2 |D|^2 - \mu_S^2 |S|^2 + \mu_{DS} (D S^\dagger + H.c) + 2\lambda_{DS} |D|^2 |S|^2 + 2(\lambda_{DH} |D|^2 + \lambda_{HS} |S|^2) H \dagger H \tag{2}$$

where $\mathcal{D}_\mu = \partial_\mu + ig_\zeta Z_\mu$ with the coupling $g_\zeta$ and charge $\zeta$ associated with $U(1)_c$. The $H$ field develops a nonzero VEV as in the SM. The non-SM parts of $V$ were discussed before (e.g., Refs. [8]). The role of DM is played by the $D$ field—sometimes dubbed the darkon—via its lighter component. To maintain the $Z_2$ symmetry and hence the DM longevity, $D$ must have zero VEV. The presence of both the $S$ field and the $\mu_{DS}$ term is essential because the latter triggers the spontaneous breakdown $U(1)_c \to Z_2$ when $S$ gets a nonzero VEV.

The parameters in $V$ should be chosen such that the vacuum has the above desired properties. Accordingly, we assume that all the $\lambda$'s in $V$ are positive to render it bounded from below. Subsequently, upon expressing the VEV's of $H$ and $S$ as $\langle H \rangle = (0, v_H^2)/\sqrt{2}$ and $\langle S \rangle = v_S/\sqrt{2}$, with $v_H > 0$, we arrive at

$$v_{H(S)}^2 = \frac{\lambda_S(H)}{\lambda_H S} - \frac{\lambda_{HS} \mu_{S(H)}^2}{\lambda_H S - \lambda_{HS}^2} \tag{3}$$

and so

$$\mu_{H(S)}^2 = \lambda_H(S) v_{H(S)}^2 + \lambda_{HS} v_{S(H)}^2 > 0. \tag{4}$$

Furthermore, writing $D = (D_R + i D_I)/\sqrt{2}$ in terms of its real and imaginary components leads to the combinations

$$m_{D_R, D_I}^2 = \mu_D^2 + \lambda_{DH} v_H^2 + \lambda_{DS} v_S^2 \pm \sqrt{2} \mu_{DS} v_S > 0, \tag{5}$$

which are the squared masses of $D_R$ and $D_I$. The mass difference between $D_R$ and $D_I$ therefore depends on the sign of $\mu_{DS}$. We will take $\mu_{DS} > 0$ so that $D_I$ acts as the WIMP DM. The results of our analysis would be the same if we took $\mu_{DS} < 0$, only that the roles of $D_R$ and $D_I$ would be interchanged.

After electroweak symmetry breaking, the remaining field $h'$ in $H = (0, v_H + h')^2/\sqrt{2}$ will mix with $s' = \sqrt{2} S - v_S$ because of the $\lambda_{HS}$ term in $\mathcal{L}$. This results in the mass eigenstates

$$h = h' \cos \theta + s' \sin \theta, \quad s = s' \cos \theta - h' \sin \theta, \tag{6}$$

with the mixing angle $\theta$ and masses $m_{h,s}$ given by

$$\tan(2\theta) = \frac{M_{H_s}^2}{M_H^2 - M_S^2}, \tag{7}$$

$$2m_{h,s}^2 = M_H^2 + M_S^2 \pm \left[(M_H^2 - M_S^2)^2 + M_{H_s}^4\right]^{1/2}. \tag{8}$$

where

$$M_{H_s}^2 = 2\lambda_{H,s} v_H v_S, \quad M_{H_s}^2 = 4\lambda_{HS} v_H v_S. \quad (9)$$

The lighter state, $h$, is the physical Higgs boson.

In the neutrino sector, the mass-generating terms have the form

$$i\lambda_{kl} \bar{\nu}_{kR} H^T \tau_2 L_{iL} - \frac{1}{2} \lambda_{kl} \bar{\nu}_{kR} (\nu_{iR})^2 S^\dagger + H.c., \tag{10}$$

where $k, l = 1, 2, 3$ are summed over, $\tau_2$ is the second Pauli matrix, and $L_{iL}$ represents a lepton doublet. They give rise to the Dirac and Majorana mass (3×3) matrices

$$M_D = \sqrt{\lambda} v_H, \quad M_{\nu_R} = \sqrt{\lambda} v_S, \quad (11)$$

respectively. Incidentally, the $\lambda'_{kl}$ term in Eq. (10) can play the same role as the $\mu_{DS}$ term in Eq. (2) for symmetry breaking. For the type-I seesaw mechanism to yield the light neutrino masses, the eigenvalues of $M_{\nu_R}$ are expected to be orders of magnitude bigger than a TeV.

Thus, assuming that $\lambda_H$ and $\lambda_S$ in Eq. (5) are roughly of similar order to the eigenvalues of $\lambda'$ in Eq. (11), we have $M_S \gg 1$ TeV and also $M_S \gg M_H$. In addition, we will pick $M_{H_s}^2 \ll M_S^2 - M_H^2$, so that $m_{h,s} \simeq M_{H,s}$ and $\theta \ll 1$, leading to $h \sim h'$ and $s \sim s'$. It follows that we can neglect the effects of the heavy $s$ on the processes of interest and the relevant interactions for $D_{R,I}$ are described by

$$\mathcal{L}_D = -\frac{1}{2} \lambda_{DH} (D_R^2 + D_I^2) \left(v_H h + \frac{1}{2} h^2\right) - \frac{1}{2} \lambda_D (\bar{D}_R^2 + \bar{D}_I^2)^2 + g_\xi \lambda_D \bar{D}_R \bar{D}_I D_I \tag{12}$$

From now on, we consider the possibility that $D_R$ and $D_I$ are nearly degenerate. As a consequence, the DM relic abundance is determined not only by the annihilation rate of $D_I$, but also by that of $D_R$ and/or both of them. In this so-called coannihilation case [8], we assume specifically that

$$\Delta = (m_{D_R} - m_{D_I})/m_{D_I} \simeq 0. \tag{13}$$

Therefore, the relevant reactions are mainly $D_{I(R)} D_{I(R)} \to$ SM particles, from Higgs-exchange and contact diagrams, the latter if $m_{D_I, D_R} > m_H$, and $D_I D_R \to Z'\nu$ to SM fermions, as the fermions carry $U(1)_c$ charges. Moreover, the $\lambda_D$ part in $\mathcal{L}_D$ is not pertinent to our purposes.

Before proceeding to our numerical calculations, we would like to make a few remarks comparing this work to those in the literature. A number of earlier analyses addressed the phenomenology of some of the elements in our model separately, such as scalar DM in the absence of a new gauge sector [8, 10, 12, 13] or the $Z'$ boson of a gauged $U(1)_{B-L}$ symmetry in conjunction with right-handed neutrinos, but no DM candidates [8]. There are also studies dealing with models which possess DM stabilized by the remnant $Z_2$ symmetry of a $U(1)_{B-L}$ that
is global [14] or local [6, 13]. The $Z_2$ symmetry could instead be simply put in by hand [8, 10–13, 16], or it could be accidental [14]. Here we would like to emphasize that the model which we have adopted incorporates the minimal mechanism for both stabilizing scalar DM with the $Z_2$ remnant of a gauged $U(1)_{B-L}$ and generating neutrino mass through the spontaneous breaking of the same group. Furthermore, the most nontrivial aspects of our analysis—in distinction to others’—are that the $Z'$ boson of the $U(1)_{B-L}$ as well as the Higgs boson contribute to the DM interactions with SM particles, and hence to DM annihilation, and that we treat the case where the $Z'$ boson plays a dominant role in determining the DM relic density. As we will demonstrate later, this allows the model to provide not only a good candidate for DM that is consistent with the latest direct-search results, but also a Higgs boson that is SM-like in its couplings to the standard particles and, for sufficiently light DM, has a significant invisible decay mode compatible with the current LHC data.

III. CONSTRAINTS ON THE NEW GAUGE COUPLING

Since $\lambda_{DH}$ and $g_\zeta$ in Eq. (12) are free parameters, we will consider different interesting combinations of their contributions to the relic density and their potential implications for Higgs and DM direct searches. Before doing so, we first constrain $g_\zeta$ using other observables measured at colliders. Since only $S$ induces the $Z'$ mass, $m_{Z'} = 2g_\zeta v_S$, there is no tree-level $Z-Z'$ mixing, implying that at tree level $g_\zeta$ has no effects on the $Z$-pole observables [18], but can affect $e^+e^- \to \ell^+\ell^-$ and hadron collisions into fermion pairs. Neglecting $Z-Z'$ kinetic mixing which can arise at loop level, we focus on $e^+e^- \to \ell^+\ell^-$ and the Drell-Yan (DY) process $pp \to \ell^+\ell^-X$, which can restrict $g_\zeta$ well.

Measurements of $e^+e^- \to \ell^+\ell^-$ for $\ell = \mu, \tau$ were performed at LEP II with center-of-mass energies from 130 to 207 GeV [19]. We employ the data on the cross section and forward-backward asymmetry. Adopting their 90% confidence-level (CL) ranges and employing the formulas given in Ref. [18], but with $s$-dependent $Z$ and $Z'$ widths [19], we extract the upper limit on $g_\zeta$ as a function of $m_{Z'}$, represented by the blue solid curve in Fig. 1.

The latest cross-section data on the DY process from the LHC [20] reveal no deviation from the SM expectations and hence no evidence of a $Z'$ boson. We can derive an upper bound on the coupling constant $g_\zeta$ using the SM cross-section, following the method of Ref. [21]. The DY cross-section is numerically estimated using the CalcHEP package [22] by incorporating new Feynman rules in the model file. We count the events in the invariant-mass window of $\pm 20\%$ around the $Z'$ mass for a luminosity of $1.1 \text{ fb}^{-1}$ according to the recent experimental analysis [20]. This number of signal events is plugged into the one-bin log likelihood $LL = 2[N \ln(N/\nu) + \nu - N]$, where $N (\nu)$ is the number of events predicted by the SM (SM plus the $Z'$ boson). A value of $LL = 2.7$ is taken for the 90% CL, and $\nu$ is solved for. The upper limit on the cross section is then derived from the solved value of $\nu$ for each $Z'$ mass. This upper limit in turn constrains $g_\zeta$, as depicted by the purple dashed curve in Fig. 1.

IV. DARK MATTER PHENOMENOLOGY

In the simplest darkon model, with a real darkon and no new gauge sector, the $m_D < m_h/2$ region has mostly been disfavored by DM direct detection results as well as the recent observation of the Higgs-like resonance at the LHC, for the darkon-Higgs coupling is required to be sizable by the relic density data, rendering the Higgs mostly invisible [23]. In our model, the presence of the $Z'$ boson alters the darkon phenomenology in important ways. Particularly, there are now $Z'$-mediated diagrams, besides the Higgs-mediated ones, contributing to both darkon annihilation and darkon-nucleon interactions. One of the consequences is that the allowed parameter space of the model can still comfortably make room for a light darkon, as we will see below.

A. Relic abundance

In evaluating the darkon contributions to the relic density $\Omega_D$, we employ the relations [6, 24]

$$\Omega_D h_0^2 = \frac{1.07 \times 10^9}{\sqrt{\sigma_{\text{eff}} v_{ \text{rel} }^2}} \text{GeV},$$

$$J = \int_{x_f}^{\infty} dx \frac{\langle \sigma_{\text{eff}} v_{ \text{rel} }^2 \rangle}{x^2},$$

$$x_f = \ln \left[ 0.038 g_{\text{eff}} m_D m_{pl} \langle \sigma_{\text{eff}} v_{ \text{rel} }^2 \rangle \left( g_* x_f \right)^{-1/2} \right],$$

where $h_0$ denotes the Hubble constant in units of $100 \text{ km/s/Mpc}$, $g_*$ is the number of relativistic degrees
of freedom below the freeze-out temperature $T_f$, $m_{\eta_1} = 1.22 \times 10^{19}$ GeV is the Planck mass, $x_f = m_D/T_f$ with $m_D = m_{D_i}$ being the WIMP mass, $\langle \sigma v_{\text{rel}} \rangle \equiv \langle \sigma v \rangle$ is the thermally averaged product of the effective darkon-annihilation cross-section and the relative speed of the darkon pair in their center-of-mass frame, and $g_{\text{eff}}$ is the darkon’s effective number of degrees of freedom in the coannihilation case. Our choice $\Delta \approx 0$ above leads to some simplification [9]. Thus, we have $g_{\text{eff}} \simeq 2$ and

$$\sigma_{\text{eff}} \simeq \frac{1}{4} (\sigma_{II} + 2 \sigma_{IR} + \sigma_{RR}) ,$$

(15)

where $\sigma_{ij}$ denotes the cross section of $D_i D_j$ annihilation into possible SM final states. Since $D_{I,R}$ have the same interactions [Eq. (12)], we have $\sigma_{RR} \simeq \sigma_{II}$ for $\Delta \simeq 0$.

The expression for $\sigma_{II}$ due to Higgs-exchange diagrams follows from its real darkon counterpart [10, 12], and so

$$\langle \sigma_{II} v_{\text{rel}} \rangle = \frac{4\lambda_{D_H}^2 v_H^2 m_D^{-1} \sum_i \Gamma(h \rightarrow X_i)}{(4m_D^2 - m_H^2)^2 + m_h^2} ,$$

(16)

where $h$ is a virtual Higgs with the same couplings as the physical $h$, but with the invariant mass $\sqrt{s} = 2m_D$, and $h \rightarrow X_i$ is any kinematically allowed decay mode of $h$. For $m_{D_i} > m_h$, the $D_i D_i \rightarrow hh$ channel needs to be included [12]. Thus the contributions of $\sigma_{II,RR}$ to $\langle \sigma v \rangle$ and $J$ are

$$\langle \sigma v \rangle_h \simeq \frac{\langle \sigma_{II} v_{\text{rel}} \rangle}{2} , \quad J_h \simeq \frac{\langle \sigma_{II} v_{\text{rel}} \rangle}{2x_f} .$$

(17)

For the $Z'$-mediated counterpart, assuming nonrelativistic darkons and $\Delta \simeq 0$, we derive

$$\sigma_{IR} \simeq \frac{g^{4}_{\text{v,rel}} \sum f N_{Z'}^2 (2m_D^2 + m^2)}{12\pi} \frac{m_D}{m_D^2 - m_{Z'}^2} (4m_D^2 - m_{Z'}^2)^2 + \Gamma_{Z',m^2}^2 ,$$

(18)

where the sum is over fermions with mass $m_f < m_D$ and $N_{Z'}$ colors and $\Gamma_{Z'}$ is the $Z'$ width. It follows that

$$\langle \sigma v \rangle_{Z'} \simeq \frac{3b_{IR}}{x} , \quad J_{Z'} \simeq \frac{3b_{IR}}{2x_f} ,$$

(19)

where $b_{IR} = \sigma_{IR}/v_{\text{rel}}$. Applying

$$\langle \sigma v \rangle = \langle \sigma v \rangle_h + \langle \sigma v \rangle_{Z'} , \quad J = J_h + J_{Z'} ,$$

(20)

in Eq. (14), we can then extract constraints on $\lambda_{D_H}$ and $g_{\xi}$ from the 90%-CL range $0.092 \leq \Omega_D h^2 \leq 0.118$ of the observed relic density [23].

In Fig. 2 we display $g_{\xi}$ as a function of $m_{D_1}$ subject to the relic data, assuming the absence of the Higgs effect and, for definiteness, $m_{Z'} = 300$ GeV, and compare it with the collider bound. This illustrates that only the resonance region, $m_{Z'} \approx 2m_D$, can satisfy both sets of constraints. Since the situation is unchanged in the presence of the Higgs contribution, hereafter we limit the $Z'$ contribution to this resonance case. If the former is nonnegligible, we can consider various combinations of $J_h$ and $J_{Z'}$ with different implications for the prediction in relation to DM direct searches.

**B. Direct detection**

The direct detection of DM is through the recoil of nuclei after it hits a nucleon $N$. In our model, with nearly degenerate $D_{I,R}$, the involved interactions are $D_{I(R)}N \rightarrow D_{I(R)}N$ and $D_{I(R)}N \rightarrow D_{I(R)}N$ via $h$ and $Z'$ exchanges, respectively, in the $t$ channel. The resulting spin-independent cross-section $\sigma_{DN}$ needs to accommodate these possibilities, but also take into account the fact that the DM local density is independent of the number of DM components. Accordingly, we find

$$\sigma_{DN} \simeq \frac{\lambda_{D_H}^2 g_{NNh}^2 m_N^2 v_H^2}{\pi m_D^2 m_h^2} + \frac{g^4_{1/2} m_{DN}^2}{\pi m_{Z'}^2} ,$$

(21)

in the nonrelativistic limit, where $g_{NNh}$ is the Higgs-nucleon effective coupling, $m_N = m_{D_I}/(m_D + m_N)$, the first term is equal in form to the cross section in the real darkon case [10, 12], and for the second term we have used $\langle N | \bar{u} \gamma^\mu u + \bar{d} \gamma^\mu e d | N \rangle = 3N \gamma^\mu N$ with vanishing contributions from the other quarks [24].

Here we look at a couple of representative examples. In the first one, the $Z'$ contribution considerably dominates $J$ in Eq. (20), and we choose $J_{Z'} = 999 J_h$ for definiteness. We illustrate the prediction for $\sigma_{DN}$ as a function of the darkon mass with $m_h = 125$ GeV and $m_{Z'} = 2m_D$ in Fig. 2(a). For this plot, $\lambda_{D_H}$ follows from $J = 10^3 J_h$ plus the relic constraint, $g_{\xi}$ varies in the range allowed by the collider data, $v_H = 246$ GeV, and we have employed the range $0.0011 \leq g_{NNh} \leq 0.0032$ for the Higgs-nucleon coupling [13, 27]. The lightly shaded (lighter orange) portions of the prediction curve indicate that in the $m_D \gtrsim 100$ GeV region the $Z'$ effect also dominates $\sigma_{DN}$ and is enhanced by a few orders of magnitude relative to the Higgs contribution. Evidently, compared to the most recent data from the leading direct searches for WIMP DM, the prediction can largely escape the strictest bounds to date. However, future direct searches...
such as XENON1T would probe this parameter space of the model more stringently.

In Fig. 3(b), we show an example where the Higgs contribution is less suppressed, \( J_h = J_{Z'} / 9 \). In this case, most of the \( m_D < 100 \text{ GeV} \) range is ruled out by the null results of some of the direct searches. Nevertheless, the prediction (red areas) is partly consistent with the possible WIMP hints reported by CoGeNT (purple areas) and CRESST-II (blue areas), although they conflict with the other experiments.

If the darkon annihilation is dominated by the Higgs contribution instead, \( J_h \gg J_{Z'} \), the allowed parameter space for \( m_D < 100 \text{ GeV} \) would be further reduced compared to that in the second example. With \( J \simeq J_h \), the relic density would comprise approximately equal parts from \( D_{I,R} \) and consequently \( \lambda_{DH} \) (or \( \sigma_{DN} \)) would be about \( \sqrt{2} \) (2) times the corresponding coupling (cross section) in the simplest darkon model. We note that in all these instances darkon masses larger than \( \sim 100 \text{ GeV} \) are still viable and will be probed by future measurements.

C. Invisible decay of the Higgs boson

For \( m_h > 2m_{D,R,D_I} \), the decays \( h \rightarrow D_{I,R}D_{I,R} \) will occur and, for non-negligible \( \lambda_{DH} \), can substantially enhance the invisible decay rate of the Higgs. Their combined branching ratio is

\[
B(h \rightarrow D_ID_I + D_RD_R) = \frac{\Gamma_{h \rightarrow D_ID_I} + \Gamma_{h \rightarrow D_RD_R}}{\Gamma_h} \tag{22}
\]

where the Higgs width \( \Gamma_h = \Gamma_{h \rightarrow SM} + \Gamma_{h \rightarrow D_ID_I} + \Gamma_{h \rightarrow D_RD_R} \) includes the SM Higgs width \( \Gamma_{h \rightarrow SM} \). We depict in Fig. 4 the branching ratios for the same \( J_h,J_{Z'} \) and \( m_h \) choices made in Fig. 3 and \( \Delta \simeq 0 \). Clearly, unless \( J_h \) is very small relative to \( J_{Z'} \), these invisible channels tend to dominate the Higgs width. This is similar to the real darkon case \([10, 12, 13]\), in which the Higgs can be hidden from sight for sufficiently low \( m_h > 2m_D \). The model can thus readily account for the possibility that the LHC does not detect any Higgs. On the other hand, if an SM-like Higgs is observed at \( \sim 125 \text{ GeV} \), the model can also accommodate it and still provides light DM with the right relic abundance via the \( Z' \) contribution. In contrast, the simplest real-darkon model with \( m_D < m_h/2 \) would be disfavored by such a discovery \([23]\). More generally, depending on the parameters, our model can explain a partially hidden Higgs, while leaving the ratios of branching fractions to SM particles essentially the same. For \( 2m_{D_1,D_R} > m_h \), the Higgs decay pattern would be SM-like.

The recently discovered particle of mass near \( 125 \text{ GeV} \) has properties consistent with those of an SM Higgs, based on the LHC information reported so far \([1]\). However, the present data still allow the new boson to have

![Figure 3: Darkon-nucleon scattering cross-section corresponding to (a) \( J = 10^3 J_h \) (orange regions) and (b) \( J = 10 J_h \) (red regions) for \( m_h = 125 \text{ GeV} \) and \( m_{Z'} = 2m_D \). The darker (lighter) portions of the predictions come from the contributions of \( h \) alone (both \( h \) and \( Z' \)). The predictions are compared to 90\%-CL upper limits from CDMS, XENON10, and XENON100 \([28]\), as well as the 90\%-CL signal (purple) region suggested by CoGeNT, a (gray) patch compatible with DAMA modulation signal at the 5\( \sigma \) level, and two 2\( \sigma \)-confidence (blue) areas representing CRESST-II data \([29]\). Also plotted is the XENON1T projected sensitivity \([30]\).](image)

![Figure 4: Total branching ratios of \( h \rightarrow D_ID_I + D_RD_R \) versus \( m_D \) corresponding to \( J = 10J_h \) and \( 10^3 J_h \) for \( m_h = 125 \text{ GeV} \).](image)
a branching ratio into invisible particles of up to a few tens of percent. All this is compatible with one of the possibilities discussed above, where the $Z'$ contribution greatly dominates the relic density. Obviously, future measurements at the LHC and DM direct searches together will test such a scenario within our model.

V. SUMMARY AND DISCUSSION

Anticipating the upcoming results of the Higgs hunt at the LHC and direct searches for DM in various underground experiments, we have considered a simple model possessing only a small number of nonstandard particles, including scalar DM—the darkon—which is stabilized by a $Z_2$ symmetry naturally arising from the spontaneous breaking of a gauged $U(1)_{B-L}$ group. The associated gauge boson, $Z'$, yields new effects on the DM relic abundance in the resonant case, $m_{Z'} \simeq 2m_D$. We use the measurements of $e^+e^- \rightarrow \ell^+\ell^-$ and $pp \rightarrow \ell^+\ell^-X$ ($\ell = \mu, \tau$) and the relic density data to constrain the new gauge coupling. Subsequently, we explore different combinations of the Higgs- and $Z'$-mediated contributions to darkon annihilation subject to the relic density constraints and evaluate the corresponding darkon-nucleon scattering cross-section, $\sigma_{DN}$, compared to direct search results. If the $Z'$ effect dominates the darkon relic density, $\sigma_{DN}$ tends to evade the present limits and may come mainly from the $Z'$-mediated contribution, depending on the darkon mass. We also discuss some implications of the invisible decay of the Higgs if the darkon is sufficiently light, $m_D < m_h/2$. In that case, we find that the Higgs invisible decay branching-fraction is still significant for most of the allowed $m_D$ values even if the Higgs contribution to darkon annihilation is very small compared to the $Z'$ contribution. This result is important because the invisible branching fraction of the Higgs-like particle recently observed at the LHC has now been estimated from current data to reach up to a few tens of percent, which implies, within the context of our model, the necessity of a highly dominant $Z'$ effect on the DM relic density. In general, the parameter space of the model has enough room to accommodate various potential outcomes of the ongoing Higgs and DM direct searches, which will therefore probe the model further. The LHC can offer additional tests via processes that can produce the $Z'$ and/or a darkon pair, the latter due to diagrams mediated by $h$ or $Z'$. The $Z'$ may be most detectable in final states containing a pair of charged leptons, whereas the darkon may be uncovered in events with missing energy.

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