Light Dark Matter Models with Isospin Violation

Xin Gao,[1,*] Zhaofeng Kang,[1,†] and Tianjun Li[1,2,‡]

[1]*Key Laboratory of Frontiers in Theoretical Physics, Institute of Theoretical Physics, Chinese Academy of Sciences, Beijing 100190, P. R. China*

[2]*George P. and Cynthia W. Mitchell Institute for Fundamental Physics, Texas A&M University, College Station, TX 77843, USA*

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Abstract

Light dark matter (DM) models with isospin violation (ISV) and large DM-nucleon spin-independent cross section σ_n may provide a way to understand the confusing DM direct detection experimental results. Combining with stringent astrophysical and collider constraints, we can further deduce the DM properties. In light of general operator analyses, we systematically study the ISV operators with $f_n/f_p < 0$, and show that the required ISV $f_n/f_p \simeq -0.7$ must arise from the DM and first-family quark couplings. Then we discuss three kinds of the ISV models classified by the mediators: a light Z' gauge boson in the extra $U(1)_X$ model, a (approximate) spectator Higgs doublet, and color triplets. In particular, the spectator Higgs doublet model can explain the Tevatron CDF $W + jj$ anomaly simultaneously. In addition, although the $U(1)_X$ gauge boson Z' which has kinetic mixing with $U(1)_Y$ gauge boson only generates $f_n = 0$, we can combine it with the conventional Higgs Yukawa couplings to achieve the proper ISV. Especially, most of our models can address the recent GoGeNT annual modulation as well as all the other DM direct detection experiments. As a concrete example, we propose the $U(1)_X$ model where the $U(1)_X$ charged light sneutrino is the inelastic DM (iDM), and dominantly decays to the light dark states such as Z' ($M_{Z'}$ <1 GeV). With ISV it is consistent with all the DM direct detection experiments and satisfies all the other constraints.

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*Electronic address: xgao@itp.ac.cn

†Electronic address: zhfkang@itp.ac.cn

‡Electronic address: tli@itp.ac.cn

I. INTRODUCTION AND MOTIVATION

The way that dark matter (DM) interacts with nucleons is a puzzle by virtue of the absence of confirmative experiments results. The recent possible progress made on direct detections may shed light on it. On the one hand, the events from the DAMA/LIBRA [1] and CoGeNT [2] experiments imply a light DM (LDM) with mass around 8 GeV and a rather large spin-independent (SI) cross section with nucleon $\sigma_n \sim 2 \times 10^{-4}$ pb and $\sigma_n \sim 5 \times 10^{-5}$ pb, respectively. On the other hand, the null results from XENON [4] and CDMS [3] experiments challenge the CoGeNT/DAMA results. Moreover, for the CoGeNT/DAMA favored DM there may be some expected signals (but not found) from astrophysics or collider experiments. So how to reconcile these results in a natural way may guide us to identify some of DM properties.

DM-nucleus interactions with isospin-violation (ISV) may provide an interesting way to reconcile various direct detection experimental results [5]. ISV is not a novel phenomena [6, 7], which arises when the DM interactions with proton and neutron have different strengths, $f_p \neq f_n$. If $f_n/f_p < 0$ the amplitudes between the DM-proton and DM-neutron destructively interfere, leading to a cancellation in the DM-nucleus amplitude. The degree of cancellation varies with the concrete nucleus used in the detector target. In this way, the strongest constraints from XENON are substantially weakened by taking $f_n/f_p = -0.7$, especially, the CoGeNT/DAMA regions may overlap [5]. But this region is excluded by the CDMS-Ge experiment which use the same atom as the CoGeNT experiment.

However, if one only considers annual modulation (the observed CoGeNT annual modulation with significance of 2.8σ [8]), inelastic DM (iDM) scenario [9] is able to enhance modulation and reduce tension between CDMS-Ge and CoGeNT. The Ref. [10] found that taking $f_n/f_p = -0.7$ and the upper-limit quenching factor $Q_{Na} = 0.43$, one can address all the confusing experimental results consistently via an iDM ∼10 GeV with mass splitting $\delta \simeq 15$ keV. On the other hand, right now the direct detection experiments may have reached the level to distinguish f_n and f_p [11]. Once determined, they definitely convey information on the DM-quark interactions. In a word, it is worthy of studying the LDM models with ISV systematically.

Inspired by the aforementioned facts, in this work we make rather thorough analyses on the LDM models with ISV, focussing on the scalar and fermionic DMs. We outline the generic characters for the LDM models:

- Proper ISV $f_n/f_p \approx -0.7$. How to get this ISV is highly non-trivial. In fact, we shall show that the conventional mediators, such as the Higgs/Z/squarks fail to accommodate this value at least in the case of a single dominated mediator. So, we need to investigate the new mediators beyond them.

- Large DM-nucleon SI scattering cross section $\sigma_n \sim 10^{-2}$ pb. In the ISV models, the nucleus amplitude is reduced by destructively interference. So σ_n is required to be about 2 orders larger than the conventional scenario. However, such large cross section will bring tension (or even be excluded) with some astrophysical or collider constraints.

- Right DM relic density $\Omega_{\rm DM}h^2 \simeq 0.11$. We may relax this constraint by allowing possible new annihilation channels beyond the Standard Model (SM) fermions as final states. In actual model building, LDM usually can annihilate to light dark sector states, which will be very helpful sometime.

- The possible minimal models which may explain the other new physics as well. Especially, we try to account for the Tavetron physics such as the CDF W+2jets anomalies [12].

In this paper, we systematically study the ISV operators with $f_n/f_p < 0$ via the effective operaror analyses. In particular, we show that the required ISV $f_n/f_p \simeq -0.7$ must arise from the couplings between the DM and first-family quarks. Moreover, we propose three kinds of the ISV models classified by the mediators: a light Z' gauge boson in the extra $U(1)_X$ model, a (approximate) spectator Higgs doublet, and color triplets. Interestingly, the spectator Higgs doublet model can explain the Tevatron $W+jj$ anomaly simultaneously. In addition, combining the $U(1)_X$ gauge boson Z' which has kinetic mixing with $U(1)_Y$ gauge boson with the conventional Higgs Yukawa couplings, we also obtain the proper ISV. We emphasize that most of our models can address the recent GoGeNT annual modulation as well as all the other DM direct detection experiments. To be concrete, we propose the $U(1)_X$ model where the $U(1)_X$ charged light sneutrino is the inelastic DM (iDM), and dominantly decays to the light dark states such as Z' ($M_{Z'}$ <1 GeV). With ISV it is not only consistent with all the DM direct detection experiments but also satisfies all the other constraints.

The paper is organized as follows. In Section II, we make some general effective operator analyses for the dark matter models with ISV. In Section III, we present three kinds of the ISV models by considering the different mediators. In Section IV, a concrete model with sneutrino iDM is discussed. Conclusion and discussions are given in Section V, and finally some useful formulas are collected in the Appendices.

II. GENERIC OPERATOR ANALYSES

Let us address the setup for generic operator analyses which simplify the discussions. The ISV mainly comes from the different $U(1)_X$ charges for u_R and d_R quarks, or from the different Yukawa couplings between the DM and the first-family quarks.

A. Effective ISV Operators for the CoGeNT/DAMA Experiments

General anatomy of the interactions between the DM and SM fields involves a large class of effective operators $\mathcal{O}_{\rm DM}\bar{f}\Gamma f$ with $\Gamma = 1, \gamma_\mu, \gamma_5 ...$ where f is the SM fermion, and $\mathcal{O}_{\rm DM}$ denote the corresponding DM bilinear operators. However, they are greatly reduced if we are just interested in the operators that are relevant to the CoGeNT/DAMA experiments. And we can further recover their $SU(3)_C \times U(1)_{EM}-$UV completion by specifying the possible mediators that connect the DM and the SM fermions, please see the Appendix A for details.

At first, we pick out the operators generating SI cross sections, and only three of them are not non-relativistic (NR) suppressed [13, 14]:

$$\bar{q}q, \quad \bar{q}\gamma_\mu q, \quad (G^a_{\mu\nu})^2. \tag{1}$$

In a concrete model, only a subset of them maybe generated simultaneously. Obviously, the gluon operators can not generate ISV and then are dropped in our discussions. Converting to the operators on the DM-SM particle interactions, we obtain the scalar and vector interactions, which are respectively given in Eqs. (2) and (3). In this paper, we use a_q and b_q to describe the operator coefficients for the scalar and vector type interactions, respectively.

$$a_{q_i}\bar{\chi}\chi\bar{q}_i q_i, \qquad a_{q_i}\phi^\dagger\phi\bar{q}_i q_i, \tag{2}$$
$$b_{q_i}\bar{\chi}\gamma^\mu\chi\bar{q}_i\gamma_\mu q_i, \quad b_{q_i}\phi^\dagger\overleftrightarrow{\partial}^\mu\phi\bar{q}_i\gamma_\mu q_i, \tag{3}$$

where χ is a fermionic DM while ϕ is a scalar DM. If χ is Majorana or ϕ is real, then the corresponding vector interactions vanishes automatically. Operators $\bar{\chi}\gamma^5\chi\bar{q}q$ and $\bar{\chi}\gamma^\mu\gamma^5\chi\bar{q}\gamma^\mu q$ also generate SI cross section, nevertheless NR suppressed. In our notation, $q = u, d$ denotes the up- and down-type quarks respectively, with i the family index (all fields are written in the mass eigenstates). We will show that only the first-family quarks are crucial to produce ISV. To suppress the potential large flavor violation, the quark bilinears are supposed to be diagonal in the flavor space. But in some cases, we will find that this assumption is not hold so naively. The operator coefficients a_q and b_q are assumed to be suppressed by high threshold scale (e.g., mass scale of mediator) $\Lambda \gg M_{\rm DM}$ unless we specify.

Next, there are other operators that do not contribute to the leading order SI cross section in the NR limit, nevertheless usually generated together with the other operators in Eq. (2) and Eq. (3) in a UV completed theory. Interestingly, they are important to determine the relic density of DM. The indirect detection signals and the annihilation rates from these operators are collected in the Appendix A. To be concrete, we present these operators in the following

$$\bar{\chi}\gamma^5\chi\bar{f}\gamma^5 f, \quad \bar{\chi}\chi\bar{f}\gamma^5 f, \quad \bar{\chi}\gamma^5\chi\bar{f}f,$$
$$\bar{\chi}\gamma^5\gamma_\mu\chi\bar{f}\gamma^5\gamma^\mu f, \quad \bar{\chi}\gamma_\mu\chi\bar{f}\gamma^5\gamma^\mu f, \quad \bar{\chi}\gamma^5\gamma_\mu\chi\bar{f}\gamma^\mu f,$$
$$\bar{\chi}\sigma^{\mu\nu}\chi\bar{f}\sigma_{\mu\nu}f, \quad \bar{\chi}\sigma^{\mu\nu}\gamma^5\chi\bar{f}\sigma_{\mu\nu}f,$$
$$|\phi|^2\bar{f}\gamma_5 f, \quad \phi^\dagger\overleftrightarrow{\partial}_\mu\phi\bar{f}\gamma^\mu\gamma^5 f... \,. \tag{4}$$

For the LDM we only have to consider the SM states lighter than $M_{\rm DM}$ which the LDM can annihilate into. So Eqs. (2)-(4) give a general effective operator description on the DM models inspired by the CoGeNT/DAMA experiments. Having established the setup, in the following we shall investigate the ISV from scalar and vector interactions.

First let us consider scalar interactions given in Eq. (2). To construct the DM-nucleon effective operators from microscopic DM-quark interactions, we must calculate the quark bilinear matrix elements in the nucleon states. In the case of scalar interaction, they are given by

$$m_q\langle n|\bar{q}q|n\rangle = m_n f_{T_q}^{(n)}, \tag{5}$$

where n denotes either the proton or neutron, and m_q the quark mass. For the light quarks, the form factors $f_{T_u}^{(p)} = 0.020 \pm 0.004$, $f_{T_u}^{(n)} = 0.014 \pm 0.003$, $f_{T_d}^{(p)} = 0.026 \pm 0.005$, $f_{T_d}^{(n)} = 0.036 \pm 0.008$, and $f_{T_s}^{(p,n)} = 0.118 \pm 0.062$ [15]. The heavy quarks $q = c, b, t$ contribute to the nucleon mass through the triangle diagram [16] as follows

$$m_q\langle n|\bar{q}q|n\rangle = \frac{2}{27}m_n\left(1 - \sum_{q=u,d,s} f_{T_q}^{(n)}\right) \equiv \frac{2}{27}m_n f_{T_G}^{(n)}, \tag{6}$$

where $f_{T_G} \simeq 0.84$ and 0.83 for the proton and neutron respectively. Through the above two equations, we get the DM-nucleon effective couplings

$$f_n \propto a_n = m_n\left[\sum_{q=u,d,s} a_q\frac{f_{T_q}^{(n)}}{m_q} + \frac{2}{27}\left(1 - \sum_{q=u,d,s} f_{T_q}^{(n)}\right)\sum_{c,b,t}\frac{a_q}{m_q}\right] \equiv \sum_{q=\text{quarks}} B_q^n a_q, \tag{7}$$

where the dimensionless quantities $B_q^n \equiv f_{T_q}^{(n)} m_n/m_q$, encoding the ISV in the nucleus itself. We use the quark masses: $m_u = 0.002$ GeV, $m_d = 0.005$ GeV, $m_s = 0.095$ GeV, $m_c = 1.25$ GeV, $m_b = 4.2$ GeV, $m_t = 172.3$ GeV. Then we get

$$B_u^p \approx 9.3, \quad B_u^n \approx 6.5, \quad B_d^p \approx 5.1, \quad B_d^n \approx 7.1, \tag{8}$$
$$B_s^{p,n} \approx 1.2, \quad B_c^{p,n} \approx 0.05, \quad B_b^{p,n} \approx 0.015, \quad B_t^{p,n} \approx 0.00035. \tag{9}$$

By the way, in Eq. (7) $f_n = a_n$ for a fermionic DM. Whereas for s scalar DM a_q has dimension +1, and if the DM-nucleus SI scattering cross section written in the form of Eq. (B1), then actually $f_n = a_n/2M_{\rm DM}$. Hereafter we will absorb the $1/2M_{\rm DM}$ into a_q for the scalar DM.

Remarkably, from Eqs. (7)-(9) it is obvious that only the DM and u/d quark interactions break isospin effectively. Immediately, we draw the conclusion: <u>if scalar interactions account for CoGeNT/DAMA experiments, the DM and first-family quark interactions must give the predominant contribution</u>. Then the ISV $f_n = \mathcal{I} f_p$ transformed into quark level is simply given by

$$\frac{a_u}{a_d} \simeq \frac{\mathcal{I}B_d^p - B_d^n}{B_u^n - \mathcal{I}B_u^p} < 0, \tag{10}$$

and the ratio is about -0.77 for $\mathcal{I} = -0.7$ (throughout this work, we shall use this ISV as a referred value, as well the referred DM mass $M_{\rm DM} = 8$ GeV). Furthermore, the general effective DM-proton couplings with ISV can be organized in such a form

$$f_p \simeq \left(\frac{B_d^p B_u^n - B_d^n B_u^p}{B_u^n - B_u^p \mathcal{I}}\right) \times a_d \simeq \frac{10^{-5}}{\sqrt{\delta_C}}\,{\rm GeV}^{-2} \tag{11}$$

It is factorized into DM-quark effective coupling and a model-independent factor casted in the bracket, which takes value -2.5 for $\mathcal{I} \simeq -0.7$. We typically require DM-proton SI cross section $\sigma_p = 0.01$ pb, in turn f_p is determined to be around the value given above. By the way, it applies to both scalar and vector interactions in a notation given in Eq. (B1).

The above conclusion excludes the models with conventional Higgs mediator, *e.g.*, in the one/two Higgs doublet SM and in the Minimal Supersymmetric Standard Model (MSSM) where the Higgs fields not only dominantly mediate interactions but also generate the SM fermion masses. In such models, we have

$$\frac{a_{u_i/d_i}}{m_{u_i/d_i}} \propto \frac{1}{v_{u/d}}, \tag{12}$$

with $v_{u/d}$ the vacuum expectation value (VEV) of Higgs field $H_{u/d}$. It is independent on the mediator-quark couplings, consequently the second and third families give main contributions by virtue of larger form factor, thus the isospin is preserved. But if the Higgs field is a spectator Higgs field whose VEV is zero or very small, its Yukawa couplings with SM fermions are free parameters. And then the light quark contributions to SI crosse section can definitely exceed the contributions from heavy quarks.

Now turn to the vector interactions. The quark bilinear matrix elements in the nucleon states are greatly simplified by virtue of the conservation of the vector current, to which the sea quarks and gluons do not contribute. As a consequence, the effective interactions between DM and nucleons are simply given by

$$\mathcal{L}_{vec} = b_n \bar{\chi}\gamma_\mu \chi \bar{n}\gamma^\mu n, \quad b_n \phi^\dagger \overleftrightarrow{\partial}_\mu \phi \bar{n}\gamma^\mu n,$$
$$f_p = b_p = 2b_u + b_d, \quad f_n = b_n = 2b_d + b_u. \tag{13}$$

Notice that even in the case where these vector interactions are generated by other interactions, *e.g.*, colored mediators discussed later, the above description still holds. So the ISV $b_n = \mathcal{I} b_p$ corresponds to the coupling ratio at the quark level

$$\frac{b_u}{b_d} = \frac{2-\mathcal{I}}{2\mathcal{I}-1}, \tag{14}$$

for $\mathcal{I} = -0.7$ it takes a value $-9:8$. Obviously, similar to the case for the scalar interactions, the ISV from vector interactions also must come from the interactions between mediator/DM and the first-family quarks. And the DM-proton effective coupling is

$$f_p = \frac{3}{2\mathcal{I}-1} b_d, \tag{15}$$

and $f_p = -1.25 b_d$ for $\mathcal{I} = -0.7$.

As given in the introduction, we should pay special attention on the inelastic DM (iDM) [9]. In the $U(1)_X$ model, the vector currents involving DM couple to the vector gauge boson non-diagonal, which in the mass basis are

$$g_{ab}\bar{\chi}_a \gamma^\mu \chi_b Z'_\mu, \quad g_{ab}\phi_a^\dagger \overleftrightarrow{\partial}^\mu \phi_b Z'_\mu, \tag{16}$$

where $g_{ab} \sim \delta_{ab}$, and χ is pseudo Dirac while ϕ is approximate complex scalar. In other words, for the ISV iDM, an $U(1)_X$ model should receive special attention and we shall consider it in details later.

Comments are in orders: (i) The scalar interaction realization of proper ISV depends on both the isospin-violations in the nucleons itself, *i.e.*, $B_{u,d}^{(p)} \neq B_{u,d}^{(n)}$, and the DM-quark interactions, whereas the vector interactions only depends on the latter. (ii) The thermal DM with right relic density has annihilation rate $\langle \sigma_{an} v \rangle \simeq 1$ pb $\sim 10^2 \sigma_p$, given $\sigma_p \sim 10^{-2}$ pb. On the other hand, due to crossing symmetry, the annihilation rate and scattering rate from the operators accounting for the direct detection experiments scale as follows

$$\frac{\langle \sigma_{an} v \rangle}{\sigma_p} \sim \frac{M_{\rm DM}^2}{\mu_p^2}. \tag{17}$$

It just gives the ratio at right order for LDM around 10 GeV. Interesting, this numerical coincidence involving three basic elements of DM, the mass, relic density as well as scattering rate. In this regard, it "justifies" the GoGeNT/DAMA inspired LDM models with ISV. Of course, the rough argument ignores the velocity suppression for the operators like $\bar{\chi}\chi\bar{q}q$, but it is still sensible. Since in a complete model the pseudo-scalar type operators $(\bar{\chi}\gamma_5\chi)(\bar{q}\gamma_5 q)$ usually do exist, and has comparable coefficients. And we will see it via concrete discussions later.

B. Some Constraints

For a light DM with mass around 8 GeV and having quite large DM-nucleon scattering rate, it suffers from a list of model independent constraints (denoted as $\mathcal{C}_{1-4}$ in the following), coming from the cosmology, astrophysics and collider as follows

$\mathcal{C}_1$: **PAMELA** The PAMELA measures the anti-proton spectrum from 1-100 GeV, and it has no deviation from background [17]. Since the LDM couples to quark substantially, it probably renders the low energy spectrum of the anti-proton excess due to crossing symmetry [18, 19]. This constraint holds when the LDM has rather large annihilation rate (> 0.1 pb) into quarks today, thus invalid for the DM annihilation with velocity suppressing. Moreover, such constraints usually can be avoided by choosing proper astrophysical parameters [18].

$\mathcal{C}_2$: **Sun Neutrino** Because of the large DM-nucleon cross section, the Sun captures DM particles with large rate. These DM particles subsequently (cascade) annihilate into neutrinos, which are expected to be observed by the Super-Kamiokande. It excludes the thermal DM with predominantly annihilation modes into $\tau\bar{\tau}/\nu\bar{\nu}/4\tau$ as well as heavy quarks modes $\bar{b}b, \bar{c}c$ [20]. We have to emphasize that such neutrino constraint is so strong that generically one has to sufficiently suppress the DM annihilation via these channels, especially directly to neutrinos.

$\mathcal{C}_3$: **CMB** The lack of distortions in the Cosmic Microwave Background (CMB) spectrum due to DM annihilations at redshifts $z \sim 500 - 1000$ may give constraint on the GeV scale DM [21]. To make DM being a thermal relic, its annihilation modes should be dominated by the μ or τ modes (and corresponding 4 leptons), alternatively the DM annihilation is velocity dependent.

$\mathcal{C}_4$: **Colliders** DM-nucleus recoil can be converted to the DM production at the hadronic colliders, and the lack of relevant signals provides another constraint [22, 23]. In the ISV scenario the Tevatron gives very strong constraint. Vector interactions (as weak as $\sigma_p \sim 0.001$ pb) have been excluded definitely while the scalar interactions also close to the exclusion line, as given in the Fig. 1. It implies that the iDM with ISV must not be viable. To evade such Tevatron constraint from the effective operator analyses, we require that the mediator is light and its mass is much smaller than DM mass. In this way, the DM-nucleon cross section is enhanced by lighter mediator mass since $\sigma_n \propto 1/M^4_{mediator}$ whereas DM production at the collider scales as $1/s$, so the constraint made in Ref. [22, 23] is evaded. In addition, Ref. [24] gives the LEP bound on the operators $\mathcal{O}_{\rm DM}\bar{e}\Gamma e$, and shows that this operator is not allowed to provide main annihilation channel of DM freeze-out.

Let us summarize the constraints and point their possible implications (a recent more quantitative similar consideration see Ref. [25]). Among them, $\mathcal{C}_1$ can be satisfied thus not very serious. But $\mathcal{C}_2$ is a universal constraint independent on the DM spin, and from it we may expect that DM annihilates into μ or light quarks since the final state $e^{\pm}$ is excluded by $\mathcal{C}_4$. It is consistent with LDM in the CoGeNT/DAMA region where LDM likely annihilates into up and down quarks with cross section around 1pb, *e.g.*, in the models given in Section III B and Section III C. Then $\mathcal{C}_3$ may disfavor it. But take the astrophysical uncertainty into account, we just regard it as a referred constraint. Last but not the least, $\mathcal{C}_4$ gives the most powerful constraint in actual model building, it picks out the models whose DM-quark interaction either is a scalar type or via a light vector mediator exchange. In the case of light mediators as in Section III A, there is an elegant way to avoid various constraints by virtue of DM dominant invisible annihilation modes (to light states), and we shall discuss it in detail in the Section IV.

III. LIGHT DARK MATTER MODELS WITH ISV

In this Section, we systematically construct the ISV models. Let us explain our convention first. The $s-$channel Feymann diagrams have the vertices between DM-DM-mediators and between quark-quark-mediators, while the $t-$channel Feymann diagrams have the vertices betweem DM-quark-mediators. Thus, our $s-$channel and $t-$channel Feymann diagrams correspond to the $t-$channel and $s/u-$channel Feymann diagrams in the direct detection experiments, respectively. For simplicity, we shall not consider the DM models where the gauge bosons are DM candidates since they are generically complicated. Thus, the dark matter particles are either a scalar or a fermion. For $s-$channel Feymann diagrams, the mediators must be a scalar or a gauge boson since they can couple to two quarks. For $t-$channel Feymann diagrams, the mediators must be a color triplet fermion if DM is a scalar. And the mediator must be a color triplet scalar or vector boson if DM is a fermion. Because the models with color triplet vector boson are complicated in general, we will not study them in this paper. Therefore, we shall consider three kinds of LDM models with ISV, where the mediators are dominantly the $U(1)_X$ gauge boson Z', a spectator Higgs doublet, and color triplets respectively. For the first and second kinds of models, mediators propagate in the $s-$channel, thereby ISV is independent on the DM-mediator couplings, which simplifies the discussion. However, the color triplets propagate in $t-$channel, and their analyses are a little bit involved.

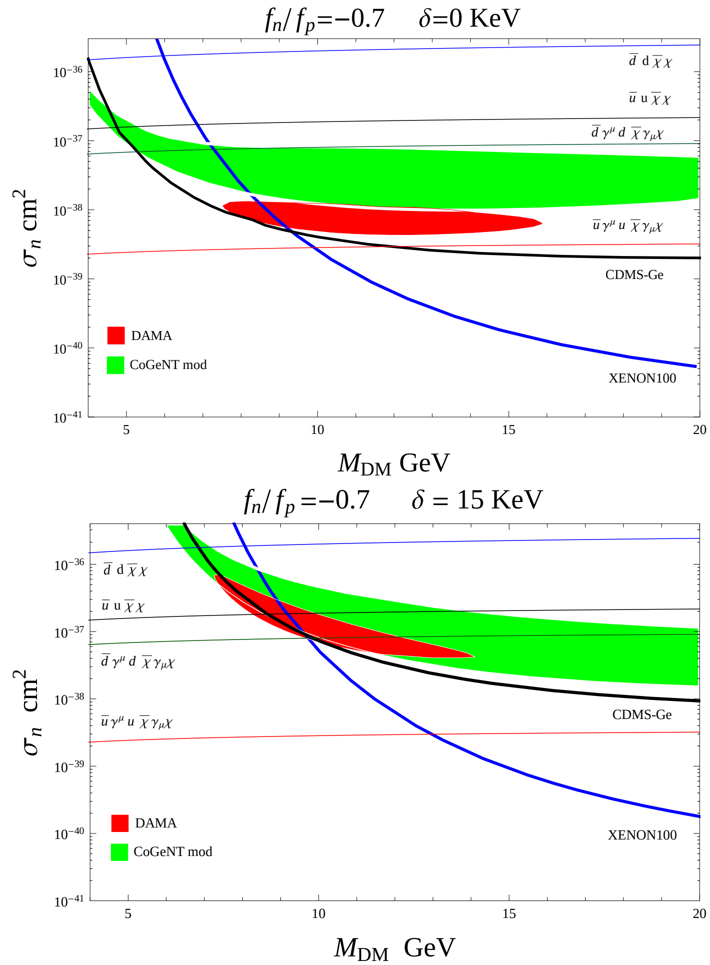

FIG. 1: Top: favored regions and exclusion contours in the $(M_{\rm DM}, \sigma_n)$ plane, ISV $f_n/f_p = -0.7, \delta = 0$; Bottom (for the annual modulation): ISV $f_n/f_p = -0.7, \delta = 15$ keV. Collider exclusion for some operators are also imposed. The data are from Ref. [10] and Ref. [23].

A. Chiral $U(1)_X$ Vector Boson

A simple way to produce ISV is to introduce an exotic $U(1)_X$ gauge boson as the mediator. From the discussions in the last Section, we only consider the light Z'. In practice, some of the following discussions that merely involve in the direct detection apply to any Z' boson with mass much larger than the transfer momentum. First, we briefly prove that the Z boson in the SM can only generate ISV with $|f_n/f_p| \ll 1$. Explicitly, the Lagrangian between the

Z boson and SM quarks

$$\mathcal{L}_{NC} \supset \frac{1}{\cos\theta_w}\left[\bar{u}_L\gamma^\mu\left(\frac{1}{2}-\frac{2}{3}\sin^2\theta_w\right)u_L+\bar{u}_R\gamma^\mu\left(-\frac{2}{3}\sin^2\theta_w\right)u_R\right.$$
$$\left.+\bar{d}_L\gamma^\mu\left(-\frac{1}{2}+\frac{1}{3}\sin^2\theta_w\right)d_L+\bar{d}_R\gamma^\mu\left(+\frac{1}{3}\sin^2\theta_w\right)d_R\right]Z_\mu\;. \tag{18}$$

Utilizing the formula Eq. (13), one can easily get

$$\frac{b_p}{b_n}=-\left(1-4\sin^2\theta_w\right)\approx -0.08 \ll 1. \tag{19}$$

We emphasize that this result depends only on the SM structure.

Now consider Z' from an exotic $U(1)_X$. We can gain some insight into the structure of the quark charges under $U(1)_X$, by investigating the ISV origin. And the discussions have nothing to do with concrete information in dark sector. We start from the Z' and SM fermion current couplings

$$\mathcal{L}_{NC}=-g_X\sum_{i,j}\bar{f}_i\gamma^\mu\left[(Q_{f_L})_{ij}P_L+(Q_{f_R})_{ij}P_R\right]f_jZ'_\mu, \tag{20}$$

where g_X is the gauge coupling of $U(1)_X$ and $Q_{f_{L/R}}$ are the charge matrices of the left- and right-handed fermions in family space. To not induce the tree-level flavor changing neutral current (FCNC), we reasonably assume they are diagonal matrices (for Abelian gauge group it must be ture). In the quark sector, transforming quarks into the mass eigenstates via $q_{L/R}\to V^\dagger_{q_{L/R}}q_{L/R}$ (the same letters are used to label gauge basis and mass basis), we have

$$\mathcal{L}'_{NC}=-\,g_X\sum_{q=u,d;i,j}\bar{q}_i\gamma^\mu\left[\left(V_{q_L}Q_{q_L}V^\dagger_{q_L}\right)_{ij}P_L+\left(V_{q_R}Q_{q_R}V^\dagger_{q_R}\right)_{ij}P_R\right]q_jZ'_\mu$$
$$\supset-\frac{g_X}{2}\sum_{q,i}\left[(Q_{q_L})_i|(V_{q_L})_{1i}|^2+(Q_{q_R})_i|(V_{q_R})_{1i}|^2\right]\bar{q}\gamma^\mu qZ'_\mu+... \tag{21}$$

In the second line, we explicitly show the first family, while dots collects axial vector current and other families.

In this paper, we consider that the charge matrices are family universe $(Q_{q_{L/R}})_i\equiv Q_{q_{L/R}}$, then the unitary of $V_{q_{L/R}}$ ensures that the SM fermion mixings will not induce flavor violation in Eq. (21). And the first-family quark charges give rise to proper ISV:

$$\frac{b_u}{b_d}=\frac{Q_{u_L}+Q_{u_R}}{Q_{d_L}+Q_{d_R}}=\frac{Q_{q_L}+Q_{u_R}}{Q_{q_L}+Q_{d_R}}. \tag{22}$$

In the second equation Q_{u_L}=Q_{d_L}= Q_{q_L} is taken since u_L and d_L are in the same $SU(2)_L$ multiplet. Clearly, the ISV effect is ascribed to the different right-handed up and down quark charges. So quarks must belong to the chiral representation under $U(1)_X$, and we are forced to work in the two Higgs doublet model (2HDM) where H_u couples to up-type quarks, while H_d only couples to the down-type quarks and charged leptons. By the way, it is easy to check that such $U(1)_X$ can not come from E_6 gauge symmetry if we want $\mathcal{I} \simeq -0.7$.

Through simple arguments, some useful conclusions are draw. The $U(1)_X$ gauge invariance of Yukawa interactions give

$$Q_{q_L} - Q_{u_R} + Q_{H_u} = 0, \quad Q_{q_L} - Q_{d_R} + Q_{H_d} = 0 \ . \tag{23}$$

If we do not introduce additional colored vector-like fermions, the $SU(3)_C^2 U(1)_X$ anomaly cancellation condition gives

$$3\left(Q_{q_L} - \frac{Q_{u_R}}{2} - \frac{Q_{d_R}}{2}\right) = 0 \ . \tag{24}$$

From these two equations, we get $Q_{H_u} = -Q_{H_d}$, and then we just need to introduce one Higgs doublet for the non-supersymmetric models. Combining them and Eq. (13), we get the charge conditions to obtain proper ISV

$$Q_{u_R} = \frac{7-5\mathcal{I}}{6(1-\mathcal{I})} Q_{H_u}, \quad Q_{d_R} = -\frac{5-7\mathcal{I}}{6(1-\mathcal{I})} Q_{H_u}, \quad Q_{q_L} = \frac{1+\mathcal{I}}{6(1-\mathcal{I})} Q_{H_u}. \tag{25}$$

And $\mathcal{I}$=-0.7 leads to a somewhat peculiar solution: $Q_{u_R} = 35/34 Q_{H_u}$, $Q_{d_R} = -33/34 Q_{H_u}$ and $Q_{q_L} = 1/34 Q_{H_u}$ [44]. However, if Eq. (24) is relaxed by introducing new colored vector-like fermions, the elegant solution may be found. But we just assume Eq. (24) is held in this paper.

Finally, the large DM-nucleon recoil cross section is easily achieved with light mediator. Assume the Dirac or complex scalar DM carries $U(1)_X$ charge $Q_{\rm DM}$, then in light of Eq. (15) and Eq. (25), we have

$$\begin{aligned} f_p &= \frac{g_X^2 Q_{\rm DM} Q_{H_u}}{M_{Z'}^2} \frac{2}{1-\mathcal{I}} \\ &\approx 1.2 \times 10^{-5} \times \left(\frac{1\,\text{GeV}}{M_{Z'}}\right)^2 \left(\frac{g_X^2 Q_{\rm DM} Q_{H_u}}{10^{-5}}\right), \end{aligned} \tag{26}$$

where $\mathcal{I} = -0.7$ is took. So typically we require a very weakly coupled (at least to quarks) $U(1)_X$.

Furthermore, the DM particles can be a scalar or a fermion. In particular, the $U(1)_X$ gauge symmetry can be broken down to a Z_2 discrete symmetry, which stabilizes the DM particles. Thus, we do not need to introduce extra discrete symmetry for DM particle to be stable.

B. (Approximate) Spectator Higgs Doublet

Higgs doublet is another typical mediator. But as argued in Section II A, if the Higgs doublet mediates DM-quark interactions meanwhile accounts for the SM fermion masses, the allowed ISV will be ignorable. So we have to introduce the Higgs doublet that is a (approximate) spectator to the electroweak symmetry breaking, namely it is only in charge of mediating the DM-SM fermion interactions. The origin of ISV arises from the Yukawa interactions involving in quark sector

$$-\mathcal{L}_Y \supset \left(y_{u,1}\bar{q}_L\epsilon H_1^\dagger u_R + y_{d,1}\bar{q}_L H_1 d_R\right) + \left(y_{u,2}\bar{q}_L\epsilon H_2^\dagger u_R + y_{d,2}\bar{q}_L H_2 d_R\right) + h.c. \ . \tag{27}$$

Here H_1 is the conventional SM Higgs fields whose neutral component acquires a VEV $\langle H_1^0\rangle \equiv v$ around 174 GeV. While $H_2 = \left(H^+, \frac{H^0+iA^0}{\sqrt{2}}\right)^T$, which carries the same SM quantum numbers as H_1, is the spectator Higgs with ignorable VEV. Thus, $y_{u/d,2}$ are free parameters, which we shall address this issue in details later. Quite interestingly, recently such a spectator Higgs doublet is also inspired to explain the CDF W+2jets anomaly [35, 36].

Now we shall consider the spectator Higgs and dark sector interactions. To keep the discussion as general as possible, we do not specify to any concrete model. The CP-even particle H^0 mediates DM-nucleon SI scattering, and as usual, H^0 probes the dark sector via the operators

$$a|\phi|^2 H^0, \quad \bar{\chi}(\alpha - \beta\gamma^5)\chi H^0, \tag{28}$$

$$a_A|\phi|^2 A^0, \quad \bar{\chi}(\alpha_A - \beta_A\gamma^5)\chi A^0, \tag{29}$$

where the CP-odd Higgs field A^0 does not generate SI cross section but opens another DM annihilation channels. In our scenario, we assume that H^0 barely mixes with the CP-even Higgs component of H_1, thus, such mixing factor can be ignored. The DM particles can be stable by introducing the Z_2 disccrete symmetry. If the DM particle is a real or complex scalar through Higgs port [38], we can realize the first term in Eq. (28) easily via the following renormalizable term

$$\lambda_\phi|\phi|^2 H_1 H_2 + h.c. \Rightarrow a = \sqrt{2}\lambda_\phi v. \tag{30}$$

If the dark matter is a fermion, to have the renormalizable interactions between the DM and H^0, one may consider a dark sector containing term $\lambda S\bar{\chi}\chi$, and the singlet scalar S further mixes with H^0 after the electroweak symmetry breaking. A more interesting possibility arises in the supersymmetric models, *e.g.*, in the next to the MSSM where the light neutralino $\widetilde{\chi}_1$ is the DM [39].

The ISV is determined by the flavor structure of H_2 Yukawa couplings with quarks. Transforming quarks into the mass eigenstates, we find that Eq. (27) turns to be

$$\mathcal{L}_Y \supset \frac{H^0 - iA^0}{\sqrt{2}} \bar{u} Y_{uu} P_R u - H^- \bar{d} Y_{du} P_R u + \frac{H^0 + iA^0}{\sqrt{2}} \bar{d} Y_{dd} P_R d + H^+ \bar{u} Y_{ud} P_R d + h.c., \tag{31}$$

where the effective Yukawa coupling matrices are defined by $Y_{qq'} \equiv V^\dagger_{q_L} y_{q',2} V_{q',R}$. Generally speaking, $Y-$matrices are not diagonal, consequently the spectator neutral Higgs mediates tree-level FCNC. But $y_{q,2}$ and V_{q_R} are free, thus, the flavor problem may be solved by choosing them properly. Integrating out H^0, the operator coefficients for the SI cross sections are

$$a_q = \frac{a}{\sqrt{2}(2M_{\rm DM})} \frac{1}{m^2_{H^0}} (Y_{qq})_{11}, \quad \frac{\alpha}{\sqrt{2} m^2_{H^0}} (Y_{qq})_{11}. \tag{32}$$

And in light of Eq. (10), the ISV $\mathcal{I} \approx -0.7$ is obtained provided that the Yukawa coupling ratio $(Y_{uu})_{11}/(Y_{dd})_{11} \approx -0.77$. Of course, we also require $(Y_{qq})_{22,33}$ sufficiently small so that the DM interact dominantly with the first-family quarks mediated by the scalar interaction.

Now we further assume the spectator Higgs is the source of the W+2jets anomaly via the process $p\bar{p} \to H^\pm \to W^\pm H^0/A^0 \to \ell^+\ell^-\nu + jj$ [35]. As a consequence the measured invariant mass of two jets fixes the charged Higgs mass $m_{H^0} \simeq 150$ GeV. For the benchmark scenario given in Ref. [35]: $(Y_{uu})_{11} \simeq 0.06$ and $m_{H^\pm} \simeq 250$ GeV, then to produce $\sigma_p \sim 0.01$ pb, we obtain

$$\begin{aligned} a\ (Y_{uu})_{11} &\sim 1.6\,\text{GeV}, \quad & a\ (Y_{dd})_{11} &\sim -2.0\,\text{GeV}, \\ \alpha\ (Y_{uu})_{11} &\sim 0.10, \quad & \alpha\ (Y_{dd})_{11} &\sim -0.13\,. \end{aligned} \tag{33}$$

For a real scalar or Majorana DM we only need half of these values. But H^0 will decay to a pair of DM, with branch decay width ratios over the one of $H^0 \to u\bar{u}$ as follows

$$\frac{\Gamma(H^0 \to 2\text{DM})}{\Gamma(H^0 \to u\bar{u})} \sim \frac{(Y_{uu})^{-4}_{11}}{3} \left(\frac{a (Y_{uu})_{11}}{M_{H_0}} \right)^2, \quad \frac{1}{3} \left(\frac{\alpha (Y_{uu})_{11}}{(Y_{uu})^2_{11}} \right)^2. \tag{34}$$

In the above estimation we have used the tree-level decay for H^0 decays to quarks, nevertheless the QCD corrections are important. For the (real) scalar DM the two channels are comparable, thus, the scenario is viable. Whereas for the fermionic DM H^0 predominantly invisible decays to DM. The situation can be improved by considering process $p\bar{p} \to H^0 \to W^\pm H^\mp \to \ell^+\ell^-\nu + jj$ [35], further inverting H^0 and $H^\pm$ mass hierarchy. Then $H^\pm \to W^\pm$+DM+DM decay width is suppressed by additional phase factor factor $1/(2\pi)^3$. Moderately increasing $(Y_{uu})_{11}$ and/or lowering m_{H^0} so as to increase the production rate of $H^\pm$, the fermionic DM may also be consistent with $W + jj$.

Finally we investigate the condition that such (quasi) spectator Higgs doublet appears. Equivalently speaking, when v_2 is small enough, the Yukawa couplings $y_{u/d,2}$ become free parameters. Roughly it requires $(Y_{uu})_{11}v_2 < m_u$, $(Y_{dd})_{11}v_2 < m_d$, where $(Y_{uu})_{11} \sim 0.06$ are taken to explain W+2jets. Thus, we have $v_2 < 0.05$ GeV. Now let us consider the renomalizable scalar potential of two Higgs doublet model

$$\begin{aligned} V =&\mu_1^2|H_1|^2 + \mu_2^2|H_2|^2 + \mu_{12}^2(H_1^\dagger H_2 + H_2^\dagger H_1) + \frac{\lambda_1}{2}|H_1|^4 + \frac{\lambda_2}{2}|H_2|^4 \\ &+ \lambda_3|H_1|^2|H_2|^2 + \lambda_4|H_1^\dagger H_2|^2 + \frac{\lambda_5}{2}(H_1^\dagger H_2)^2 + \frac{\lambda_5^*}{2}(H_2^\dagger H_1)^2, \end{aligned} \tag{35}$$

where all parameters are assumed to be real for simplicity. The small v_2 can be achieved by a small Higgs mixing term $\mu_{12}^2 \lesssim \mathcal{O}(10)$ GeV^2, which can be seen by the tadpole equations

$$\frac{\partial V}{\partial H_0} = \mu_{12}^2 v + \lambda_2 v_2^3 + \lambda_3 v^2 v_2 + \lambda_4 v^2 v_2 + \lambda_5 v^2 v_2 = 0, \tag{36}$$

from which it gives $v_2 \sim -\mu_{12}^2 v/((\lambda_3+\lambda_4+\lambda_5)v^2+\mu_2^2) \sim \mathcal{O}(10)$MeV with $\lambda_{3,4,5} \sim \mathcal{O}(0.1)$. In conclusion, it is not difficult to get an approximate spectator Higgs doublet to account for ISV as well the W+jj anomaly, given a scalar DM.

C. Color Triplets

The color triplet mediators are distinguished from the Z' and Higgs mediators since they mediate DM-quark interaction in the $t-$channel. In addition, they usually generate hybrid interactions, namely generating scalar and vector interactions simultaneously. It is more convenient to study such models from complete Lagrangian directly. In this paper, we concentrate on the fermionic DM particle and scalar color triplet mediators. The discussions for the scalar DM particle and fermionic color triplet mediators are similar, and will not be studied here. Moreover, to stabize the DM particles, we consider the discrete Z_2 symmetry. In particular, the DM particle and the color triplet mediators should be Z_2 odd.

We consider the most general mediators structure, each quark type $u_{L/R}$ and $d_{L/R}$ is accompanied by a color triplet mediator $\widetilde{q}_{L/R}$, just like the corresponding squarks in the MSSM. Since we are aiming at ISV, only the first family will be considered. Then the model preserving $SU(3)_C \times U(1)_{em}$ gauge symmetry takes a general form, adopting notation easily matches to the MSSM:

$$\begin{aligned} \mathcal{L} =& - m_\chi \bar{\chi}\chi - \sum_{\alpha,\beta=L/R} m_{u_{\alpha\beta}}^2 \widetilde{u}_\alpha \widetilde{u}_\beta^\dagger - \sum m_{d_{\alpha\beta}}^2 \widetilde{d}_\alpha \widetilde{d}_\beta^\dagger \\ & - \sum_{q=u,d}^{\alpha} \left[\lambda_{q_\alpha}\bar{\chi}(1+\gamma_5)q\widetilde{q}_\alpha^\dagger + \lambda'_{q_\alpha}\bar{\chi}(1-\gamma_5)q\widetilde{q}_\alpha^\dagger + h.c.\right]. \end{aligned} \tag{37}$$

Since $\widetilde{q}_L$ and $\widetilde{q}_R$ carry identical $SU(3)_C \times U(1)_{EM}$ quantum numbers, we allow triplets left-right (L-R) mixing as well as quark-mediators coupling with opposite chirality (L-R coupling), namely $\lambda_{q_L} \neq 0$ and $\lambda'_{q_R} \neq 0$. They are sources of chiral symmetry breaking. The triplet mass eigenstates are denoted by $\widetilde{q}_{1,2}$, related to gauge eigenstates by $\widetilde{q}_\alpha = \sum_l F_{q_\alpha l} \widetilde{q}_l$:

$$F_{q_L 1} = \cos\theta_q, \quad F_{q_L 2} = \sin\theta_q, \quad F_{q_R 1} = -\sin\theta_q, \quad F_{q_R 2} = \cos\theta_q. \tag{38}$$

$$\tan\theta_q = x_q - \sqrt{1 + x_q^2} < 0, \quad x_q \equiv \left(m^2_{\widetilde{q}_L} - m^2_{\widetilde{q}_R}\right)/2m^2_{\widetilde{q}_{LR}}. \tag{39}$$

The corresponding mass eigenvalues are given by $m_{\widetilde{q}_{1,2}}$. Then the interactions can be rewritten in a from

$$\begin{aligned} \mathcal{L} \supset & -\bar{\chi}\left(\alpha_l^q + \beta_l^q \gamma_5\right) q \widetilde{q}_l + h.c., \\ \alpha_l^q = & \sum_\alpha \left(\lambda_{q_\alpha} + \lambda'_{q_\alpha}\right) F_{q_\alpha l}, \quad \beta_l^q = \sum_\alpha \left(\lambda_{q_\alpha} - \lambda'_{q_\alpha}\right) F_{q_\alpha l}. \end{aligned} \tag{40}$$

In light of Eq. (A9) and Eq. (A12), integration out Φ leads to the effective operators involving SI cross section in the form of Eq. (2) and Eq. (3). And the operator coefficients

$$\begin{aligned} a_u = & -\frac{1}{m^2_{\widetilde{u}_l}} \mathrm{Re}\left(\lambda^*_{u_\alpha} \lambda'_{u_\beta} F^*_{u_\alpha l} F_{u_\beta l}\right), \\ b_u = & -\frac{1}{2m^2_{\widetilde{u}_l}} \left(\lambda_{u_\alpha} \lambda^*_{u_\beta} + \lambda'_{u_\alpha} \lambda'^*_{u_\beta}\right) F_{u_\alpha l} F^*_{u_\beta l}. \end{aligned} \tag{41}$$

Also, the expressions for a_d and b_d are obtained by replacing u with d in the above equations. There are two interesting limits. One is the chiral limit $a_{u/d} \to 0$, arises when both L-R mixing and L-R coupling are ignorable, then the scalar interactions vanish [45], leaving the pure vector interactions. In other words, interactions mediated by triplets in the chiral limit, the iDM scenario can be realized even in the absence of a vector gauge boson. This novel phenomena should receive special attention, but such models can not give negative $\mathcal{I}$ since b_u and b_d take the same sign. In contrast, in the Majorana limit where the DM is a Majorana fermion (or real scalar), the vector interactions vanish. In conclusion, colore triplets must generate ISV via the scalar interaction, and chiral limit must be avoided. In general it is expected that we both the scalar and vector interactions, as well as $f_n = a_n + b_n$.

L-R couplings usually are small, since the SM quarks and corresponding mediators are chiral and then are distinguished. In the MSSM, the lightest supersymmetric particle (LSP) is usual neutralino. The Higgsino component in the LSP neutralino has such coupling from superpotential, which nevertheless is tied to the chiral symmetry breaking in SM and then is tiny. Of course, one can relax this constraint by recurring to neutralino from spectator Higgs doublet, but the following discussion will not much different to the previous subsection, so we do not take it into account and only consider the L-R mixing as unique chiral symmetry breaking source. So we have $\lambda_{q_L} = \lambda'_{q_R} = 0$, and

$$a_q = -\frac{\lambda'_{q_L}\lambda_{q_R}}{m^2_{\widetilde{q}_1}} \times \left(\frac{m^2_{\widetilde{q}_2} - m^2_{\widetilde{q}_1}}{2m^2_{\widetilde{q}_2}} \sin 2\theta_q\right) \equiv -\frac{\lambda'_{q_L}\lambda_{q_R}}{m^2_{\widetilde{q}_1}} F_q, \tag{42}$$

with $q = u, d$. Clearly, if the two states are highly degenerate or L-R mixing is small, the effective coupling will be suppressed. So we consider the general case with $F_q \sim \mathcal{O}(1)$, as a result, according to Eq. (11) with $\mathcal{I} = -0.7$

$$f_p = 0.5 \times 10^{-5} \times \left(\frac{\lambda'_{d_L}\lambda_{d_R}}{1}\right)\left(\frac{500\,\mathrm{GeV}}{m_{\widetilde{q}_1}}\right)^2 \left(\frac{F_q}{1}\right) \mathrm{GeV}^{-2}, \tag{43}$$

for complex scalar double this estimated value.

An application of this result is the supersymmetric model with light $U(1)_X-$gaugino LSP, and the first-family squarks mediate DM-quark interactions (the second and third family contributions can be suppressed due to heavier masses or smaller $U(1)_X$ charges). Relevant terms are

$$\begin{aligned}\mathcal{L} \supset &-\frac{\sqrt{2}}{2} Q_{q_L} g_X \widetilde{B}_X (1-\gamma_5)\, q\widetilde{q}_L^\dagger \\ &-\frac{\sqrt{2}}{2} Q_{u_R} g_X \widetilde{B}_X (1+\gamma_5)\, u\widetilde{u}_R^\dagger - \frac{\sqrt{2}}{2} Q_{d_R} g_X \widetilde{B}_X (1+\gamma_5)\, d\widetilde{d}_R^\dagger + h.c.,\end{aligned} \tag{44}$$

where Q denotes $U(1)_X$ charge. The couplings defined in Eq. (37) are inherited from gauge couplings

$$\lambda'_{d_L} = \lambda'_{u_L} = \frac{\sqrt{2}}{2} g_X Q_{q_L}, \quad \lambda_{d_R} = \frac{\sqrt{2}}{2} g_X Q_{d_R}, \quad \lambda_{u_R} = \frac{\sqrt{2}}{2} g_X Q_{u_R}. \tag{45}$$

But actually a_q are only semi-quantized, by virtue of extra parameters $F_q/m^2_{\widetilde{q}_1}$. And now we have

$$\frac{a_u}{a_d} = \frac{Q_{u_R}}{Q_{d_R}} \left(\frac{m^2_{\widetilde{d}_1} F_u}{m^2_{\widetilde{u}_1} F_d}\right) \simeq -0.77. \tag{46}$$

From Eq. (39), it is seen that Q_{u_R}/Q_{d_R} <0 is still a necessary condition to get $\mathcal{I} < 0$, implying a chiral $U(1)_X$.

In practice, the MSSM with $U(1)_Y$ bino dominated LSP just fits this scenario. However, it fails to explain CoGeNT/DAMA quantitatively. In light of Eq. (43) and Eq. (45), even we have $F_d \sim 1$, to make squark masses lie above collider lower bound ($\gtrsim$500 GeV), we still need large quark gauge coupling $g_X Q_q \sim \mathcal{O}(1)$ (not excluded by collider for sufficiently heavy Z'). Whereas the (s)quark $U(1)_Y$ gauge couplings are too small. We have to stress that, exotic $U(1)_X$ is still interesting enough, since it realizes ISV in an economic way, asides from introducing $U(1)_X$ vector multiplet. Moreover, it does not trigger flavor problem.

Finally comment on another interesting aspect of the existence of color triplets at low energy based on SM. In the model given by Eq. (37), the Yukawa couplings $\lambda_{q/u/d}$ generically are complex c-number, so they introduce new physical CP phases. Recall that the SM fails to give sufficient baryon asymmetry, then these new color triplet mediators (If they are scalars, no anomalies will be introduced.) not only account for proper ISV, but also potentially provides adequate enough baryon asymmetry via heavy color triplet decays. Amazingly, for a Dirac dark matter, it is possible to realize asymmetric DM by the triplet common decay so as to resolve the coincidence puzzle $\Omega_{\rm DM}h^2 : \Omega_b h^2 \approx 5 : 1$. By the way, non-annihilating asymmetric Dirac DM automatically satisfies astrophysical constraints. We believe it is a minimal unified framework to realize ISV, asymmetric DM as well as baryogenesis, and we leave it for future work [37].

D. Dual Mediators

In the previous discussions we concentrate on ISV arises due to a single-type mediator, and no conventional mediators succeed in giving proper ISV. However, as proved in Section III A, the SM $Z-$boson actually only mediates DM-neutron interaction. So combining Z boson with Higgs field (without ISV effect) in principle can produce proper ISV. In fact, a complex sneutrino DM ν_1 [30, 31] just fits this scenario. Unfortunately, Z mediator plays no role in the ISV scenario. To satisfy the constraint from the $Z-$boson invisible decay (into a pair of DMs) width [33], we obtain the model independent constraint on the vertex $g_{Z_{11}} Z^\mu \widetilde{\nu}_1 \overleftrightarrow{\partial}_\mu \widetilde{\nu}_1^*$, whose effective coupling must be smaller than 0.023, *i.e.*, $g_{Z_{11}} < 0.023$. Then we have $f_n \approx -M_Z^{-2} g\, g_{Z_{11}}/4\cos\theta_w \sim -0.5 \times 10^{-6}\,{\rm GeV}^{-2}$, far to enough in the ISV scenario.

Another mediator frequently appears in the literatures is an exotic X_μ boson of $U(1)_X$ under which the SM fields are neutral, whereas $U(1)_X$ and $U(1)_Y$ have kinetic mixing term so that X_μ can mediate DM-quark interaction. As noticed in the Ref. [7], this kind of X_μ only mediates DM-proton interactions. Let us show this point explicitly. We consider the gauge kinetic sector of $U(1)_X \times U(1)_Y$

$$\mathcal{L}_{gauge} = -\frac{1}{4} F_Y^{\mu\nu} F_{Y\mu\nu} - \frac{1}{4} F_X^{\mu\nu} F_{X\mu\nu} + \frac{\theta}{2} F_Y^{\mu\nu} F_{X\mu\nu} \; . \tag{47}$$

We are interested in the small mixing limit $\theta \ll 1$, since it will not significantly affect the Z invisible decay width even with $U(1)_X$ charged LDM. Moreover, we consider the collider favored light X_μ scenario therefore $M_X < M_{\rm DM} \ll M_Z$. Then, after the electroweak symmetry breaking, working in the mass eigenstate basis (Z_μ, A_μ, X_μ), we obtain the interactions between the X_μ and the SM current at the leading order of θ [34]

$$\begin{aligned}\mathcal{L}_{coupling} \supset & \theta X_\mu \left(\cos\theta_w J^\mu_{em} + \mathcal{O}(M_X^2/M_Z^2) J_Z^\mu\right),\\ J^\mu_{em} = & g\sin\theta_w \left[\frac{2}{3}\bar{u}\gamma^\mu u + \left(-\frac{1}{3}\right)\bar{d}\gamma^\mu d + (-1)\,\bar{e}\gamma^\mu e\right].\end{aligned} \tag{48}$$

Apparently, the kinetic mixing just induces the coupling between the X_μ and electromagnetic current J^μ_{em}. So X_μ behaves like the photon that also only mediates DM-proton interaction.

We have shown that X can violate isospin. Next we consider the DM and light gauge boson gauge interactions, including the following DM-X_μ effective couplings

$$-\mathcal{L}_{\rm DM} \supset g_{\rm DM} X_\mu \phi^* \overleftrightarrow{\partial}^\mu \phi, \quad g_{\rm DM} X_\mu \bar{\chi}\gamma^\mu \chi. \tag{49}$$

So the DM-proton interaction effective operator coefficients are given by

$$b_p = \left(\frac{g\ \sin 2\theta_w}{2}\right)\left(\frac{g_{\rm DM}\theta}{M_X^2}\right). \tag{50}$$

The X_μ alone is not adequate to produce right $\mathcal{I}$, so we introduce the conventional Higgs as another mediator. The relevant discussion has been represented in Section III B, here we just identify H^0 as the SM Higgs h, then the DM-quark effective couplings are given by Eq. (32). In this case in the Eq. (7) $a_q/m_q \propto 1/v$ is family universal, and we get the isospin conserved contribution

$$\begin{aligned}a_p \approx a_n \approx & \frac{m_n}{\sqrt{2}(2M_{\rm DM})}\frac{a}{v}\frac{1}{m_h^2}\left(f^{(n)}_{T_s} + 3\times\frac{2}{27}f^{(n)}_{T_G}\right)\\ = & 0.15\times 10^{-5}\left(\frac{10\,\text{GeV}}{M_{\rm DM}}\right)\left(\frac{100\,\text{GeV}}{m_h}\right)^2 \text{GeV}^{-2}.\end{aligned} \tag{51}$$

We only include the s quark and heavy quarks contribution that gives a value $\simeq 0.31$ in the bracket. Moreover, we assume that the $\lambda_\phi |H|^2|\phi|^2$ term gives the coupling constant a as well as $\lambda_\phi = 1$. Perturbitivity and naturalness (large λ_ϕ contribute sub-TeV mass to the DM) do not favor large λ_ϕ. But the two Higgs DM model with large $\tan\beta$ enhancement can solve this numerical problem. For example, consider a simplified case where ϕ couples to Higgs fields mainly via a term $\lambda_\phi|\phi|^2 H_d H_u$, then in the Eq. (51) via H_d^0 mediation (we assume the heavier CP-even Higgs is dominated by H_d^0)

$$\frac{a}{v} \to \frac{\sqrt{2}\lambda_\phi v_u}{v_d} = \sqrt{2}\lambda_\phi \tan\beta. \tag{52}$$

and now it is required that $\lambda_\phi \tan\beta/(m_{H_d^0}/100\,\text{GeV})^2 \sim 1$. In conclusion, the total DM-nucleon coupling $f_p = a_p + b_p$, $f_n = a_n \approx a_p$, and the proper ISV is produced as follows

$$\frac{b_p}{a_p} = \frac{\mathcal{I}}{1-\mathcal{I}}, \tag{53}$$

which is readily to be realized.

Analysis on fermionic DM can be employed similarly, with the replacement $a/2M_{\rm DM} \to \alpha$. In general, the numerical problem is exacerbated, and maybe somewhat light m_h (even with large $\tan\beta$) is required to enhance the cross section, just as required in the NMSSM with neutralino LSP [39].

IV. SNEUTRINO IDM WITH LIGHT Z' MEDIATOR

Interestingly, the light iDM models with ISV can explain the CoGeNT and DAMA annual modulation consistently, as well as reconcile them with all other results [10]. In particular, we can evade the other astrophysical constraints naturally in quite a few models, for example, the spectator Higgs mediator models. To be concrete, we shall propose sneutrino iDM model with a very light and very weakly coupled Z' mediator.

A. Motivation for iDM

In the first place, in light of the general analyses made in Section II B, the Z' should be lighter than the DM. We even require $M_{Z'} < 1$ GeV to escape from various astrophysical exclusion elegantly. It is much lighter than the consideration like in Ref. [27], nevertheless natural since the Z' mass scale closes to the LDM mass scale [46]. The light Z' mediator allows very weakly coupling $g_X \ll 1$, in this way the Tevatron [22, 23] and LEPII [40] constraints are satisfied. Moreover, kinetic block mechanism makes the DM main annihilation products to be soft $2e$ or/and 2μ today. Then the constraints from Sun neutrino, CMB as well as PAMELA can be readily satisfied.

We explicitly show how this mechanism works for a scalar iDM. At the early Universe, the DM has two kinds of comparable annihilation channels: DM+DM$\to q\bar{q}$ via Z' mediation, as well as the invisible annihilation modes DM+DM$\to XX$, with X the sub-GeV hidden states such as Z' (we take it as an example). Then the total annihilation rate is

$$\sigma_{an}|v| \sim \sum_{\pm} \frac{\frac{1}{4}g_X^4 Q_{\rm DM}^2 \left(Q_{f_L} \pm Q_{f_R}\right)^2}{16\pi}\frac{c_f}{M_{\rm DM}^2}v^2 + \left(\frac{g_X^4 Q_{\rm DM}^4}{16\pi}\frac{1}{M_{\rm DM}^2} + ...\right), \tag{54}$$

where Q_f is understood to be effective charge in the mass eigenstates. Terms in the bracket collect DM annihilation into hidden states. Provided that the $U(1)_X$ charges for the SM fermions and DM particle are comparable, today the second term overwhelmingly dominates (since $v \sim 10^{-3}$), subsequently the Z' cascade decays into $e/\mu/\gamma$ have no dangerous modes, thus the strong tension with astrophysics is resolved.

In contrast, the Dirac iDM annihilation modes DM+DM→ $f\bar{f}$ have no v^2 suppressing, so it is quite unnatural to suppress dangerous modes to desired level. In fact, direct annihilations into the SM fermions (assuming family universal Z') are almost fixed by cross symmetry. According to Eq. (26), to get the adequately large DM-proton (inelastic) scattering rate $\sigma_p \sim \mathcal{O}(0.1)$ pb, we need $g_X^2 Q_{H_u} Q_{\rm DM} \sim 10^{-4}$. On the other hand, from Eq. (54) we know that the direct DM annihilations into the SM fermions have a rate about $\mathcal{O}(10^{-2})$ pb. While the Dirac DM annihilations via $s-$wave have even larger rate $\mathcal{O}(10^{-1})$ pb. This result is derived from cross symmetry, thus, the Dirac DM annihilation rate can be reduced by no way, except by lowering Z' mass even orders lighter. Or one can consider only the u and d quark charges under $U(1)_X$ to avoid Sun neutrino constraint, but it leads to a weird model. In conclusion, scalar iDM with sub-GeV Z' mediator moreover dominantly decays to sub-GeV light hidden states is a preferred model for the GoGeNT annual modulation.

B. Light Sneutrino iDM from Low Scale Seesaw Mechanism

In the MSSM extended with low-scale seesaw mechanism, (light) sneutrino is a natural scalar iDM candidate (We also present a quite simple non-supersymmetric model in Appendix C.). To realize ISV we further consider a light $U(1)_X-$extension, and the right-handed neutrino (RHN) is charged under $U(1)_X$. Then it is natural to attribute the origin of Majorana mass scale to the $U(1)_X$ breaking scale. The relevant parts in the minimal model are

$$W \supset y_{ij}^N L_i N_j H_u + \frac{\lambda_i}{2} S N_i^2 + \mu H_u H_d, \tag{55}$$

$$-\mathcal{L}_{soft} \supset \left(m_{\widetilde{N}_i} |\widetilde{N}_i|^2 + m_S^2 |S|^2 \right) + A_0 \left(y_{ij}^N \widetilde{L}_i \widetilde{N}_j H_u + \frac{\lambda_i}{2} S \widetilde{N}_i^2 + h.c. \right), \tag{56}$$

For simplicity, all parameters are assumed to be real, and soft mass square terms for the MSSM singlets S and N_i (we simply consider one family only) are around GeV scale or even smaller. To reduce the parameters, we assume a common trilinear term A_0. Because it controls the mass splitting of $\widetilde{N}_i$, we further assume it to be very small on our purpose. Such soft parameter pattern is natural in the gauge mediated supersymmetry breaking (GMSB) scenario. S carries $U(1)_X$ charge $Q_S = -Q_N/2$, and it develops a VEV with $\langle S \rangle \equiv v_s \sim \mathcal{O}(100)$ GeV. Thus, it breaks the exotic $U(1)_X$ gauge symmetry at low scale and gives the Majorana mass terms for the RHN $M_N = \lambda v_s$. Since we want to naturally have a sneutrino state with mass around 8 GeV to be the LSP, we expect $M_N \sim \mathcal{O}(10)$ GeV, and then the neutrino Yukawa couplings $y^N \sim \mathcal{O}(10^{-7})$ are irrelevant to our discussion.

This model is able to realize the above scenario readily. First, let us consider the $U(1)_X$

spontaneously breaking. The scalar potential of the S is given by

$$\begin{aligned} V_S =& V_F + V_D + V_{soft}, \\ V_F =& |\lambda_i S N_i + y^N_{ji} L_j H_u|^2, \\ V_D =& \frac{g_X^2}{2}\left(Q_{H_u}|H_u|^2 + Q_{H_d}|H_d|^2 + Q_S|S|^2 + \sum_f Q_f|\widetilde{f}|^2\right)^2, \end{aligned} \tag{57}$$

and $V_{soft} = -\mathcal{L}_{soft}$ is given by Eq. (56). So the crucial part of the scalar potential of S can be casted in a form of simple ϕ^4 field theory

$$V_S \supset -\mu_S^2|S|^2 + \frac{\kappa}{4}|S|^4, \tag{58}$$

where

$$\mu_S^2 = \left(g_X^2 Q_{H_u} Q_S \cos 2\beta\right) v^2 - m_S^2, \quad \kappa = 2g_X^2 Q_S^2, \tag{59}$$

where $Q_{H_d} = -Q_{H_u}$ is used. Interestingly, a simple way to trigger $U(1)_X$ breaking is available even with positive m_S^2: the first term in the μ_S^2 is negative provided

$$Q_{H_u} Q_S < 0, \tag{60}$$

then μ_S^2 in principle can be as (positive) small as will. It leads to

$$v_s = \sqrt{\frac{2}{\kappa}}\mu_S \sim \sqrt{\cos 2\beta |Q_{H_u}/Q_S|} \times v, \tag{61}$$

at the weak scale. In the above estimation we have set $m_S^2 \to 0$. In other words, in our model $D-$term is adequate to break the $U(1)_X$ gauge symmetry without resorting to any exotic requirements.

Next we examine the mass spectrum. The Z' mass is determined by v_s. We have to stress that without S's VEV one linear combination of Z'_μ and Z_μ is still massless after the electroweak symmetry breaking. The point is that $Q_{H_u} = -Q_{H_d}$ and then the Higgs mixing term $B\mu H_u H_d$ is allowed. Now, ignoring the small $Z - Z'$ mixing effect, at the leading order we have

$$M_{Z'} \approx \sqrt{2} g_X |Q_S| v_s \simeq \sqrt{2|Q_{H_u}||Q_S|\cos 2\beta} \times g_X v \ . \tag{62}$$

Thus, $M_{Z'}$ can naturally at the GeV order by choosing $g_X \sim 10^{-2}$. On the other hand, at the leading order, the CP-even Higgs field from the singlet S (denoted as h_s) has degenerate mass with the Z'. However, in this minimal model, LSP tends to be singlino $\widetilde{S}$, which has a light (degenerate with $M_{Z'}$) Dirac mass term with the $U(1)_X$ gaugino $\widetilde{\lambda}$. So the heavier the Majorana gaugino mass term of $\widetilde{\lambda}$, the lighter the Majorana mass fermion $\widetilde{S}$. To solve this problem, we introduce extra singlets S' (they are maybe required to cancel the $U(1)_X$ gauge anomalies), with terms $M_i S S_i$ or $S^2 S_i$. As a result no light (lighter than $\widetilde{N}$) $R-$parity odd particles are left in the particle spectrum.

In this model, DM has many annihilation channels, and we simply assume the dominant mode is DM+DM$\rightarrow$ $2h_s$ via the contact interactions from $|\lambda|^2|\widetilde{N}||S|^2$. Right DM relic density requires $\lambda \sim 0.1$, which is consistent with $v_s \sim \mathcal{O}$ (100) GeV and $M_N = \lambda v_s \sim 10$ GeV.

Finally, the sneutrino LSP is dominated by the RHN sparticle $\widetilde{N}$. Taking into account that the mass splitting between CP-even and CP-odd states due to non-zero A_0, we get an iDM with mass and splitting

$$M_{\rm DM}^2 = m_{\widetilde{N}_1}^2 = \lambda^2 v_s^2 + m_{\widetilde{N}}^2, \quad \delta \simeq \frac{\lambda v_s}{m_{\widetilde{N}_1}} A_0. \tag{63}$$

So to get $\delta \sim 10^{-5}$ GeV, A_0 is expected to be around this scale. It is maybe upset by virtue of fine-tuning, but a version equipped with inverse seesaw mechanism [29] can resolve this tuning problem [30]. This model contains an extra singlet $\bar{N}$,

$$W \supset y_N L H_u N + m_N N \bar{N} + M_{\bar{N}} \bar{N}^2/2. \tag{64}$$

Then the sneutrino mass splitting can be related to the origin of small neutrino mass rather than merely the soft mass terms [30]:

$$\delta \sim m_\nu \frac{m_N}{m_D} + m_\nu \frac{m_N}{m_D} \frac{B}{m_D}, \tag{65}$$

where we have assumed the Dirac mass term $m_D = y_N v_u$ is much lighter than the second Dirac mass term m_N: $m_D \ll m_N \simeq m_{\widetilde{N}_1}$. And B is the bilinear soft mass for $\widetilde{\bar{N}}$. Even set $B = 0$, in principle $m_N/m_D \sim 10^6$ just give the right order for δ.

V. DISCUSSION AND CONCLUSION

Light dark matter models with the ISV $f_n/f_p \approx -0.70$ and large DM-nucleon spin-independent cross section $\sigma_n \sim \mathcal{O}(0.01)$ pb may provide a way to understand the confusing direct detection experimental results. Combining with stringent astrophysical and collider constraints, we can further deduce the DM properties. In this work, we investigated the possible origin of ISV. General anatomy on the effective operators generating SI cross section is made, and we found that ISV essentially arises from the DM and first-family quarks couplings. To further explore their UV origin, we propose three kinds of models with the following mediator structure:

Z' from $U(1)_X$ The $U(1)_X$ must be chiral, and the light Z' is strongly favored.

Spectator Higgs Conventional Higgs doublet mediates interactions preserving isospin. So we have to introduce (approximate) spectator Higgs doublet whose couplings to the SM quarks are free parameters. Such a Higgs doublet can be used to explain Tevatron CDF $W + jj$ anomaly.

Color triplets Combining the squarks in the MSSM with exotic $U(1)_X$ (quarks strongly charged under it), we found that the light $\widetilde{B}_X$ LSP can generate proper ISV via the first-family squark mediation. This scenario is economic furthermore suffers no flavor problem.

Exotic Z' plus Higgs For a SM-neutral $U(1)_X$ having kinetic mixing with $U(1)_Y$, its light gauge boson X_μ only mediates DM-proton interaction. Combining it with a conventional Higgs mediator, we can obtain the right ISV.

Inspired by the CoGeNT annual modulation, we propose a light sneutrino as the iDM. It is charged under a $U(1)_X$, and the light Z' ($M_{Z'}$ <1 GeV) leads to a proper ISV. Moreover, this model naturally reconciles the ISV scenario with the stringent constraints from the astrophysical (such as Sun neutrino) and collider (Tevertron mono-jet search) experiments.

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Appendix A: Operators after Integrating Out Mediators

The models on the DM-SM fermions interactions can be classified based on the propagators mediating DM and SM particle interactions. In this appendix, we borrow some results from the Ref. [13]. For the scalar DM, it interacts with SM fermions by exchanging a Z' boson, a (real) Higgs doublet h and the colored fermion Q, the corresponding Lagrangian is given by

$$\mathcal{L} = -\frac{1}{4}F'^{\mu\nu}F'_{\mu\nu} + \frac{1}{2}m_{Z'}^2\, Z'^{\mu}Z'_{\mu} + a\phi^{\dagger}\overleftrightarrow{\partial}_{\mu}\phi Z'^{\mu} + \bar{q}\gamma^{\mu}(\alpha - \beta\gamma^5)q\; Z'_{\mu}, \tag{A1}$$

$$\mathcal{L} = \frac{1}{2}(\partial h)^2 - \frac{1}{2}m_h^2 h^2 - a\phi^{\dagger}\phi h - \bar{q}(\alpha - \beta\gamma^5)qh, \tag{A2}$$

$$\mathcal{L} = \bar{Q}\left(i\partial - m_Q\right)Q - \bar{q}\left(\alpha - \beta\gamma^5\right)Q\phi^{\dagger} - h.c.. \tag{A3}$$

Integrating out the heavy propagators via equation of motion, we obtain the effective operators generating SI cross section (other operators belongs to Eq. (4), we do not list here)

$$\mathcal{L}_{eff} \supset -\frac{a\alpha}{m_{Z'}^2}(\phi^{\dagger}\overleftrightarrow{\partial}_{\mu}\phi), \tag{A4}$$

$$\mathcal{L}_{eff} \supset \frac{a\alpha}{m_h^2}\phi^{\dagger}\phi\;\bar{q}\,q,\;\;, \tag{A5}$$

$$\mathcal{L}_{eff} \supset \frac{1}{m_Q}\left(|\alpha|^2 - |\beta|^2\right)\bar{q}q\phi^{\dagger}\phi + \frac{i}{m_Q^2}\left(|\alpha|^2 + |\beta|^2\right)\bar{q}\gamma^{\mu}q\phi^{\dagger}\overleftrightarrow{\partial}_{\mu}\phi. \tag{A6}$$

For the real scalar DM, vector interactions disappear. Interactions of fermionic DM can be described analogous to scalar DM

$$\mathcal{L} = -\frac{1}{4}\mathcal{F}'^{\mu\nu}\mathcal{F}'_{\mu\nu} + \frac{1}{2}m_{Z'}^2\, Z'^{\mu}Z'_{\mu} + \bar{\chi}\gamma^{\mu}(\alpha - \beta\gamma^5)\chi Z'_{\mu} + \bar{q}\gamma^{\mu}(\widetilde{\alpha} - \widetilde{\beta}\gamma^5)qZ'_{\mu}, \tag{A7}$$

$$\mathcal{L} = \frac{1}{2}(\partial h)^2 - \frac{1}{2}m_h^2 h^2 - \bar{\chi}(\alpha - \beta\gamma^5)\chi h - \bar{q}(\widetilde{\alpha} - \widetilde{\beta}\gamma^5)qh, \tag{A8}$$

$$\mathcal{L} = |\partial\Phi|^2 - m_{\Phi}^2|\Phi|^2 - \bar{\chi}(\alpha - \beta\gamma^5)q\Phi - h.c.\;, \tag{A9}$$

where Φ denotes the scalar color triplet mediators. And the corresponding effective operators are

$$\mathcal{L}_{eff} \supset -\frac{1}{m_{Z'}^2}\alpha\widetilde{\alpha}\;\bar{\chi}\gamma^{\mu}\chi\;\bar{q}\gamma_{\mu}q, \tag{A10}$$

$$\mathcal{L}_{eff} \supset \frac{\alpha\widetilde{\alpha}}{m_h^2}\bar{\chi}\chi\;\bar{q}q, \tag{A11}$$

$$\mathcal{L}_{eff} \supset \frac{1}{4m_{\Phi}^2}\left[\left(|\alpha|^2 - |\beta|^2\right)\bar{\chi}\chi\bar{q}q + \left(|\alpha|^2 + |\beta|^2\right)\bar{\chi}\gamma_{\mu}\chi\bar{q}\gamma^{\mu}q\right]\;. \tag{A12}$$

When the DM is a Majorana fermion or a real scalar, the vector interaction vanishes.

Appendix B: Scattering and Annihilating

In this appendix we briefly introduce the formula involving direct detections and give the relevant annihilation rates. Ignoring small ISV from the form factor of proton and neutron, the DM-nucleus SI scattering cross section at zero momentum transfer (not the actual cross section) can be written in a form [41]

$$\sigma_0 = \frac{\delta_C\mu_N^2}{\pi}\left[Zf_p + (A - Z)f_n\right]^2, \tag{B1}$$

where A is the atomic mass of the nucleus while Z is its atomic number. The reduced mass $\mu_N = M_{\rm DM} m_N/(M_{\rm DM} + m_N)$, and $\delta_C = 4$ for the self-conjugate particle like Majorana and real scalar DM, otherwise $\delta_C = 1$. In the reference DM-proton scattering cross section is used frequently

$$\sigma_p = \frac{\delta_C \mu_p^2}{\pi} f_p^2, \tag{B2}$$

where μ_p the DM-proton reduced mass. And f_p from scalar interaction in Eq. (2) and vector interaction in Eq. (3) are respectively given by

$$\begin{aligned} \text{Fermionic DM}: \quad f_p &= \sum_q B_q^p a_q; \qquad\qquad f_p = 2a_u + a_d, \\ \text{Scalar DM}: \quad f_p &= \sum_q B_q^p \frac{a_q}{2M_{\rm DM}}; \quad f_p = 2a_u + a_d. \end{aligned} \tag{B3}$$

The DM can annihilate into the SM particles, which determines the final relic density after the annihilation freeze-out. And the thermal average annihilation cross section, expanded with relative velocity v_{rel} (subscript will be omitted) takes a form of

$$\langle \sigma v \rangle_{F.O.} = a + b\langle v^2 \rangle = (a + 3b/x_f) \tag{B4}$$

where $x_f \equiv M_{\rm DM}/T_f = 3/\langle v^2 \rangle$ with T_f the DM decoupling temperature, for the weakly interactive massive particle (WIMP) typically having $x_f \sim 20 - 30$. The DM relic density can be formulated as

$$\Omega h^2 \approx \frac{1.07 \times 10^9 x_f {\rm GeV}^{-1}}{M_{\rm Pl}\sqrt{g_*}\langle \sigma v \rangle_{F.O.}}, \tag{B5}$$

with g_* as the effective relativistic degree of freedom when the DM decouples. The actual effective annihilation rate, which is used to determine relic density and calculate signals from DM annihilation today, is given by

$$\langle \sigma v \rangle_{F.O} = T_{\rm DM} \times (a + 3b/x_f) \ , \tag{B6}$$

wher $T_{\rm DM} = 1/2$ for the complex DM while $T_{\rm DM} = 4$ for the self-conjugate DM [42].

The a and b can be extracted out from partial wave expansion of the cross section times the relative velocity $\sigma v = a + bv^2$ (regarded as a rough thermal averaged cross section). Here the relative velocity is $v = 2\sqrt{1 - 4M_{\rm DM}^2/s}$ with s the Mandelstam variable. For the fermionic DM, σv from operators involved given by [43],

$$a_f \bar{\chi}\chi \bar{f} f: \qquad \frac{c_f}{16\pi} \times 2a_f^2 M_{\rm DM}^2 \beta_f^3\, v^2, \tag{B7}$$

$$\frac{G_{P,f}}{\sqrt{2}} \bar{\chi}\gamma^5\chi \bar{f}\gamma^5 f: \qquad \frac{c_f}{4\pi} \times G_{P,f}^2 M_{\rm DM}^2 \beta_f, \tag{B8}$$

$$\frac{G_{PS,f}}{\sqrt{2}} \bar{\chi}\gamma^5\chi \bar{f} f: \qquad \frac{c_f}{4\pi} \times G_{PS,f}^2 M_{\rm DM}^2 \beta_f^3, \tag{B9}$$

$$\frac{G_{SP,f}}{\sqrt{2}} \bar{\chi}\chi \bar{f}\gamma^5 f: \qquad \frac{c_f}{16\pi} \times G_{SP,f}^2 M_{\rm DM}^2 \beta_f, \tag{B10}$$

$$b_f \bar{\chi}\gamma_\mu\chi \bar{f}\gamma^\mu f: \qquad \frac{c_f}{4\pi} \times 2b_f^2 M_{\rm DM}^2 \beta_f \left(2 + z_f\right), \tag{B11}$$

$$\frac{G_A}{\sqrt{2}} \bar{\chi}\gamma^5\gamma_\mu\chi \bar{f}\gamma^5\gamma^\mu f: \qquad \frac{c_f}{4\pi} \times G_{A,f}^2 M_{\rm DM}^2 \beta_f \left[z_f + \frac{1}{12}\left(4 - z_f\right) v^2\right], \tag{B12}$$

$$\frac{G_{AV,f}}{\sqrt{2}} \bar{\chi}\gamma^5\gamma_\mu\chi \bar{f}\gamma^\mu f: \qquad \frac{c_f}{48\pi} \times G_{AV,f}^2 M_{\rm DM}^2 \beta_f^2\, v^2, \tag{B13}$$

$$\frac{G_{VA,f}}{\sqrt{2}} \bar{\chi}\gamma_\mu\chi \bar{f}\gamma^5\gamma^\mu f: \qquad \frac{c_f}{2\pi} \times G_{VA,f}^2 M_{\rm DM}^2 \beta_f, \tag{B14}$$

$$\frac{G_T}{\sqrt{2}} \bar{\chi}\sigma^{\mu\nu}\chi \bar{f}\sigma_{\mu\nu} f: \qquad \frac{c_f}{4\pi} \times G_{T,f}^2 M_{\rm DM}^2 \beta_f \left(7 + z_f\right), \tag{B15}$$

where the final state velocity $\beta_f \equiv \sqrt{1 - z_f}$ and $z_f \equiv m_f^2/M_{\rm DM}^2$, the color factor $c_f = 3$ for quarks otherwise 1. The scalar DM and relevant operators σv are given by

$$a_f |\phi|^2 \bar{f} f: \qquad \frac{c_f}{8\pi} \times 2a_f^2 \beta_f^3, \tag{B16}$$

$$\frac{F_{Vf}}{\sqrt{2}} \phi^\dagger \overleftrightarrow{\partial}_\mu \phi \bar{f}\gamma^\mu f: \qquad \frac{c_f}{4\pi} \times F_{Vf}^2 M_{\rm DM}^2 \beta_f \left[\frac{2}{3}\left(2 + z_f\right) v^2\right], \tag{B17}$$

$$\frac{F_{SPf}}{\sqrt{2}} |\phi|^2 \bar{f}\gamma_5 f: \qquad \frac{c_f}{8\pi} \times F_{SPf}^2 \beta_f, \tag{B18}$$

$$\frac{F_{VAf}}{\sqrt{2}} \phi^\dagger \overleftrightarrow{\partial}_\mu \phi \bar{f}\gamma^\mu\gamma^5 f: \qquad \frac{c_f}{4\pi} \times F_{VAf}^2 M_{\rm DM}^2 \beta_f \left[\frac{2}{3}\left(2 - z_f\right) v^2\right]. \tag{B19}$$

Appendix C: A Non-Supersymmetric Scalar iDM Model

In the non-supersymmetric case, there is a simple way to realize iDM at renormalizable level. The dark sector consists of two SM singlets $\phi_{1,2}$ with mass hierarchy $\mathcal{O}(8\,\text{GeV})^2 \sim m_{\phi_1}^2 \ll m_{\phi_2}^2$. The relevant scalar potential for iDM generation is quite simple

$$-V \supset \left(m_{\phi_1}^2|\phi_1|^2 + m_{\phi_2}^2|\phi_2|^2\right) + \left(\eta_1\phi_1\phi_2^* S^2 + \eta_2\phi_2^2 S^2 + h.c\right). \tag{C1}$$

We need to arrange charge assignment properly so that the dark sector conserves the Z_2 symmetry to protect DM stable. Singlet S breaks $U(1)_X$ at $v_s \sim \mathcal{O}(100)$ GeV as the model given in the text, which also induces a quartic mass term for the heavy state ϕ_2

$$-V \supset \eta_2 v_s^2 \phi_2^2 + c.c., \tag{C2}$$

and breaks the dark global $U(1)$ symmetry (acting on ϕ_i only). It renders small mass splitting between the CP-even and CP-odd states of ϕ_1. To extract out the mass splitting analytically, we diagonlize the mass matrix for $\phi_{1,2}$ by the unitary matrix U_ϕ, where we ignore the $U(1)$ breaking mass term. Then we obtain

$$\phi_1 \to \cos\theta_{12}\phi_1' + \sin\theta_{12}\phi_2', \quad \phi_2 \to -\sin\theta_{12}\phi_1' + \cos\theta_{12}\phi_2', \tag{C3}$$

where the fields with prime are in the (approximate) mass eigenstates. And the mixing angle

$$\theta_{12} \simeq \eta_1 \frac{v_s^2}{m_{\phi_2}^2} \ll 1 \tag{C4}$$

is invalid if $(\eta_1 v_s^2)^2 < m_1^2 m_2^2$ (assure the positivity) and $|\eta v_1 v_2|, m_1^2 \ll m_2^2$. With it, substituting Eq. (C3) into Eq. (C1), then it is not difficult to get the mass splitting for the ϕ_1

$$\delta = 2\eta_2 \sin\theta_{12}^2 \times \frac{v_s^2}{m_{\phi_1}} \approx 2\eta_2\eta_1^2 \left(\frac{v_s}{m_{\phi_2}}\right)^4 \frac{v_s^2}{m_{\phi_1}}. \tag{C5}$$

Thus, we can obtain $\delta \sim 10^{-5}$ GeV in many ways, even by setting the m_2 at the interesting TeV scale.

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[45] In the case with real scalar DM and fermionic triplet mediators, the mediate Dirac masses, *i.e.*, the L-R mixings provide the chiral symmetry breaking sources. For an example please see Ref. [5].
[46] Another interesting application of this sub-GeV scale Z' scenario is the kinetically suppressing antiproton product, when we consider cosmic ray anomaly from heavy DM decaying into Z' that cascade decays to charged leptons [7, 28].