Study of a WIMP dark matter model with the updated results of CDMS II

Lian-Bao Jia and Xue-Qian Li*

Department of Physics, Nankai University, Tianjin 300071, People's Republic of China

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The new observation of CDMS II favors low mass WIMPs. Taking the CDMS II new results as input, we consider a standard model (SM) singlet: the darkon as the dark matter (DM) candidate, which can be either a scalar, fermion, or vector. It is found that the simplest scenario of DM + SM conflicts with the stringent constraint set by the LHC data. New physics beyond the SM is needed, and in this work, we discuss an extended standard model $SU_L(2) \otimes U_Y(1) \otimes U(1)'$ where $U(1)'$ only couples to the darkon. The new gauge symmetry is broken into $U_{em}(1)$ and two neutral bosons Z^0 and Z', which results in mixtures of W_μ^3, B_μ, and X_μ. Following the literature and based on the CDMS data, we conduct a complete analysis to verify the validity of the model. The cross section of the elastic scattering between the darkon and nucleon is calculated, and the DM relic density is evaluated in the extended scenario. It is found that by considering the constraints from both the cosmology and collider experiments, one can reconcile all the presently available data only if Z' is lighter than Z^0.

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I. INTRODUCTION

Recently, the CDMS Collaboration reported that three WIMP-candidate events were observed [1] by using silicon detectors. With a final surface-event background estimate of $0.41^{+0.20}_{-0.08}(\mathrm{stat})^{+0.28}_{-0.24}(\mathrm{syst})$, they indicated that the highest likelihood occurs for a WIMP mass of 8.6 GeV and a spin-independent (SI) WIMP-nucleon cross section of 1.9×10^{-41} cm^2. This observation seems to contradict the results of XENON100 [2]. Hooper [3] reanalyzed the data of XENON100 and reached a different conclusion; namely, the two experimental results could be reconciled. Therefore, in this work, we take the CDMS results as input to study the dark matter. We will test a viable model proposed in the literature by checking whether both the astronomical observation and constraints from the collider experiments can be simultaneously satisfied in this scenario.

As is well known, none of the standard model (SM) particles can meet the criterion to stand as dark matter (DM) candidates. Many particles beyond the SM are proposed, for example, the primordial black holes, axions, heavy neutrinos, the lightest supersymmetric neutralino, etc. Among them, the darkon model, namely, a SM singlet scalar [4–9] which interacts with the SM particles by exchanging a Higgs boson only, probably is the simplest version for the dark matter candidates. The spin-independent cross section for the darkon-nucleon elastic scattering might be measured by the earth detectors. The typical recoil energy is $\Delta E_R \sim (\mu v)^2/m_A$, where μ is the dark matter-nucleus reduced mass, v is the DM velocity, and m_A is the target nucleus mass. The WIMPs with not very heavy masses will weaken the bounds in the detector search, and in this work, we are more concerned with the low mass WIMPs (mass around 10 GeV).

The successful operation of the LHC, where the Higgs boson signals have been observed [10,11], provides a possible means to directly detect the dark matter particles on Earth, if they indeed exist. This means that all the proposed dark matter candidates and the possible new interactions by which the DM particles interfere with our detector would withstand the stringent test on the earth colliders. Namely, if the proposed DM particles, especially the lighter ones, are not observed at the LHC as expected, then the concerned model fails or needs to be modified. As indicated in Ref. [7], if the mass of the darkon is lighter than half the Higgs mass, then the Higgs would decay into a darkon pair, which is a channel with invisible final products, and the simplest version of the scalar darkon + SM may fail. That is to say, if the darkon's coupling to the Higgs is not much smaller than 1, then a large partial width is expected, and it obviously contradicts the measured value of the invisible width of the SM Higgs. As a possible extension of the scalar darkon + SM version, the two-Higgs-doublet model was discussed in Refs. [7,8,12], and there seems to be a large parameter space to accommodate both the LHC data on the Higgs and the CDMS observations.

We also find that the scenario of the darkon + SM, no matter if the darkon is a scalar, fermion, or vector, definitely fails; thus, a new interaction beyond the SM is needed. Alternatively, we propose an interaction beyond the SM as the darkon + SM + an extra $U(1)'$. The extended gauge group $SU_L(2) \otimes U_Y(1) \otimes U(1)'$ breaks into $U_{em}(1)$, and two neutral bosons Z^0 and Z' are the result. Z^0 and Z' are mixtures of W_μ^3, B_μ, and X_μ, which is the gauge boson of the newly introduced $U(1)'$, while the photon remains massless. In this scenario, to be consistent with the CDMS and LHC data simultaneously, we should assume that the

*lixq@nankai.edu.cn

coupling between the darkon and Higgs boson is very small, and the interaction by exchanging the Higgs between the detector material and darkon can be safely ignored. Therefore, the possible decays of the Higgs boson into darkons are almost forbidden, and one cannot expect to measure the mode at the LHC at all. The scattering between the darkon and nucleons is due to exchanging the gauge bosons Z^0 and Z'. Definitely, such an interaction may also exist in the decays of quarkonia; i.e., if the measurement of heavy quarkonia, such as bottomonia, are very precise, one may observe their decays into invisible final products besides the SM neutrino-antineutrino pairs. However, it is estimated that the branching ratios for such decays of heavy quarkonia are too small to be reliably measured in any of our present facilities. Besides, when the bottomonia are lighter than the new invisible final products, these decays are also forbidden. Therefore, this proposed darkon $+$ SM $+$ $U(1)'$ is safe with respect to the present experimental constraints. Moreover, the observed relic density of dark matter in our Universe sets one more constraint on our model parameter space.

This work is organized as follows. After this introduction, we first consider the simple version of a scalar, fermionic, and vectorial darkon within the framework of the SM $+$ darkon, then we derive the formulas of the cross section between the nucleon and darkon, as well as the decay width of the Higgs into invisible darkons. We further derive the corresponding formulas for the aforementioned extended version of the darkon $+$ SM $+$ $U(1)'$. Then in the following section, we numerically evaluate the cross sections of darkon-nucleon elastic scattering with the two scenarios. We indicate that the simple version does not satisfy the constraint set by the LHC data as long as we take the CDMS data as input, but in the extended version, there is a parameter space to accommodate both of the experimental measurements. The last section is devoted to our brief summary and discussion.

II. DARKON + SM

In this work, as the CDMS data suggested, we focus on low mass WIMPs. The WIMP particle could be an $SU_c(3) \times SU_L(2) \times U_Y(1)$ singlet, i.e., either a scalar, fermionic, or vectorial darkon [7,13,14]. In the scenario of the darkon $+$ SM, the elastic scatting between the darkon and the detector material is realized via a t-channel Higgs exchange, as described in Fig. 1.

A. Scalar darkon

Let us first consider a scalar-type WIMP DM, namely, a scalar darkon. This type of DM has been discussed in Ref. [7], and here, for completeness, we first repeat some relevant procedures. The Lagrangian is written as [4–7]

$$\mathcal{L} = \mathcal{L}_{\mathrm{SM}} - \frac{\lambda_S}{4} S^4 + \frac{1}{2} \partial^\mu S \partial_\mu S - \frac{m_0^2}{2} S^2 - \lambda S^2 H^\dagger H. \quad (1)$$

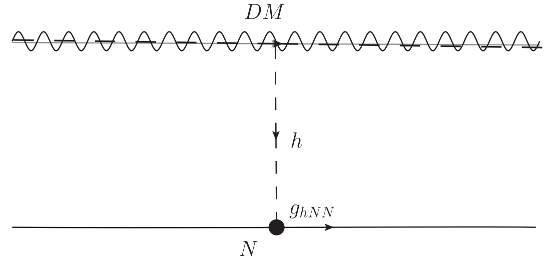

FIG. 1. The elastic scattering between dark matter and a nucleon with a Higgs boson exchange.

Here, λ_S, m_0, and λ are the free parameters to be determined by fitting the data. It has been indicated in earlier works that the scalar darkon field has no mixing with the Higgs field, and this can avoid fast decaying into the SM particles because dark matter particles must be sufficiently stable and survive from the big bang to today. From Eq. (1), the SM singlet scalar darkon can be further written as

$$\mathcal{L}_S = -\frac{\lambda_S}{4} S^4 + \frac{1}{2} \partial^\mu S \partial_\mu S - \frac{m_0^2 + \lambda v^2}{2} S^2 - \frac{1}{2} \lambda S^2 h^2 - \lambda v S^2 h. \quad (2)$$

The Higgs-nucleon coupling g_{hNN} is needed in calculating the scatting process, $\mathcal{L}_{hNN} = -g_{hNN} \bar{N} N h$. Here we adopt the value of g_{hNN} given by He *et al.* [8],

$$g_{hNN} \bar{N} N = \langle N | \frac{k_u}{v} (m_u \bar{u} u + m_c \bar{c} c + m_t \bar{t} t) + \frac{k_d}{v} (m_d \bar{d} d + m_s \bar{s} s + m_b \bar{b} b) | N \rangle, \quad (3)$$

and $g_{hNN} \simeq 1.71 \times 10^{-3}$. The cross section of the scalar DM-nucleon elastic scatting is [4–6]

$$\sigma_{\mathrm{el}} \simeq \frac{\lambda^2 v^2 g_{hNN}^2 m_N^2}{\pi (p_D + p_N)^2 m_h^4}. \quad (4)$$

Here, p_D, p_N are the momenta of the initial DM and nucleon. For the low energy elastic scatting, $(p_D + p_N)^2 \simeq (m_D + m_N)^2$, and m_D, m_N are masses of the DM and nucleon, respectively. Substituting the darkon mass 8.6 GeV and the cross section 1.9×10^{-41} cm^2 as given by CDMS II into Eq. (4), we can fix the effective coupling of the Higgs darkon.

The Higgs signals have been observed at the LHC [10,11] and $m_h = 125$ GeV, so by the data of CDMS, $\lambda \approx 0.148$ is determined. The partial width of the Higgs decaying into two scalar darkons is

$$\Gamma_{h \to SS} = \frac{\lambda^2 v^2}{8 \pi m_h} \sqrt{1 - \frac{4 m_D^2}{m_h^2}}. \quad (5)$$

Substituting the Higgs mass into the equation, $\Gamma_{h\to SS} \approx 0.418$ GeV is obtained. The main decay channel in the SM is $h \to b\bar{b}$. In the Born approximation, the width of this channel is [15,16]

$$\Gamma_{\rm Born}(h \to b\bar{b}) = \frac{3G_F}{4\sqrt{2}\pi} M_h m_b^2 \beta_b^3. \tag{6}$$

Here, $\beta = \sqrt{1 - 4m_b^2/M_h^2}$, and G_F is the Fermi coupling constant. With $G_F = 1.16 \times 10^{-5}\ \mathrm{GeV}^{-2}$ and $m_b(\overline{\rm MS}) = 4.18$ GeV, we can obtain $\Gamma_{\rm Born}(h \to b\bar{b}) \approx 0.00427$ GeV. Thus, the branching ratio $B_h \to$ invisible would be too large.

B. Fermionic and vectorial darkons

In the spin-$\frac{1}{2}$ darkon + SM scenario, the effective interaction can be written as

$$\mathcal{L}_{\rm int} = -\lambda \bar{\psi}_D \psi_D h. \tag{7}$$

The cross section of the low energy fermionic darkon-nucleon elastic scattering is

$$\sigma_{\rm el} \simeq \frac{\lambda^2 m_D^2 g_{hNN}^2 m_N^2}{\pi (p_D + p_N)^2 m_h^4}. \tag{8}$$

The partial width of the Higgs decaying into two darkon spinors is

$$\Gamma_{h\to DM} = \frac{\lambda^2 m_h}{8\pi} \sqrt{1 - \frac{4m_D^2}{m_h^2}}. \tag{9}$$

In this case, the invisible decay width is unbearably large when λ is at the order of unity. This means that such spin-$\frac{1}{2}$ darkon + SM scenario must also be abandoned.

For a vectorial darkon, the effective Lagrangian can be written as

$$\mathcal{L}_{VH} = \lambda V^\mu V_\mu H^\dagger H. \tag{10}$$

The cross section of the elastic scattering between a vectorial darkon and a nucleon via a Higgs boson exchange is

$$\sigma_{\rm el} \simeq \frac{\lambda^2 v^2 g_{hNN}^2 m_N^2}{\pi (p_D + p_N)^2 m_h^4}. \tag{11}$$

The numerical results for the vectorial darkon are similar to the two cases above for the scalar and fermionic darkons; namely, with the darkons possessing a low mass of order 10 GeV and the spin-independent cross section as determined by the CDMS data, the partial width of the Higgs decaying into invisible final products would be too large to be tolerated.

The above results indicate that the simplest scenario of the darkon + SM, no matter if the SM singlet darkon is a scalar, fermion, or vector, cannot reconcile the observation of the CDMS and LHC data. Then, one should invoke an extended version of the SM, i.e., a darkon + SM + BSM (beyond standard model) scenario. However, the BSM, which can be applied to explain the CDMS observation and the LHC data simultaneously, is a problem. There are many different proposals, and below we will investigate a naturally extended version of the SM; i.e., we introduce an extra $U(1)'$ gauge field, which is broken, and a new vector boson Z' is induced.

III. DARKON + SM + $U(1)'$

For the low mass darkon model, the simple version, darkon + SM, where darkons interact with the SM particles in the detector by exchanging a Higgs boson at the t channel, definitely fails to reconcile the observation of the CDMS and the LHC data, and, therefore, it needs to be modified. To tolerate the CDMS and LHC data, besides the two-Higgs-doublet model mentioned above, and, alternatively, Ref. [17] discussed the sneutrino dark matter that interacts dominantly with the detector material via exchanging the SM Z boson.

In this work, we study the effects of an extended SM by adding an extra $U(1)'$ [18–24], which only interacts with the darkons (no matter if they are scalar, fermionic, or vectorial darkons) in the gauge group as $SU_L(2) \otimes U_Y(1) \otimes U(1)'$ whose gauge bosons are, respectively, $W_\mu^\pm$, W_μ^3, B_μ, and X_μ (for more discussions about this model, see, e.g., Refs. [25–29]). The extended symmetry later breaks into $U_{em}(1)$. As a consequence, besides the regular charged $W^\pm$, two neutral gauge bosons Z^0 and Z' gain masses after the symmetry breaking, while the photon remains massless.

It is noted that a small mixing between the SM Z and X results in the physical mass eigenstates Z^0 and Z'. Since the mixing is required to be very small, the resultant $Z^0 = \cos\varphi Z + \sin\varphi X$ is almost the SM Z boson, whereas Z' is overwhelmingly dominated by X. Concretely, after $SU(2)_L \times U(1)_Y \times U(1)'$ breaking, one has

$$\begin{pmatrix} A_\mu \\ Z_\mu^0 \\ Z'_\mu \end{pmatrix} = \begin{bmatrix} \cos\theta_w & \sin\theta_w & 0 \\ -\sin\theta_w \cos\varphi & \cos\varphi\cos\theta_w & \sin\varphi \\ \sin\theta_w \sin\varphi & -\sin\varphi\cos\theta_w & \cos\varphi \end{bmatrix} \begin{pmatrix} B_\mu \\ W_\mu^3 \\ X_\mu \end{pmatrix}. \tag{12}$$

Assuming X_μ of $U(1)'$ only couples to the darkon but not to the SM particle, whereas Z_μ only couples to the SM particles, the interaction between the darkon and SM particles must be realized via the small mixing. Namely, the effective interaction amplitude between the darkon and protons or neutrons in the earth detector must be proportional to $\sin\varphi \cdot \cos\varphi$. To be consistent with the experiments, φ should be very small, i.e., $\sin\varphi \ll 1, \cos\varphi \sim 1$.

Since the new effective vertex $V_{DV(A)D}$ is a coupling between the scalar, fermionic, or vectorial darkon with the gauge boson, the Lorentz structure is well determined even though the coupling constants might be model dependent. Freytsis and Ligeti [30] listed all the possible operators and indicated which one(s) was suppressed by q^2 or v^2 where q was the exchanged momentum and v was the speed of the dark matter relative to the earth detector. Thus, in this work, we are only concerned with the unsuppressed spin-independent scattering processes which may correspond to the recently observed events. Below, we will focus on the fermionic darkon and give all the details, but for completeness, we also briefly discuss the cases for the scalar and vectorial darkons.

A. Fermionic darkon

Let us consider the fermionic darkon first. The axial-vector component of the gauge boson may induce a fermionic darkon-nucleon interaction, which is not suppressed by q^2 or v^2, even though this coupling would result in a spin-dependent cross section [30]. For easily handling, here we consider a right-handed darkon with the vertex $i\lambda\gamma^\mu \frac{1+\gamma^5}{2}$ to interact with the SM particles via exchanging the Z boson. The darkon-nucleon elastic scattering cross section is calculated for two cases $m_{Z'} \gg m_{Z^0}$ and $m_{Z'} \ll m_{Z^0}$, respectively, and the corresponding DM relic density is also computed.

1. The case of $m_{Z'} \gg m_{Z^0}$

In this case, the darkon-nucleon elastic scattering occurs mainly via exchanging Z^0. The fugacity speed of the WIMP is about 220–544 km/s [31]. For the low energy Z^0-nucleon interaction, the hadronic matrix element can be expressed as [32–34]

$$\langle p', s'|J_\mu^{Z^0}|p, s\rangle = \sqrt{\frac{G_F}{\sqrt{2}}}\bar{\mathcal{U}}_N(p', s')\left[G_A^z\gamma_\mu\gamma^5 + F_1^z\gamma_\mu + F_2^z\frac{i\sigma_{\mu\nu}q^\nu}{2M_N}\right]\mathcal{U}_N(p, s). \tag{13}$$

Here, $\mathcal{U}_N$, M_N, and q are the nucleon's wave function, mass, and momentum transfer, respectively. G_A^z, F_1^z, and F_2^z are the relevant form factors. Those form factors can be determined by the data of the elastic scattering between the neutrino and nucleon, since for this neutral current scattering process, only the Z^0 exchange is dominant (the new boson Z' is suppressed by a factor of $\sin^4\varphi$ in this process).

Here we adopt the way given in Refs. [32,33] to define the form factors. In the form of quark currents, the hadronic matrix element is written as

$$\langle p', s'|J_\mu^{Z^0}|p, s\rangle = \sqrt{\frac{G_F}{\sqrt{2}}}\bar{\mathcal{U}}_N(p', s')\sum_i[\bar{q}_i\gamma_\mu(1-\gamma^5)t_z q_i - 2Q_i\sin^2\theta_w\bar{q}_i\gamma_\mu q_i]\mathcal{U}_N(p, s). \tag{14}$$

The form factors are written as

$$G_A^z = -\frac{G_A^3\tau_3}{2} + \frac{G_A^s}{2}, \tag{15}$$

$$F_1^z = (1-2\sin^2\theta_w)F_1^3\tau_3 - 2\sin^2\theta_w F_1^1 - \frac{F_1^s}{2}, \tag{16}$$

$$F_2^z = (1-2\sin^2\theta_w)F_2^3\tau_3 - 2\sin^2\theta_w F_2^1 - \frac{F_2^s}{2}, \tag{17}$$

where the isospin factor $\tau_3 = +(-)$ for the proton (neutron) and

$$F_j^1 = \frac{F_j^p + F_j^n}{2}, \tag{18}$$

$$F_j^3 = \frac{F_j^p - F_j^n}{2}, \tag{19}$$

with $j = 1, 2$.

Defining $Q^2 = -q^2$, since $Q^2/m_N^2 \ll 1$, for the darkon-nucleon scattering via exchanging the Z^0 boson, we can set the values of the form factors at $Q^2 = 0$. At $Q^2 = 0$, $F_1^p = 1$, $F_1^n = 0$, $F_2^p = 1.7928$, $F_2^n = -1.9130$ [33]. In the limit of $Q^2 = 0$, the parameters corresponding to the strange part are $G_1^s(0) = \Delta s$, $F_1^s(0) = 0$, $F_2^s(0) = \mu_s$ [33–35], and here we take the fitted results $G_1^s(0) = -0.15 \pm 0.07$, $F_1^s(0) = 0$, $F_2^s(0) = 0$, $M_A = 1.049 \pm 0.019$, ($\chi^2 = 9.73$ at 13 d.o.f.) [33,34]. The Particle Data Group (PDG) average value of G_A^3 is $G_A^3 = 1.2701 \pm 0.0025$ [36]. So, at $Q^2 = 0$, the form factors are $G_A^z \approx -0.710\,(0.560)$ for the proton (neutron) and $F_1^z = 0.5–2\sin^2\theta_w$ (-0.5) for the proton (neutron). The contribution from the F_2^z term is suppressed at $Q^2 = 0$. When considering the conservation of the vector currents and just using the valence quarks in the nucleon, the same result can be obtained for the vector form factor F_1^z.

As the darkon is nonrelativistic, in the limit $\frac{p^\mu}{m} \to (1, \varepsilon)$, the darkon-nucleon elastic scattering cross section with the Z^0 exchanged at the t channel can be written as

$$\sigma_{\rm el} \simeq \frac{\sqrt{2} G_F \lambda^2 \sin^2\varphi m_D^2 m_N^2 (3 G_A^{z\,2} + F_1^{z\,2})}{4\pi (p_D + p_N)^2 m_{Z^0}^2}. \qquad (20)$$

It is noted that G_A^z is spin dependent, and F_1^z is spin independent. For large mass target nuclei, such as the silicon, germanium, and xenon targets, the spin-independent interaction is enhanced by the atomic number A^2 (but not exactly, see below for details) in the target nucleus, so the spin-independent interaction is more sensitive than the spin-dependent case, as discussed in Ref. [30]. Thus, we can drop the spin-dependent term G_A^z but just keep the spin-independent term F_1^z for large mass target nuclei scattering. For the proton, $F_1^z(p) = F_1^z(p) = 0.5–2\sin^2\theta_w \approx 0.038$, while for the neutron $F_1^z(n) = -0.5$. Thus, the darkon-neutron scattering is dominant, and the scattering cross section of the darkon nucleus via exchanging a neutral gauge Z boson should be proportional to $(A - Z)^2$ instead of A^2. Thus, a factor of about 0.25 might exist, and when analyzing the data to extract the information about the dark matter-nucleon interaction, this factor should be considered.

Substituting the CDMS II results for the darkon-neutron elastic scattering, $m_D \sim 8.6$ GeV and the elastic cross section $\sigma_{\rm el} \sim 4 \times 1.9 \times 10^{-41}$ cm^2 into the relevant formulas, we obtain $\lambda^2 \sin^2\varphi \approx 6.88 \times 10^{-3}$. To require the coupling constant $\alpha_D = \frac{\lambda^2}{4\pi} < 1$, the upper limit of λ is $\sqrt{4\pi}$.

In fact, the LEP data set a stringent constraint on the coupling and mixing. The width of Z^0 decaying into invisible products is $\Gamma(\text{invisible}) = 499.0 \pm 1.5$ MeV [36]. It is assumed in our scenario that after subtracting the main contribution of neutrinos from the measured width, the rest can be attributed to the darkon products. Thus, we can use the data to estimate the range of φ with some unavoidable uncertainties. The width of Z^0 decaying into a darkon pair is formulated as

$$\sin^2\varphi \Gamma_D = \frac{\lambda^2 \sin^2\varphi (m_{Z^0}^2 - m_D^2)}{24\pi m_{Z^0}} \sqrt{1 - \frac{4 m_D^2}{m_{Z^0}^2}}. \qquad (21)$$

Then, the total width of Z^0 decaying into invisible products is

$$\cos^2\varphi \Gamma_{\nu\bar{\nu}} + \sin^2\varphi \Gamma_D \le \Gamma_{\nu\bar{\nu}} + \sin^2\varphi \Gamma_D \approx 505.7 \text{ MeV}, \qquad (22)$$

and this value is larger than the experimentally measured value (the central value) for invisible products. If the mixing angle $\sin^2\varphi$ is reduced to an order of 0.01, this could satisfy the LEP data. However, this mixing angle is too large to be accepted because the SM electroweak sector would be seriously affected to conflict with all the previous well-done measurements.

Another constraint comes from the observed density of dark matter in our space.

The motion of the darkon is nonrelativistic; the invariant mass of a darkon pair can be approximated as $\sqrt{s} \simeq 2m_D$ where m_D is the darkon mass. In order to get the DM relic density, we need to calculate the DM annihilation cross section. In Ref. [6], the scalar-mediated (Higgs) $2 \to 2$ annihilation cross section of the DM pair into the SM particles is given. However, as discussed in the Introduction, we choose an alternative scenario where the coupling of the Higgs boson with the darkon is too small to make any substantial contributions to the darkon-nucleon scattering and the dark matter annihilation.

Here, the annihilation cross section of the darkons is dominated by the process that a darkon pair annihilates into a virtual gauge boson (Z^0 or Z'), which later transits into SM final states. Considering the case that the intermediate boson has a narrow width compared with its mass at the pole, the cross section is written as

$$\sigma_{\rm ann} = \frac{1}{2}\sigma_{\rm ann}^{\rm Dirac} = \frac{1}{2\beta_i (2s_1+1)(2s_2+1)} \frac{1}{(s-M^2)^2 + M^2\Gamma^2} \lambda^2 \sin^2\varphi \cos^2\varphi \times \left[2(s - m_D^2) \frac{\tilde{\Gamma}_f}{\sqrt{s}} + \left(\frac{s}{M^2} - 1\right)^2 \frac{m_{Z^0}^2 G_F}{2\sqrt{2}} \frac{N_c \beta_f c_a^2 m_f^2 m_D^2}{\pi s} \right]. \qquad (23)$$

A factor of $\frac{1}{2}$ appears for the fermionic dark matter, which is composed of a particle and antiparticle simultaneously, and the annihilation only occurs between the particle and its antiparticle (similarly, the factor of $\frac{1}{2}$ exists for the complex scalar DM, while this factor is equal to 1 for real scalar, Majorana fermionic DM). $s = (p_1 + p_2)^2$, M is the mass of the intermediate boson, and Γ is the total width of the intermediate boson. s_1, s_2 are the darkon spin projections. $\tilde{\Gamma}_f$ is the rate of the virtual boson transiting into the SM fermions (quarks or leptons); to obtain it, one only needs to replace the intermediate boson mass by $\sqrt{s}$ in the calculations. N_c is the color factor. c_a is the axial-vector current parameter, here $c_a^2 = 1$. $\beta_i = \sqrt{1 - 4m_D^2/s}$ and $\beta_f = \sqrt{1 - 4m_f^2/s}$ are the kinematic factors.

In the case of $m_{Z'} \gg m_{Z^0}$, the annihilation of a darkon pair into SM particles is dominated by Z^0 with the mixing component, namely, via darkon + darkon $\to Z^0 \to$ SM. Using Eq. (23), we can get the annihilation cross section. The DM relic density Ω_D is determined by the thermal dynamics of the big bang cosmology. The approximate values of the relic density and freeze-out temperature are [37,38]

$$\Omega_D h^2 \simeq \frac{1.07 \times 10^9 \text{ GeV}^{-1} x_f}{\sqrt{g_*} m_{Pl} \langle \sigma_{\rm ann} v_{\rm rel} \rangle}, \qquad (24)$$

$$x_f \simeq \ln \frac{0.038 g m_{Pl} m_D \langle \sigma_{\rm ann} v_{rel} \rangle}{\sqrt{g_* x_f}}. \qquad (25)$$

Here, h is the Hubble constant in units of 100 km/(s · Mpc), and $m_{Pl} = 1.22 \times 10^{19}$ GeV is the Planck mass. $x_f = m_D/T_f$ with T_f being the freezing temperature, g_* is the number of relativistic degrees of freedom with masses less than T_f. $\langle \sigma_{\rm ann} v_{\rm rel} \rangle$ is the thermal average of the annihilation cross section of the DM pair transiting into SM particles, v is the relative speed of the DM pair in their center-of-mass frame, and g is the number of degrees of freedom of the DM. In this work, the DM particle is assumed to be the darkon, which we describe above. The thermal average of the effective cross section is [39]

$$\langle \sigma_{\rm ann} v_{\rm rel} \rangle = \frac{1}{8 m_D^4 T K_2^2(\frac{m_D}{T})} \int_{4m_D^2}^{\infty} ds \sigma_{\rm ann} \sqrt{s}(s - 4m_D^2) \times K_1\left(\frac{\sqrt{s}}{T}\right), \tag{26}$$

where $K_i(x)$ is the modified Bessel functions of order i.

We calculate the cross section of the low mass darkon pairs (the mass of the darkon is supposed to be of the order 10 GeV) annihilating into SM leptons and quarks (except the top quark) via Z^0 exchange. x_f is obtained by solving Eq. (25) iteratively. The effective degrees of freedom g_* vary with the freeze-out temperature T_f, and we take the data of the Gondolo-Gelmini effective degrees of freedom in MicrOMEGAs 3.1 at $T_{\rm QCD} = 150$ MeV [40]. For $m_D \sim 8.6$ GeV, the DM density is $\Omega_D h^2 \approx 0.593$. The current PDG value for cold DM density is $\Omega_{\rm cdm} h^2 = 0.111(6)$ [36]. Thus, in the case of $m_{Z'} \gg m_{Z^0}$, the DM relic density is superabundant. Therefore, this scenario is not consistent with both the LEP data and the observed DM relic density, so that should be abandoned.

Below we turn to another possibility that $m_{Z'} \ll m_{Z^0}$.

2. The case of $m_{Z'} \ll m_{Z^0}$

Now, let us consider the case of $m_{Z'} \ll m_{Z^0}$. If the pole mass of Z' is just slightly above $2m_D$, the annihilation cross section of the darkon pair can be enhanced. The darkon-nucleon elastic scattering occurs mainly via exchanging Z' in this case, and the cross section is similar to the case for $m_{Z'} \gg m_{Z^0}$ and can be rewritten as

$$\sigma_{\rm el} \simeq G_F \frac{m_{z^0}^2}{m_{z'}^2} \frac{\sqrt{2}\lambda^2 \sin^2\varphi m_D^2 m_N^2 (3G_A^{z\,2} + F_1^{z\,2})}{4\pi (p_D + p_N)^2 m_{z'}^2}. \tag{27}$$

Taking the CDMS II results for the darkon-neutron elastic scattering as our input, we get $\lambda^2 \sin^2\varphi \approx 6.88 \times 10^{-3} \times (m_{z'}^4/m_{z^0}^4)$. As $\cos\varphi \sim 1$, the width of Z' decaying into a darkon pair is

$$\Gamma_D{}' \simeq \frac{\lambda^2 (m_{z'}^2 - m_D^2)}{24\pi m_{z'}} \sqrt{1 - \frac{4m_D^2}{m_{z'}^2}}. \tag{28}$$

For the LEP constraint, using Eq. (21) and rewriting Eq. (22), we can obtain that when $(m_{z'}^4/m_{z^0}^4) < 0.167$, the width of Z^0 decaying into the neutrinos plus darkons is within the experimental tolerance range. This can be satisfied when Z' is lighter than half of the Z^0 mass.

At the leading order, the annihilation of a darkon pair into SM particles is determined by Z' with the mixing component, and the cross section is calculated by Eq. (23). When $m_{Z'} < 2m_D$, the annihilation of a darkon pair into SM particles can also pass the constraints set by the aforementioned collider experiment and astronomical observation.

Define $2m_D/m_{Z'} = \xi$. By fitting the data, in the case where the Z' mass is near $2m_D$, we obtain $\lambda^2 \sin^2\varphi \simeq 8.7 \times 10^{-6}$ ($\xi = 1$) in the darkon-neutron SI elastic cross section. The dependence of the elastic scatting cross section on m_D is shown in Fig. 2, where m_D varies within a range of 5 GeV $\leq m_D \leq$ 12 GeV, and ξ takes the values of 0.7, 1, and 1.25. $\lambda^2 \sin^2\varphi = 1 \times 10^{-7}$ is given as a comparison. For $m_D \sim 8.6$ GeV, fitting the results of CDMS, we have $\lambda^2 \sin^2\varphi \simeq 8.7 \times 10^{-6}/\xi^4$.

The dependence of the darkon relic density $\Omega_D h^2$ on ξ ($2m_D/m_{Z'}$) is depicted in Fig. 3, where m_D is set to be 8.6 GeV and ξ varies from 0.55 to 1.35. $\lambda^2 \sin^2\varphi \simeq 8.7 \times 10^{-6}/\xi^4$. The solid square and empty dotted curves are for $\lambda = 0.5$ and 1.0, respectively. When $\xi > 1$, the curve for $\lambda = 0.5$ is close to the curve for $\lambda = 1.0$. It can be seen that there is a parameter space allowed by the present data.

As a comparison, the dependence of the darkon relic density $\Omega_D h^2$ on m_D and $\lambda^2 \sin^2\varphi$ is shown in Fig. 4 where m_D is set as 6, 8.6, and 10 GeV and $\lambda^2 \sin^2\varphi = 8.7 \times 10^{-6}$ and 1×10^{-7}. We take $\lambda = 1$ here and let ξ vary from 0.7 to 1.25. The solid curves and empty curves are for the case $\lambda^2 \sin^2\varphi = 8.7 \times 10^{-6}$ and 1×10^{-7}, respectively. The square curves, dotted curves, and diamond curves

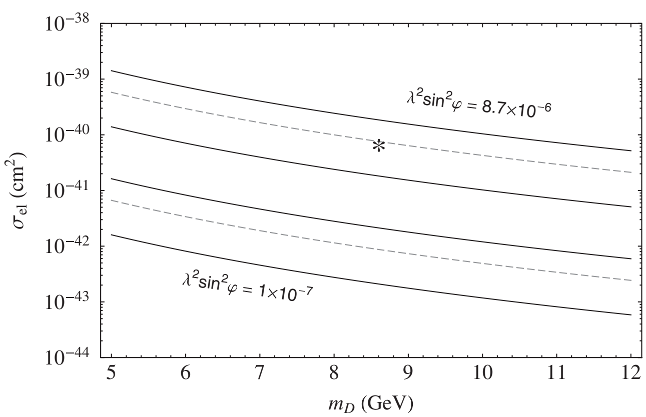

FIG. 2. Darkon-neutron SI elastic cross section $\sigma_{\rm el}$ as a function of the darkon's mass. m_D varies in a range 5 GeV $\leq m_D \leq$ 12 GeV. $2m_D/m_{Z'} = \xi$, for ξ equal to 0.7, 1, 1.25. $\lambda^2 \sin^2\varphi = 8.7 \times 10^{-6}, 1 \times 10^{-7}$. The dashed curve is in the case $\xi = 1$, the upper solid curve is $\xi = 1.25$, and the lower solid curve is $\xi = 0.7$. The $*$ is the reserved CDMS II observed event.

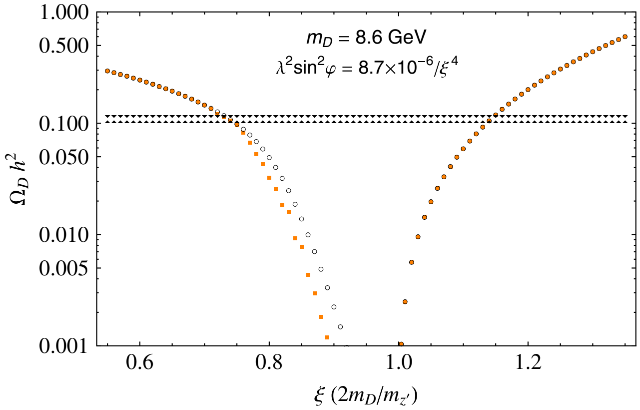

FIG. 3 (color online). The darkon relic density $\Omega_D h^2$ as a function of ξ $(2m_D/m_{Z'})$ near the Z' pole when $m_D = 8.6$ GeV, for ξ in a range from 0.55 to 1.35 and $\lambda^2\sin^2\varphi = 8.7\times10^{-6}/\xi^4$. The solid square curve is for the case of $\lambda = 0.5$, and the empty dotted curve is $\lambda = 1.0$. The triangle and reverse triangle curves are the cold dark matter relic density 0.111(6) today.

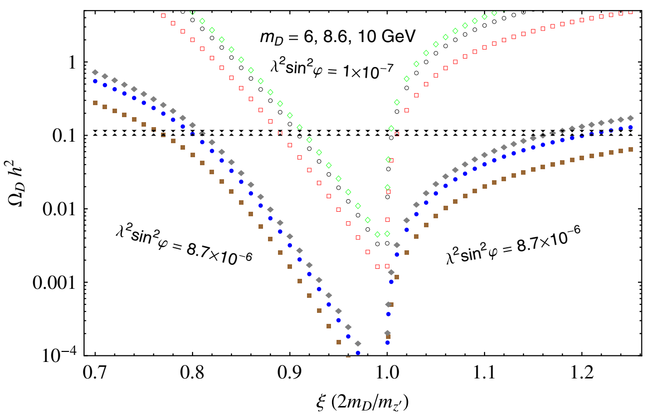

FIG. 4 (color online). The darkon relic density $\Omega_D h^2$ as a function of ξ $(2m_D/m_{Z'})$ near the Z' pole when $m_D = 6$, 8.6, and 10 GeV, for ξ values varying from 0.7 to 1.25. $\lambda = 1$ is taken here. The solid curves are for the case of $\lambda^2\sin^2\varphi = 8.7\times10^{-6}$, and the empty curves are for the case of $\lambda^2\sin^2\varphi = 1\times10^{-7}$. The square curves, dotted curves, and diamond curves (solid, empty) correspond to the case m_D equal to 6, 8.6, and 10 GeV, respectively. The triangle and reverse triangle curves are the cold dark matter relic density 0.111(6).

correspond to the case m_D equal to 6, 8.6, and 10 GeV, respectively.

B. Scalar and vectorial darkons

Now let us consider the scalar-darkon case. The effective vertex is a vector coupling $-i\lambda(k+k')^\mu$, as shown in Fig. 5 (left). As mentioned, the darkon-nucleon scattering induced by this interaction is an unsuppressed SI process [30]. In the limit $\frac{p^\mu}{m} \to (1,\epsilon)$, the darkon-nucleon elastic scattering cross section by exchanging Z^0 is written as

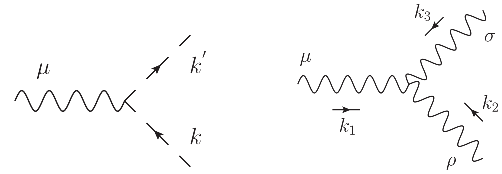

FIG. 5. Vertexes of scalar (left) and vectorial (right) darkons.

$$\sigma_{\rm el} \simeq \frac{\sqrt{2}G_F\lambda^2\sin^2\varphi m_D^2 m_N^2 F_1^{z\,2}}{\pi(p_D+p_N)^2 m_{Z^0}^2}. \tag{29}$$

$F_1^z = 0.5$–$2\sin^2\theta_w$ (-0.5) for the proton (neutron). Thus, the darkon-neutron scattering is dominant. Similar results can be obtained for the fermionic darkon case.

Instead, for the case where the Z' exchange is dominant, one can modify Eq. (29) by simply multiplying a factor of $m_{z^0}^4/m_{z'}^4$.

For the vectorial darkon, the vertex is $-i\lambda[g^{\mu\rho}(k_2-k_1)^\sigma + g^{\rho\sigma}(k_3-k_2)^\mu + g^{\sigma\mu}(k_1-k_3)^\rho]$, corresponding to the effective interaction $B_\mu^\dagger\partial^\nu B^\mu\bar{q}\gamma_\nu q$, as shown in Fig. 5 (right), which contributes an unsuppressed SI cross section. In the limit $\frac{p^\mu}{m} \to (1,\epsilon)$, the darkon-nucleon elastic scattering cross section with Z^0 exchange dominance can be written as

$$\sigma_{\rm el} \simeq \frac{\sqrt{2}G_F\lambda^2\sin^2\varphi m_D^2 m_N^2 F_1^{z\,2}}{\pi(p_D+p_N)^2 m_{Z^0}^2}. \tag{30}$$

In the case where the Z' exchange is dominant, the elastic cross section can be obtained by multiplying Eq. (30) by a factor of $m_{z^0}^4/m_{z'}^4$. The obtained result is the same as that for the scalar-darkon-nucleon elastic scattering.

It is noted that for the fermionic, scalar, and vectorial darkon-nucleon elastic scattering via the Z-boson exchange, there exists SI darkon-neutron scattering, which is not suppressed by either q^2 or v^2. In this case, the proton contributions are suppressed so that the main contributions to the SI scattering come from the interaction between the darkon and neutron. Therefore, the xenon target, which has more neutrons than protons, is more sensitive compared with the silicon and germanium targets. As explained above, if we accept the claim of XENON10 [41] and XENON100 [2] that for low energy WIMPs, null results have been obtained, then the CDMS results should be dubious. However, as suggested by Hooper [3], a reanalysis may imply that the peculiar events observed at the XENON100 might be explained as dark matter candidates to be reconciled with the CDMS data. With the further progress of the XENON experiments, more information will be obtained for the low mass WIMPs with masses of order 10 GeV.

IV. CONCLUSIONS AND DISCUSSION

Taking the recent new results of the CDMS II experiments searching for WIMPs with masses of order 10 GeV as input and considering some constraints from the LHC, LEP, and astronomical observation, etc., we discuss a simple WIMP candidate: the darkon, which can be a scalar, fermion, or vector. We have found that in the simplest scenario of the standard model plus SM singlet DM (the darkon), one cannot simultaneously satisfy the CDMS II's observation and the LHC data, and this result is consistent with the former result implied in Ref. [7]. Thus, one must extend the SM to include new physics beyond the standard model. Here we consider the extended gauge group $SU_L(2) \otimes U_Y(1) \otimes U(1)'$, which later breaks into $U_{em}(1)$ to result in two heavy neutral gauge bosons Z^0 and Z'.

The darkon $+$ SM $+$ $U(1)'$ scenario must undergo stringent tests from the cosmology observation and the LHC data. Namely, all the CDMS II results, dark matter density in our Universe, and the data of Z^0 decaying into invisible products, which were obtained by LEP experiments, must not conflict.

Our numerical results indicated that in this scenario, all the constraints can be satisfied only if Z' is lighter than Z^0. Under this assumption, the darkon $+$ SM $+$ $U(1)'$ model can withstand all the constraints set by the presently available data. Moreover, it is noted that as long as $m_{Z'} \sim 2m_D$, the model can accommodate an even smaller scattering cross section and lighter darkons.

Indeed, we should further test the validity of this mechanism. If in the future, we can precisely measure the branching ratios of heavy quarkonia, such as botomonia decaying into invisible products or the invisible decays of the SM Z boson and Higgs sector physics, we would be able to determine which one, e.g., the decays of the two-Higgs-doublets mechanism or the extra $U(1)'$ gauge group, is more reasonable. We hope that in the future there will be more precise direct and indirect detection of dark matter.

In the world today, there are many laboratories other than the XENON and CDMS collaborations directly searching for dark matter. For example, the China Jin-Ping underground laboratory [42] just began the search, and the China Dark-Matter experiment is using a 1 kg Ge detector and will develop a 10 kg and 1 ton detector for the project. We are expecting that the worldwide cooperation will eventually reveal the epoch mystery.

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[1] R. Agnese *et al.* (CDMS Collaboration), Phys. Rev. Lett. **111**, 251301 (2013).
[2] E. Aprile *et al.* (XENON100 Collaboration), Phys. Rev. Lett. **109**, 181301 (2012).
[3] D. Hooper, J. Cosmol. Astropart. Phys. 09**2013** 035.
[4] V. Silveira and A. Zee, Phys. Lett. **161B**, 136 (1985).
[5] J. McDonald, Phys. Rev. D **50**, 3637 (1994).
[6] C. P. Burgess, M. Pospelov, and T. ter Veldhuis, Nucl. Phys. **B619**, 709 (2001).
[7] X. -G. He, T. Li, X.-Q. Li, and H.-C. Tsai, Mod. Phys. Lett. A **22**, 2121 (2007).
[8] X.-G. He, T. Li, X.-Q. Li, J. Tandean, and H.-C. Tsai, Phys. Rev. D **79**, 023521 (2009).
[9] Y. Mambrini, Phys. Rev. D **84**, 115017 (2011).
[10] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[11] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[12] X. -G. He and J. Tandean, Phys. Rev. D **88**, 013020 (2013).
[13] O. Lebedev, H. M. Lee, and Y. Mambrini, Phys. Lett. B **707**, 570 (2012).
[14] A. Djouadi, O. Lebedev, Y. Mambrini, and J. Quevillon, Phys. Lett. B **709**, 65 (2012).
[15] L. Resnick, M. K. Sundaresan, and P. J. S. Watson, Phys. Rev. D **8**, 172 (1973).
[16] J. R. Ellis, M. K. Gaillard, and D. V. Nanopoulos, Nucl. Phys. **B106**, 292 (1976).
[17] K. -Y. Choi and O. Seto, Phys. Rev. D **88**, 035005 (2013).
[18] D. E. Brahm and L. J. Hall, Phys. Rev. D **41**, 1067 (1990).
[19] B. de Carlos and J. R. Espinosa, Phys. Lett. B **407**, 12 (1997).
[20] J. Erler, P. Langacker, and T.-j. Li, Phys. Rev. D **66**, 015002 (2002).
[21] B. Kors and P. Nath, Phys. Lett. B **586**, 366 (2004).
[22] V. Barger, C. Kao, P. Langacker, and H. -S. Lee, Phys. Lett. B **600**, 104 (2004).
[23] B. Kors and P. Nath, J. High Energy Phys. 07 (**2005**) 069.
[24] P. Langacker, Rev. Mod. Phys. **81**, 1199 (2009).
[25] J. Kumar and J. D. Wells, Phys. Rev. D **74**, 115017 (2006).
[26] W.-F. Chang, J. N. Ng, and J. M. S. Wu, Phys. Rev. D **74**, 095005 (2006); W.-F. Chang, J. N. Ng, and J. M. S. Wu, Phys. Rev. D **79**, 039902(E) (2009).
[27] D. Feldman, Z. Liu, and P. Nath, Phys. Rev. D **75**, 115001 (2007).
[28] Y. Mambrini, J. Cosmol. Astropart. Phys. 09 (**2010**) 022.
[29] M. T. Frandsen, F. Kahlhoefer, S. Sarkar, and K. Schmidt-Hoberg, J. High Energy Phys. 09 (**2011**) 128.
[30] M. Freytsis and Z. Ligeti, Phys. Rev. D **83**, 115009 (2011).
[31] M. C. Smith, G. R. Ruchti, A. Helmi, R. F. G. Wyse, J. P. Fulbright, K. C. Freeman, J. F. Navarro, G. M. Seabroke *et al.*, Mon. Not. R. Astron. Soc. **379**, 755 (2007).

[32] G. Garvey, E. Kolbe, K. Langanke, and S. Krewald, Phys. Rev. C **48**, 1919 (1993).
[33] G. T. Garvey, Prog. Part. Nucl. Phys. **34**, 245 (1995).
[34] W. M. Alberico, S. M. Bilenky, and C. Maieron, Phys. Rep. **358**, 227 (2002).
[35] H. Forkel, Phys. Rev. C **56**, 510 (1997).
[36] J. Beringer *et al.* (Particle Data Group), Phys. Rev. D **86**, 010001 (2012).
[37] E. W. Kolb and M. S. Turner, *The Early Universe* (Westview Press, Boulder, CO, 1990).
[38] K. Griest and D. Seckel, Phys. Rev. D **43**, 3191 (1991).
[39] P. Gondolo and G. Gelmini, Nucl. Phys. **B360**, 145 (1991).
[40] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, arXiv:1305.0237; http://lapth.cnrs.fr/micromegas/.
[41] J. Angle *et al.* (XENON10 Collaboration), Phys. Rev. Lett. **107**, 051301 (2011).
[42] K.-J. Kang *et al.* (CDEX Collaboration), Front. Phys. China **8**, 412 (2013).