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Cancellation in dark matter-nucleon interactions: The role of non-standard-model-like Yukawa couplings

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ABSTRACT

Extensive searches to probe the particle nature of dark matter (DM) have been going on for some decades now but, so far, no conclusive evidence has been found. Among various options, the Weakly Interacting Massive Particles (WIMP) remains one of the prime possibilities as candidates for DM near the TeV scale. Taking a phenomenological view, such null results may be explained for a generic WIMP in a Higgs-portal scenario if we allow the light-quark Yukawa couplings to assume non-Standard Model (non-SM)-like values. This follows from a cancellation among different terms in the DM-nucleon scattering which can, in turn, lead to a vanishingly small direct-detection cross section. It might also lead to isospin violation in the DM-nucleon scattering. Such non-SM values of light-quark Yukawa couplings may be probed in the high luminosity run of the LHC.

1. Introduction

In the search for new physics (NP), a large class of current and future experiments are dedicated to test the particle nature of dark matter (DM), especially by probing the ones known as the Weakly Interacting Massive Particles (WIMPs) (see e.g., [1,2]). Among them, the direct detection (DD) experiments rely upon the scattering of a DM particle off a detector nucleus, thereby looking for the recoiled nucleus. The list of current experiments in this direction includes Xenon1T [3], LUX [4], PandaX-II [5], SuperCDMS [6,7] (also see [8]). The WIMP-nucleus scattering can be divided into spin-independent (SI) and spin-dependent parts. Because of its coherent nature, the SI part is enhanced by a factor of A^2, where A is the mass number of the nucleus, making the SI scattering more relevant to the experiments (see, e.g., [9,10]).

In the coming years, the Xenon-based detectors LZ [11] or XENONnT [12] are expected to give stringent bounds on the SI DD cross sections. The projected sensitivity to probe SI WIMP-nucleon scattering reaches to about 1.4×10^{-12} pb for $M_{\rm DM} \sim$ 40-50 GeV. For $M_{\rm DM} \sim 1$ TeV, the upper bound on the DM-nucleon scattering cross section goes to about 10^{-11} pb. Obviously, such precise measurements would put strong constraints on the parameter spaces of several simple extensions of the Standard Model (SM), like the Z-portal [13–17], torsion-portal [18], Higgs-portal (H-portal) [19–29], Z'-portal [30–38], or the pseudoscalar-portals [39–41] etc. (see [42] for a review). Among these, only the torsion-portal, Z'-portal with only axial couplings and the pseudoscalar-portal models are somewhat preferred for fermionic DMs as the current null result in the SI DM-nucleon scattering can be accommodated. Otherwise, most popular TeV-scale DM models are already (or expected to be) under pressure from the existing (and future) SI DD limits, as the DM mass is pushed towards TeV values.

There have been some attempts to identify ways that would allow the WIMP scenarios to explain the non-observation of any DM signal above the neutrino floor. The Higgs mediated models have drawn significant attention in this regard [19–29,43–51]. They are relevant in many favoured beyond-the-SM (BSM) scenarios. The SM-like Higgs scalar can lead to large contributions at the microscopic level through its coupling with the DM and quarks. One may find parts of the parameter spaces where the couplings of the DM to Z or the Higgs boson may be highly suppressed or even zero identically. Similarly, in the models with an extended Higgs sector, destructive interference between light and heavy CP-even Higgs exchanges may lead to a cancellation in the SI scattering cross section. These are commonly known as the “Blind spots”, since, in such cases, the DD experiments would not be able to probe the DM initiated recoils [52–62]. Similarly, in a simple H-portal dark-matter model, where a complex scalar is added to the SM, a softly broken symmetry might ensure that the DM detection cross sec-

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tion vanishes at the tree level for zero momentum transfer. Here, the imaginary part of the complex scalar plays the role of the DM [63]. Ref. [64] studies such a pseudo-Nambu-Goldstone WIMP DM in the context of composite Higgs models which provide with a compelling ultraviolet motivation. Isospin-violating DM (IVDM) [65–69] is another interesting scenario where one can deal with the same concern by allowing non-identical couplings between the DM and nucleons. Since the limits on the DM-nucleon scattering cross section assume isospin conservation, they cannot be directly used and hence, effectively become much relaxed [69–71].

In this paper, we pursue a *simple-minded* phenomenological analysis where a vanishing DM-nucleon scattering cross section can be attained easily. We look for a H-portal model that can account for the null result in the DD experiments. However, instead of tuning the Higgs-DM interaction strength, we focus on the Higgs-nucleon interaction which is ultimately controlled by the Higgs-quark interactions, i.e., the Yukawa couplings of the SM quarks. Here, we consider different signs and non-SM values of the quark Yukawa coupling(s) to make the net SI DM-nucleon/nucleus scattering cross section small. As an illustrative example, we consider the possibility of a singlet scalar DM. Earlier, it has been shown that the Yukawa interactions of the SM quarks leave room for a partial cancellation in their contribution to the Higgs-nucleon coupling within the two-Higgs-doublet extension of the SM along with a real singlet scalar DM [52,53]. Implications of non-SM-like Higgs Yukawa couplings to the light quarks on the H-portal DM phenomenology have also been studied in Ref. [72] where the authors showed that the DD scattering rate can increase by up to four orders of magnitude. Here, in particular, we will show that in NP models, it is possible to have partial or even complete cancellation of the Higgs-nucleon couplings in the presence of higher dimensional operators at the TeV scale. Another potential benefit can be found, since the DM-nucleon effective coupling may become isospin-violating, in general. This conclusion can be generalized to two Higgs doublet models including the Supersymmetric (SUSY) ones.

The main difficulty in modifying the $hq\bar{q}$ interactions lies within the quark masses — in the SM, the Yukawa couplings of the fermions are proportional to the respective masses, i.e., $y_q \propto m_q/v$, where v is the vacuum expectation value (VEV) of the SM Higgs. Any deviation coming from the higher-order terms, particularly for the light fermions, are suppressed and unobservable in the SM. However, in the presence of new physics at the TeV scale, such alignments between the quark masses and Yukawa matrices can be relaxed (see [73,74] for discussions within the context of effective operators and [75], for the radiative generation of such a higher dimensional operator). Hence, the quark Yukawa couplings to the physical Higgs boson h, especially for the light ones (u, d, c, s), may take any value permissible by the measurements at the LHC. For the third-generation fermions and the vector bosons, these couplings are rather tightly constrained and have already been measured within 10%-20% of their SM predictions at the LHC [76]. But, the same for the first two generations of light quarks are not yet measured directly — significant deviations from their SM values are still allowed [77,78] and can offer interesting phenomenology [75,79–82]. Interestingly, for the light quarks, negative Yukawa couplings are also allowed [76,77] and have drawn some interests in recent times [83,84]. This, in the present context, is highly welcome as it allows the SI DM-nucleon scattering cross section to become vanishingly small through a cancellation among the different quark contributions to the Higgs-nucleon interaction. Note that our results are valid for any TeV scale DM model with the Higgs scalar as the dominant source for SI scatterings, since the small SI DM-nucleon scattering cross section does not depend on how the DM interacts with the Higgs scalar.

This paper is organized as follows. We first review the prerequisites for calculating the SI DM DD cross section in Sec. 2. Then in Sec. 3, we discuss the general framework where the quark Yukawa couplings to the physical Higgs field can be attuned. We illustrate how a specific type of higher dimensional operator could help. Then we discuss a TeV scale extension of the SM with some vector-like (VL) quarks where the desired dimension-6 operator is generated effectively at the tree level. A real scalar is assumed to be the DM in this model. We also discuss the existing LHC constraints and the possibility of probing the required non-SM light-quark Yukawa couplings at the colliders. In Sec. 4, we present our numerical results. There we consider the Yukawa couplings as free parameters and show how a negative coupling (in this case, of the second-generation quarks) can be helpful in explaining the null results in the DD experiments. Finally, we conclude in Sec. 5.

2. Spin independent direct detection cross section: Higgs exchanges

For a real scalar dark matter ϕ, the SI DM-nucleon elastic scattering cross section at zero momentum transfer can be written as,

$$\sigma_{\rm SI}^{\phi-N} = \frac{m_r^2 f_N^2}{4\pi M_\phi^2}, \tag{1}$$

where, M_ϕ is mass of the DM and m_r is the DM-nucleon reduced mass. At the microscopic level, the effective DM-quark scattering can be read from the following interaction term:

$$\mathscr{L}_q^{\rm SI} = f_q(\phi\phi)(\bar{q}q), \tag{2}$$

where q denotes a quark and f_q, the corresponding scattering coefficient. In a simple model of H-portal DM, we can set $f_q = \lambda_\phi\, y_q/m_H^2$ where y_q is the Yukawa coupling of the quark q and λ_ϕ is the effective DM coupling to the SM Higgs boson. Similarly, the DM-nucleon interaction can be expressed as,

$$\mathscr{L}_N^{\rm SI} = f_N(\phi\phi)(\bar{N}N), \tag{3}$$

where $f_N = \lambda_\phi\, \lambda_N/m_H^2$ and $\lambda_N = m_N \sum_q (y_q f_q^{(N)}/m_q)$ with m_N being the nucleon mass. The effective interaction of the DM particle with a nucleon (proton or neutron) can be obtained from the expectation value of the operator in Eq. (2) with respect to the initial and final nucleon states ($N \equiv p$ or n) [85]. Here, one may use the fact that nucleon mass is determined from the trace of the energy-momentum tensor. Generically, for $q = u, d, s$ the factor $f_q^{(N)}$ can be expressed as

$$\langle N|m_q\bar{q}q|N\rangle = m_N f_q^{(N)}. \tag{4}$$

For the heavier quarks, $f_q^{(N)}$ can be evaluated through one-loop contributions due to scattering off gluons [85,86]:

$$f_{c,b,t}^{(N)} = \frac{2}{27} f_G^{(N)} = \frac{2}{27}\left(1 - \sum_{q=u,d,s} f_q^{(N)}\right). \tag{5}$$

In fact, to the leading order, the effective $HG^{\mu\nu}G_{\mu\nu}$ vertex at small momentum transfer can also be used for the above computation [87]. The dominant QCD corrections should also be taken into account in that case. Altogether, one can cast the effective DM-nucleon scattering coefficient as [29,88],

$$\frac{\lambda_N}{m_N} = \sum_{q=u,d,s} f_q^{(N)} \frac{y_q}{m_q} + \frac{2}{27} f_G^{(N)} \sum_{q=c,b,t} \frac{C_q y_q}{m_q}. \tag{6}$$

The parameters, $f_q^{(N)}$ $(q \in u, d, s)$ can be determined from lattice QCD calculations [89]. For the heavier quarks, the leading order QCD correction becomes $C_q = 1 + 11\alpha_s(m_q)/4\pi$. We use the following values of $f_q^{(N)}$ [89,90],

$$f_u^p = 0.0153, \quad f_d^p = 0.0191, \quad f_s^p = 0.0447,$$
$$f_u^n = 0.0110, \quad f_d^n = 0.0273, \quad f_s^n = 0.0447 \tag{7}$$

which lead to $f_G^{(N)} \sim 0.92$ (ignoring the differences between nucleons $N = p,\ n$).[1] It should be noted that the above numerical values are subject to some uncertainties as they are evaluated using the hadronic data.

As indicated in the Introduction, we allow the light-quark Yukawa couplings to deviate from their respective SM values, and more importantly, allow them to attain negative values which are not in violation with any experimental observation so far. We use this freedom to delineate the tentative range for y_s or y_c in terms of the other SM-like Yukawa couplings so that $\lambda_N = 0$ can be achieved. Since C_q is very close to 1, here, we may simply assume $C_q = 1$ to get a qualitative picture. Thus, by substituting the values for $f_q^{(N)}$ in Eq. (6), one would get a typical regime for the Yukawa couplings where λ_N would be vanishingly small, i.e.,

$$y_s = -\frac{m_s}{f_s^{(N)}}\left(f_u^{(N)}\frac{y_u}{m_u} + f_d^{(N)}\frac{y_d}{m_d}\right); \quad y_c = -m_c\left(\frac{y_b}{m_b} + \frac{y_t}{m_t}\right). \tag{8}$$

In this illustrative example, we have allowed only y_s and y_c to take non-SM values for simplicity and further assumed that y_s cancels the light quark contribution and y_c cancel the heavier quark contribution separately. Obviously, this assumption is only a choice that we make for the illustration and not a requirement. One may tune any one or more light-quark Yukawa couplings to attain $\lambda_N = 0$. The above equation simplifies to

$$y_s = -0.770\, y_s^{\rm SM}\,; \quad y_c = -2\, y_c^{\rm SM}\,, \tag{9}$$

for $N = p$. Notably, for this set of parameters, $\lambda_p \neq \lambda_n$, implying that the DM-neutron scattering cross section will not vanish identically, and hence, a degree of isospin violation would be observed. Furthermore, these changes are potentially insignificant to modify the effective ggh vertex in the SM.

We will further discuss about this possibility in Sec. 4. A similar cancellation condition can be achieved for the neutron as well by using the respective form factors where one finds $y_s = -0.857\, y_s^{\rm SM}$. If one only allows either y_c or y_s to take values such that the DM-nucleon effective coupling in Eq. (6) vanishes, then a relatively larger negative values would be required, which may need some fine-tuning. Similarly, y_u and/or y_d can also be considered to be negative to achieve the same result.

3. Non-SM light-quark Yukawa couplings: examples and experimental tests

In this section, we illustrate how the non-SM-like Yukawa couplings can be generated through dimension-6 operators at the tree level in an effective theory framework. The current LHC measurements prevent large variations in $y_{t,b}$ [94] but leave space for deviations in the Yukawa couplings of the first two generations of quarks (for large changes to the top Yukawa couplings, see [79]). A discussion on the collider tests of non-SM values of y_q follows in the next section.

3.1. Non-SM-like light-quark Yukawa couplings and higher dimensional operators

In this example, we include a particular type of effective dimension-6 operators at some NP scale Λ in the quark Yukawa interaction Lagrangian,

$$\mathscr{L} \supset -Y_u \bar{q}_L \tilde{H} u_R - Y_d \bar{q}_L H d_R + \Delta\mathscr{L}_{eff} + H.c., \tag{10}$$

where

$$\Delta\mathscr{L}_{eff} = \frac{H^\dagger H}{\Lambda^2}\left(Y_H^u \bar{q}_L \tilde{H} u_R + Y_H^d \bar{q}_L H d_R\right). \tag{11}$$

Here, $Y_H^{u,d}$ are the Wilson coefficients determined by the details of the NP model. In general, Y_q and Y_H^q are assumed to be 3×3 matrices in the generation space. The SM Higgs doublet is denoted by H ($\tilde{H} \equiv i\tau_2 H^*$), q_L is the left-handed $SU(2)$ quark doublet and u_R, d_R are the right-handed up- and down-type quarks, respectively. After the electroweak symmetry breaking (EWSB), using $H = \begin{pmatrix} 0 \\ h + v \end{pmatrix}$ with $v \simeq 174$ GeV, one obtains the quark mass matrix M_q and the corresponding Yukawa coupling matrix in the mass basis by considering unitary rotations of both left- and right-handed quark fields. For simplicity, we ignore flavour mixings and assume diagonal NP Yukawa couplings, i.e., $Y_H^q = \mathbb{I}_{3\times 3}$ in the same basis. Though this is not true in general, it suffices for the present purpose. Now, we see that for a quark (q), its physical mass (m_q) and Yukawa coupling (y_q) to the physical Higgs boson (h) become non-aligned, i.e.,

$$m_q = v\left(Y_q - \varepsilon Y_H^q\right), \tag{12}$$
$$y_q = \left(Y_q - 3\varepsilon Y_H^q\right) = \frac{m_q}{v} - 2\varepsilon Y_H^q \tag{13}$$

where, $\varepsilon \equiv (v/\Lambda)^2$. Assuming $\Lambda \sim$ TeV and $Y_H^q \simeq \mathscr{O}(1)$, a few comments are in order.

- It is clear that, when a higher-order operator [like the one shown in Eq. (11)] is added to the SM Lagrangian, the fermion mass and Yukawa coupling become two independent quantities in the physical basis. In other words, it gives us the freedom to modify the quark-Higgs Yukawa couplings without perturbing the quark masses.
- The sign of the Yukawa couplings y_q depends on the sign of the Wilson coefficients Y_H^q. In particular, for the first two generations of quarks (u, d, s, c), one may find that $m_q/v \ll \varepsilon Y_H^q$, thus the respective Yukawa couplings may naturally become negative.
- To achieve the correct size and sign of m_q with a negative y_q [see e.g., Eq. (8)], one may use Eqs. (12) and (13) to obtain,

$$Y_H^q\left(\frac{v}{\Lambda}\right)^2 > \frac{m_q}{2v}. \tag{14}$$

  This sets an upper bound on the NP scale Λ — for a perturbative choice of $Y_H^q \sim \mathscr{O}(1)$, the NP scale Λ should not be much larger than a few TeV. For example, Eq. (14) implies that Λ should be lower than about 2.9 TeV if we take the charm quark (i.e., set $m_q = m_c^{\rm SM}$ and $Y_H^q = 1$). On the contrary, $y_q > 0$ can only set a lower bound on Λ. Thus the choice of negative values for the Yukawa couplings is more natural and predictive compared to the positive values.
- Usually, if one considers with full generality, Y_q and Y_H^q cannot be diagonalized simultaneously in the mass basis. As a result, Higgs mediated flavour changing neutral couplings among the SM quarks, which are otherwise extremely constrained by

[1] One gets slightly different values from chiral perturbation theory [29,91–93].

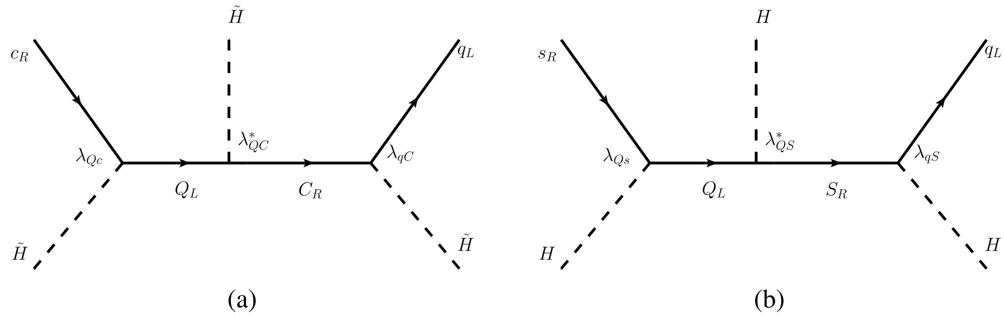

Fig. 1. Dimension-6 operators relevant for (a) $\tilde{H}^\dagger\tilde{H}c_R \to q_L\tilde{H}$ and (b) $H^\dagger H s_R \to q_L H$ process initiated at the tree level after integrating out heavy VL quarks.

experiments, may appear. However, such flavour changing couplings may be suppressed if a definite flavour symmetry or some flavour selection rules are applied [82]. Since it hardly has any impact in the present case, we will not discuss it anymore.

3.2. Singlet scalar DM and negative light-quark Yukawa couplings

Here, we consider a specific realization where the aforesaid dimension-6 operators are generated through some underlying NP that includes new heavy VL particles. VL fermions have a rich phenomenology and can be found in various popular NP scenarios, e.g., in some SUSY extensions [95–98], composite Higgs models [99–102], warped extra-dimension models [103–108], little Higgs models [109–112], etc.

In general, the effective Lagrangian can be considered to have three parts:

$$\mathscr{L}_{eff} = \mathscr{L}_{\rm SM} + \mathscr{L}_{\rm NP} + \mathscr{L}_{\rm DM} \tag{15}$$

where $\mathscr{L}_{\rm SM}$ is the usual SM part, $\mathscr{L}_{\rm NP}$ includes the new physics terms (except the DM part) responsible mainly for the nonalignment of the Yukawa couplings of the physical Higgs boson to SM quarks once the heavy fields are integrated out, and $\mathscr{L}_{\rm DM}$ refers to the terms involving DM interactions. Usually, these three parts may couple with each other in specific scenarios. For example, here we will see the NP-SM couplings producing the necessary deviations in Yukawa couplings at the tree level.

As mentioned in Sec. 2, we consider a real singlet scalar ϕ as the DM particle while the SM Higgs acts as the portal. An additional Z_2 symmetry ensures the stability of ϕ. Apart from the usual kinetic term in $\mathscr{L}_{\rm SM}$, the potential term for the model can be written as,

$$V = \frac{1}{2}\mu_\phi^2\phi^2 + \lambda_{H\phi}(H^\dagger H)\phi^2. \tag{16}$$

In principle, there should be a self-interacting ϕ^4 term also. However, we can safely ignore it, since it will have no effect in determining either the relic density or the DD cross section. After EWSB, ϕ doesn't get any VEV. Hence, the ϕ-mass term can be defined as,

$$M_\phi = \sqrt{\mu_\phi^2 + 2\lambda_{H\phi}v^2}. \tag{17}$$

The interaction terms connecting the dark sector with the SM are $(2\lambda_{H\phi}v)h\phi^2$ and $\lambda_{H\phi}h^2\phi^2$. The phenomenology of this model has been widely studied. But as discussed earlier, it is quite difficult to satisfy the current direct detection bounds with this simple extension and this motivates us to propose $\mathscr{L}_{\rm NP}$.

In $\mathscr{L}_{\rm NP}$, we introduce the VL fermions and argue that the non-SM-like Yukawa couplings for the light quarks can be generated at the tree level once the heavy VL fermion degrees are integrated out. Based on our previous discussion it is clear that we want non-SM-like Yukawa couplings, specifically for the $2^{\rm nd}$ generation quarks to the SM Higgs scalar. Thus, we consider a SM-like set up to specify the underlying NP theory [79,82] that includes only one generation of VL quarks: an $SU(2)$ VL quark doublet $Q = (C, S)(3, 2, 1/6)$ and the corresponding up-type and down-type $SU(2)$ singlets $C(3, 1, 2/3)$ and $S(3, 1, -1/3)$, carrying the same quantum numbers as the SM quark doublets and singlets. Further, we assume the VL quarks are in their mass basis, with $M_{Q,C,S} \gtrsim 2$ TeV, putting them well above the current LHC bounds. In general, we may write down the NP Lagrangian for the interactions among the VL quarks and the SM states, as follows.

$$\begin{aligned} -\mathscr{L}_{\rm NP} = \; & \left(\lambda_{QC}\,\bar{Q}_L\tilde{H}C_R + \lambda_{QS}\,\bar{Q}_L H S_R\right) \\ & + \left(\lambda_{qC}\,\bar{q}_L\tilde{H}C_R + \lambda_{qS}\,\bar{q}_L H S_R\right) \\ & + \left(\lambda_{Qc}\,\bar{Q}_L\tilde{H}c_R + \lambda_{Qs}\,\bar{Q}_L H s_R\right) + H.c. \end{aligned} \tag{18}$$

Here q_L refers to the second-generation SM quarks. This Lagrangian can lead to the desired dimension-6 operators in Eq. (11) after integrating out heavy the VL quarks. The dimension-6 couplings Y_H^c, Y_H^s, as defined in Eq. (11), can be obtained from the diagrams shown in Fig. 1 and are given below[2]:

$$Y_H^c = \lambda_{qC}\lambda^*_{QC}\lambda_{Qc}\ , \quad \Lambda = \sqrt{M_C M_Q}\ , \tag{19}$$

$$Y_H^s = \lambda_{qS}\lambda^*_{QS}\lambda_{Qs}\ , \quad \Lambda = \sqrt{M_S M_Q}\ . \tag{20}$$

Thus, if all the VL quarks have masses $M_C \sim M_S \sim M_Q \sim 2$ TeV and the new physics couplings $\lambda_{NP} \sim \mathscr{O}(1)$, the Yukawa couplings of the second-generation quarks can be considered for modification, as shown in Eq. (13).

3.3. Tests of the non-standard light-quark Yukawa couplings at the LHC and beyond

The SM Higgs couplings to the massive vector bosons, the third generation quarks and the τ-lepton and its effective couplings to two gluons or two photons are known with some accuracy (roughly 10% – 20%) [76]. Recently the ATLAS and CMS collaborations announced the first observation of a Higgs decaying to two μ's [113,114]. However, how the Higgs couples to the other light fermions has never been probed experimentally – the LHC is yet to observe a direct Higgs decay to a pair of electrons or any of the first two generation of quarks.

[2] In general, the NP couplings are all 3×3 matrices in the flavour space, though for the present purpose, it suffices to assume $\lambda_{NP} = \mathbb{I}_{3\times3}$. The SM quark fields are also assumed to be in their physical mass basis. Usually, with the VL quarks the CKM matrix is extended and the SM 3×3 CKM block is, in principle, no longer unitary, though the deviation is marginal [82]. This results into 3×3 diagonal matrices for the Wilson coefficients $Y_H^{u_i}, Y_H^{d_i}$ in Eq. (11).

Because of the Higgs mechanism, the Yukawa coupling of a fermion is proportional to its mass in the SM. This makes the Yukawa couplings of the light fermions harder to probe. However, with the high luminosity run of the LHC (HL-LHC), there could be some chance of probing them. In recent times, several possibilities have been explored for testing these couplings at the LHC [83,84,115–129], especially for c and also for s-quarks, in parallel to the ongoing experimental efforts [130–135]. Some of these proposals consider looking at rare Higgs decays to light flavoured mesons (like J/ψ or ϕ, etc.) [115,117,120,126]. Even though these can offer clean signals, such strategies suffer from low signal rates due to the small decay rates involved. There are channels with relatively larger signal cross sections [like $pp \to W/Zh \to W/Z(cc)$ [116,118,119], $pp \to hc$ [121], $pp \to hh \to (cc)(\gamma\gamma)$ [128] etc.] that require light-quark jet tagging. Ref. [125] uses a refined triggering strategy and some machine learning techniques to probe $pp \to h \to c\bar{c}\gamma$ and estimates the HL-LHC reach as $\left|y_c/y_c^{\rm SM}\right| < 8$. The process $pp \to h\gamma$ is used in Ref. [129]. Ref. [123] points out the possibility of using the charge asymmetry in $pp \to W^{\pm}h$ to constrain the light-quark Yukawa couplings. Ref. [127] employs a combination of the above ideas. Ref. [124] looks at the Higgs $p_{\rm T}$ distribution in $pp \to h + j(j_b) \to \gamma\gamma + j(j_b)$. In the gluon-gluon fusion process, the Higgs is mostly produced centrally, i.e., about $y_h \approx 0$. For a non-negligible y_u, the Higgs would also be produced via $u\bar{u}$ fusion. However, since only u is a valance quark of proton, not $\bar{u}$, production from the $u\bar{u}$ fusion would peak in forward region, i.e., around higher $|y_h|$. Hence, Ref. [122] considers the idea of using both $p_{\rm T}$ and rapidity distributions of the Higgs to obtain bounds on the light Yukawa couplings. Ref. [78] considers triple heavy vector boson production as a probe of the first-generation light-quark Yukawa couplings. These proposals are somewhat competitive in nature as the projected limits obtained for either the 300 fb^{-1} LHC or HL-LHC are similar. Here, as an estimation, we quote the projected reach in the y_q values ($q = u, d, s, c$) at the LHC with 3000 fb^{-1} of integrated luminosity from Ref. [77]:

$$\begin{aligned} &|y_u| < 560\, y_u^{\rm SM}, \quad |y_d| < 260\, y_d^{\rm SM}, \\ &|y_s| < 13\, y_s^{\rm SM}, \quad |y_c| < 1.2\, y_c^{\rm SM}. \end{aligned} \tag{21}$$

Notice that the limits are put on the absolute values of the Yukawa couplings as the processes involved are insensitive to the sign of the light-quark Yukawa couplings. There are, however, some processes that have some sensitivity on the signs. In the process where a Higgs is produced with a jet ($pp \to hj$), the shape of the $p_{\rm T}$ distribution of the Higgs depends on the production mode. Ref. [83] utilizes this to estimate that the HL-LHC can restrict $y_c/y_c^{\rm SM}$ to be within $[-0.6, 3]$. Here, the interference between c- and t-quark loops gives rise to a term linear in y_c in the cross section of the $gg \to hj$ subprocess, making it somewhat sensitive to the sign of y_c. Ref. [84] also considers the $p_{\rm T}$ distribution of Higgs. They consider the process $pp \to h \to 4\ell$ at the next-to-leading order (NLO) where the quark fusion and gluon fusion subprocesses can interfere. For the HL-LHC, they find $-1550 < y_u/y_u^{\rm SM} < 700$ and $-800 < y_d/y_d^{\rm SM} < 300$.

There are other ways to probe the light-quark Yukawa couplings than at the LHC or HL-LHC. For example, Ref. [136] indicates that FCNC transitions in the Kaon sector could restrict the down-quark Yukawa coupling to be within the rage $0.4 < |y_d/y_d^{\rm SM}| < 1.7$. It is possible to probe the light-quark Yukawa couplings by measuring isotope shifts, that are affected by Higgs exchange, in atomic clock transitions [137]. Ref. [72] points out that a discovery of H-portal dark matter could let us put bounds on these couplings.

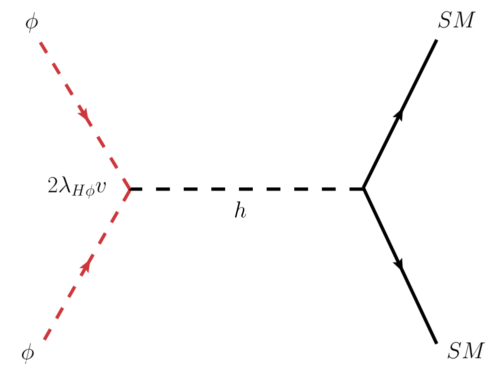

Fig. 2. Dark matter annihilation into SM particles through Higgs mediated s-channel process. In t-channel, this diagram contributes to $\sigma_{\rm SI}$.

The future colliders, especially the leptonic ones, could also offer us a better handle in measuring the Higgs couplings in general [77]. This is mainly because the lepton colliders are clean and allow us to reconstruct the processes far more accurately than their hadron counterparts. For example, Ref. [138] considers probing via hadronic event shapes at lepton colliders and shows that light-quark Yukawa couplings greater than $|0.09 \times y_b^{\rm SM}|$ might be excluded in an e^+e^- collider of centre-of-mass energy 250 GeV with an integrated luminosity of 5 ab^{-1}.

In the subsequent analysis, we compute $\sigma_{\rm SI}$ with non-SM values of y_c and y_s, especially with negative values of y_c and y_s, so that the SI scattering cross section may become vanishingly small. Hence, the future generation experiments like XENONnT or LZ could only assert our proposal through their blindness to find any signal in the $\sigma_{\rm SI}$-$M_{\rm DM}$ plane. But, on the other hand, parts of the parameter space with non-SM values for y_c (this includes the negative values as well) can be tested at the HL-LHC as claimed in Refs. [77,83]. This complementarity between the DM and LHC searches might help in testing our proposal.

4. Relic density and direct detection of DM

For numerical analysis, we use the code micrOMEGAs [88,139] to evaluate the relic density and DD cross section. The dominant QCD corrections in the SI DM-nucleon scattering are already included in the code. In our model, the main two free parameters are M_ϕ and $\lambda_{H\phi}$. Additionally, we consider y_c and y_s also as free parameters. In our computation, we set $\lambda_{H\phi} = 0.02$. The valid parameter space should comply with the observed relic abundance data [140,141],

$$\Omega_{\rm DM} h^2 = 0.1198 \pm 0.0012. \tag{22}$$

For the singlet scalar DM ϕ, one can easily solve the Boltzmann equation to get the corresponding relic abundance. The Boltzmann equation is given by,

$$\frac{dn}{dt} + 3\mathscr{H} n = -\langle \sigma_{eff}\, v \rangle (n^2 - n_{eq}^2) \tag{23}$$

where $\mathscr{H}$ is the Hubble constant and $\langle \sigma_{eff}\, v \rangle$ is the thermal averaged cross section of the DM annihilation to the SM particles. In this scenario, only the s-channel process shown in Fig. 2 keeps the DM in thermal equilibrium.

The variation of $\Omega_\phi h^2$ as a function of the DM mass is depicted in Fig. 3. The dependence of the relic density on the variations of charm and strange quark Yukawa couplings is negligible. The blue dotted line in the figure represents the central value of the DM relic density obtained from the PLANCK data. We see that, in addition to the resonance dip, this model is also able to satisfy the

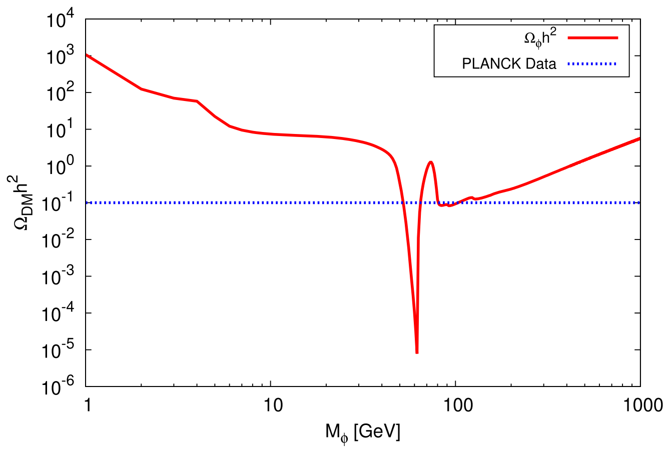

Fig. 3. Variation of relic abundance as a function of M_ϕ.

astrophysical data nicely for $M_\phi \sim 100 \pm 10$ GeV. Even though the present DD bounds exclude a real singlet scalar DM in this region, the non-SM Yukawa couplings allow us to retain it, as we see from Fig. 4.[3]

Fig. 4 validates our qualitative conclusions from Sec. 2. Here, we show the variation of SI DD cross section as a function of M_ϕ. We consider two situations: (i) both y_c and y_s can assume non-SM values while all other Yukawa couplings are fixed to their SM values and (ii) only y_c can assume non-SM values. One may simply use Eq. (6) to find y_q such that the DM-nucleon SI scattering cross-section becomes small enough to evade limits from the future-generation experiments. In the first scenario, absolute values of y_c and y_s can be much closer to their SM values while in the second case, we find that a somewhat larger value, i.e., $y_c \approx -2.9\, y_c^{\rm SM}$ would be necessary. Overall, the negative values of light-quark Yukawa couplings help to accommodate the DD bounds. For example, we see in Fig. 4(a) that there is an $\mathcal{O}(10^{-9})$ suppression between the SM line (red) and the exact cancellation line (blue). To illustrate, we can see that at $M_\phi = 100$ GeV, $\sigma_{\rm SI}^{\phi-p}(y_c = y_c^{\rm SM},\ y_s = y_s^{\rm SM}) = 1.384 \times 10^{-9}$ pb, while $\sigma_{\rm SI}^{\phi-p}(y_c = -1.875\ y_c^{\rm SM},\ y_s = -0.770\ y_s^{\rm SM}) = 1.043 \times 10^{-18}$ pb. At the backdrop of our non-SM Yukawa couplings, we have the VL quarks with mass about 2 TeV. Thus, taking the NP scale $\Lambda \sim 2$ TeV, and with $m_{c,s} = m_{c,s}^{\rm SM}$, the said non-SM values of y_c and y_s can be achieved for $Y_H^c = 1.76$ and $Y_H^s = 0.07$, respectively. Here the numerical value for y_c is slightly away from that in Eq. (9), due to the inclusion of QCD corrections. Similar results are obtained in Fig. 4(b) as well. For illustrating the functional dependence of $\sigma_{SI}^{\phi-p}$ with the 2nd generation Yukawa couplings, we present Fig. 4(c) and Fig. 4(d). In the first plot, variation is shown over $y_s/y_s^{\rm SM}$ when $y_c = -1.875\ y_c^{\rm SM}$ and in the 2nd plot, the same variation is shown over $y_c/y_c^{\rm SM}$ when all other Yukawa couplings are SM-like, respectively. One may easily find out the correct numerical value to obtain the desired cancellation in the DM-nucleon cross-section. Fig. 4(e) shows the variation of $\sigma_{SI}^{\phi-n}$ with respect to M_ϕ, for the same set of y_c and y_s where $\lambda_p \to 0$, i.e., with $y_c = -1.875\ y_c^{\rm SM}$ and $y_s = -0.770\ y_s^{\rm SM}$, keeping all other quark Yukawa couplings fixed at their SM values. The blue line in Fig. 4(e) shows that even though $\sigma_{SI}^{\phi-n}$ can't be vanishingly small for this set of parameter points, the DM-neutron scattering cross section can be below the proposed direct detection bounds for $M_\phi \geq 50$ GeV. This, clearly, reflects isospin violation to a large extent, following from the constraint $\lambda_p \to 0$.

[3] There can be specific realizations of WIMP models where a sub-TeV DM can be viable [142,143].

However, the IVDM scenario can be realized in a more general way and a large violation can be observed for moderate to large non-SM-like first-generation Yukawa couplings. Since $f_{u,d}^{(N)}$ values are different for neutron and proton, one may easily get a parameter space where $f_p \neq f_n$. A particular useful scenario appears when f_n/f_p takes a negative value, as it then offers some significant cancellations in the DM-nucleus scattering cross section. Though a larger value of y_d or y_u can easily lead to $f_n/f_p > 0$, but negative values can only be achieved in a narrow domain which can be computed using Eq. (6). For example, considering only y_d to assume non-SM values, the following range for the same coupling can be observed.

$$-\left(\frac{X^p}{f_d^p}\right) < \frac{y_d}{y_d^{\rm SM}} < -\left(\frac{X^n}{f_d^n}\right), \tag{24}$$

where,

$$X^{(N)} = \sum_{q=u,s} f_q^{(N)} + \frac{2}{27} f_G^{(N)} \sum_{q=c,b,t} C_q\,, \qquad [N = p, n].$$

Note that $X^p > X^n$ for $f_q^{(N)}$ shown in Eq. (7). For numerical estimation, we can consider a benchmark ratio, e.g., $f_n/f_p \approx -0.7$. This particular value has some importance to relax tensions between different results for low DM mass. One may easily check that, $y_d/y_d^{\rm SM} \approx -11.6$ can lead to the above ratio. Finally, we note that in the limit of vanishing DM-nucleon couplings, the two-nucleon currents might become important [144].

5. Conclusion

Direct searches of WIMPs as a form of dark matter have been underway for a long time. Due to the present and projected experimental sensitivities towards the spin-independent direct dark matter detection cross section, the available parameter spaces in simple H-portal dark matter models are significantly reduced or threatened to be ruled out. However, there exist a few small regions where the DM interactions with nucleons can be very tiny, thus escaping the ever impinging bounds from the DM searches. One such example is the so-called “Blind spots” where either the DM couplings with Higgs scalar vanish, or there is some destructive interference among diagrams involving different neutral scalars. In this paper, we have realized another route in the same direction, where the Higgs boson couplings with the nucleons become vanishingly small. Apparently, such a requirement can be realized easily, if the light-quark Yukawa couplings are allowed to assume non-SM values in presence of some new physics; in particular, negative values — a possibility allowed by the current experiments. Adopting a phenomenological effective theory perspective, we consider some higher dimensional operators or, in particular, dimension-6 operators involving the SM fields that let the light-quark Yukawa couplings to be negative without disturbing the respective quark masses. However, in a specific theory, precise cancellation of the DM-nucleon SI direct detection cross section may be realized only within a limited ranges of the parameters. Importantly, the new physics scale is bounded and for perturbative values of new couplings, can only be of the order of a few TeV. We also consider a specific realization of this with vector-like quarks with masses about 2 TeV. In this set-up, we consider a real SM-singlet scalar as the DM candidate. In the absence of any discovery, generally, such a simple DM set up would be excluded completely for $M_{\rm DM} \leq 1$ TeV from the projected sensitivity of the proposed LZ or XENONnT experiments. Here we observe that resultant SI DM-nucleon scattering cross section can be made vanishingly small for all values of DM mass. Needless to say, our observation would be unchanged for any other DM candidate in

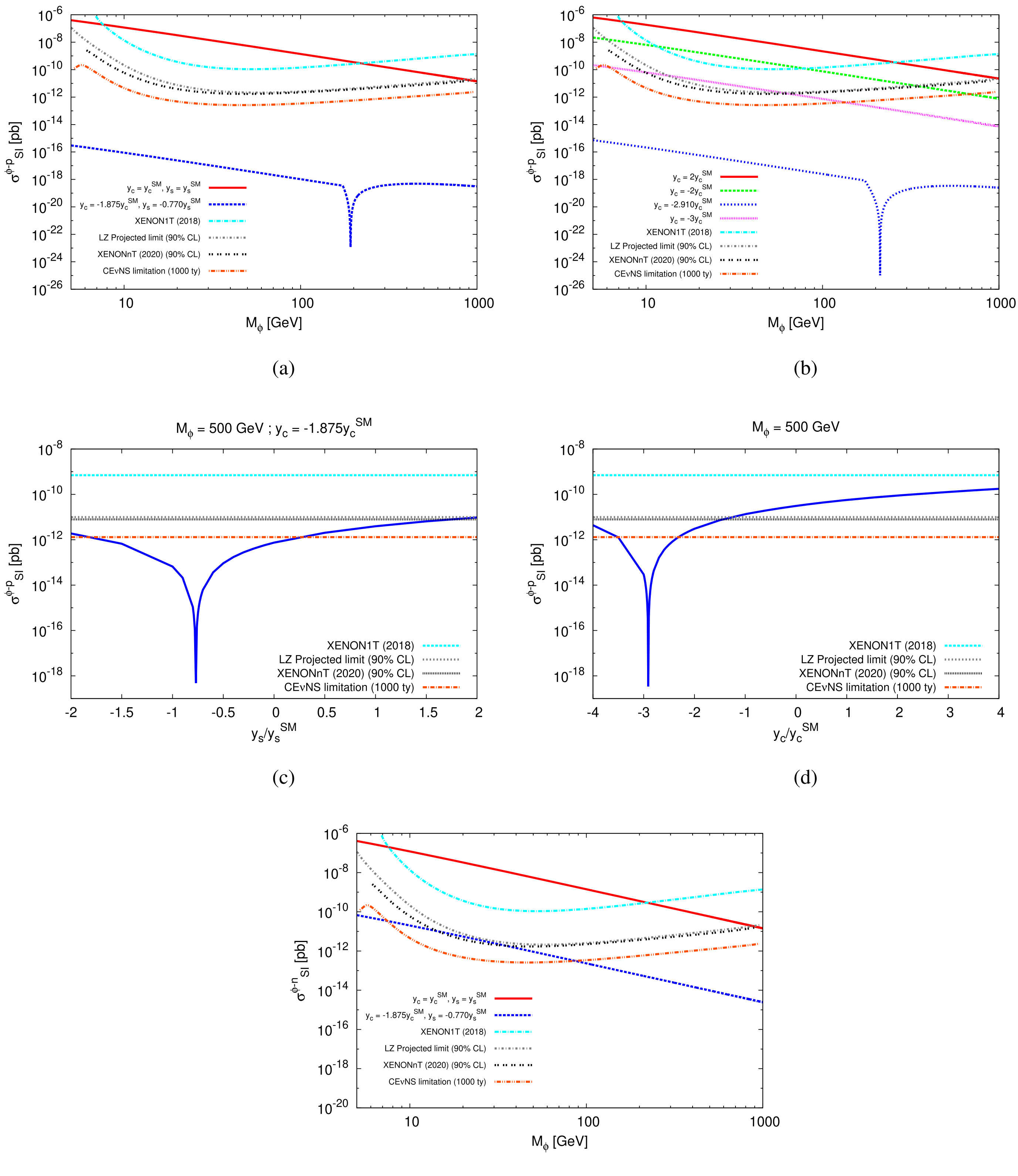

Fig. 4. Variation of the $\sigma_{\mathrm{SI}}^{\phi-p}$ with respect to M_ϕ and y_s or y_c. In (a), y_c and y_s are assumed to have non-SM values while all other Yukawa couplings are fixed to their SM values and in (b), only y_c can assume non-SM values. In (c), the variation of $\sigma_{\mathrm{SI}}^{\phi-p}$ has been plotted as function of y_s/y_s^{SM} when only y_c assumes non-SM value and in (d), only y_c can have non-SM values, while $M_\phi = 500$ GeV. In (e), we show the variation of $\sigma_{\mathrm{SI}}^{\phi-n}$ as a function of M_ϕ for the non-SM values of y_c and y_s for which $\sigma_{\mathrm{SI}}^{\phi-p}$ becomes minimum.

the Higgs-portal models or in the models where the Higgs produces the dominant contribution to the direct detection process, since our argument does not depend on the DM-DM-Higgs couplings and, hence, it can take $\mathcal{O}(1)$ values as well. Isospin-violation in the DM-nucleon scattering can also be realized for negative values of the first-generation light-quark Yukawa couplings. Usually, the exclusion limits are obtained assuming isospin-conserving dark matter and hence can be much relaxed when the DM couples differently to protons and neutrons. Probing such non-standard values of the Yukawa couplings of the first two generations of quarks is hard even for the HL-LHC, though there exist studies that aim to narrow down the allowed values of the light-quark Yukawa couplings, even the negative values. For example, it has been argued that, at the LHC, it might be possible to pin the charm-quark Yukawa coupling within $[-0.6, 3]y_c^{\rm SM}$ with 3000 fb^{-1} of integrated luminosity in a largely model independent manner. Hence, even though the future generation of dark matter search experiments based on the dark matter-nucleon scattering is blind to our proposal, it might be tested at the HL-LHC.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

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