Heavy-quark and Gluon -philic Real Dark Matter

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ABSTRACT: We explore the phenomenological viability of the real spin half, zero and one dark matter candidates which primarily interact with the third generation heavy-quarks and gluons through the twenty-two gauge invariant higher dimensional effective operators. The lower-limits on the corresponding Wilson coefficients are obtained from the relic density constraint $\Omega_{DM}h^2 \approx 0.1198$ and their contributions to the thermal averaged annihilation cross-sections are found to be consistent with the predicted upper bound on the annihilation cross-section in the $b\bar{b}$ mode from FermiLAT and H.E.S.S. experiments.

The tree-level gluon-philic and one-loop induced heavy-quark-philic DM - nucleon direct-detection cross-sections are analysed. The non-observation of at least one recoiled Xe nucleus event for spin-independent DM-nucleus scattering in XENON-1T sets the upper limits on the eighteen Wilson coefficients. Our analysis validates the real DM candidates for the large range of accessible mass spectrum below 2 TeV for all but one interaction induced by the said operators.

KEYWORDS: Dark-matter, Effective Operators, heavy-quark, Wilson Coefficient

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1 Introduction

Regardless of the unequivocal astrophysical evidence from rotation velocity curves [1] via mass-to-luminosity ratio, Bullet Cluster [2], and precision measurements of the Cosmic Microwave Background (CMB) from WMAP [3] and PLANCK [4] satellites etc., we have no clue about the fundamental nature of dark-matter (DM), which forms the dominant matter component of the Universe. The dedicated experiments for the direct-detection of DM such as LUX [5], XENON-1T [6], DarkSide50 [7], PandaX-4T [8] and CRESST-III [9], PICO-60 [10], PICASSO [11] are designed to measure the momentum of the recoiled atom and/or nucleus due to the scattering of DM particles off the sub-atomic constituents of the detector material [12–14]. We have yet to observe a single event. PandaX-4T [8] and PICO-60 [10] have recently lowered the upper-limits of the measured sensitivities at 90% C.L. corresponding to (a) spin-independent cross-sections at $3.3 \times 10^{-47}$ cm$^2$ for 30 GeV DM and (b) spin-dependent cross-sections at $2.5 \times 10^{-41}$ cm$^2$ for 25 GeV DM respectively. Efforts are being made to understand the nature of DM interactions by indirectly detecting DM resulting from DM pair annihilations to SM particles using space based facilities such as Fermi-LAT [15], PLANCK data [4], MAGIC [16] and some ground based large neutrino detectors such as HESS [17], Ice Cube [18], ANTARES [19], Super-Kamiokande [20] etc.

Past and ongoing experiments have constrained a plethora of viable UV complete dark matter models formulated by writing renormalizable Lagrangians with heavy non-SM mediators (spin 0$^\pm$, 1/2, 1, 2) facilitating interactions between DM and SM particles [21–31].
The models with Higgs bosons as mediators have been excluded by the recent collider experiments [32]. The analysis has also been performed in the domain of the effective field theory (EFT) for the EW-Boson-philic DM operators [33–35]. Recently, the GAMBIT collaboration performed a global analysis for signatures of SM gauge singlet Dirac fermion DM at LHC in the EFT set-up with simultaneous activation of fourteen effective operators constructed in association with light-quarks, gluons and photons (up-to mass dimension seven) where the authors have disfavored the exclusive contribution of DM $\leq 100$ GeV [36]. The collider signatures of the effective Lepto-philic DM operators have been studied for the proposed ILC through $e^+e^- \rightarrow \gamma + E_T$ [37–42] and $e^+e^- \rightarrow Z^0 + E_T$ [43, 44] channels. Sensitivity analysis for DM-quark effective interactions at LHC have been performed [45–49] in a model-independent way for the dominant (a) mono-jet $+ E_T$, (b) mono-$b$ jet $+ E_T$ and (c) mono-$t$ jet $+ E_T$ processes.

Various top-philic DM inspired models have been studied and constrained [50–56] by the cosmological relic density criteria, direct-detection and indirect-detection experiments. The interactions of heavy-quark-philic DM are also constrained by the ongoing collider experiments. The CMS [57, 58] and ATLAS [59] collaborations investigated the dominant scalar-mediated production of $t$-quark-philic DM particles in association with a single top quark or a pair of $t\bar{t}$. CMS collaboration recently conducted an exclusive search analysis for the $b$-philic DM in reference [60]. Many authors in the literature have explored mono-jet $+ E_T$ mode at the LHC for the viable signatures of the scalar current induced heavy-quark-philic DM interactions [49, 61–63]. A comprehensive search analysis for heavy scalar mediated top-philic DM models for mono-jets $+ E_T$, mono-$Z$, mono-$h$ and $t\bar{t}$ pair productions at LHC can be found in [64, 65].

The phenomenology of spin 1/2, 0 and 1 real DM deserves a special mention in the context of their contributions to the DM-nucleon scattering events in the direct-detection experiments. In the absence of the contributions from vector operators for Majorana, real scalar and real vector DM, the spin-independent DM-nucleon scattering cross-sections are found to be dominated by the respective scalar and dimension-8 twist operators [66–73]. The pseudo-scalar interactions are found to be spin-dependent and are therefore suppressed. Similar studies have also been undertaken for lepto-philic real DM candidates [37, 38].

In this context, it is worthwhile to investigate the WIMP DM phenomenology induced by heavy-quark-philic and gluon-philic scalar, axial-vector, and twist-2 real-DM operators and explore whether they satisfy the relic density criteria and other experimental constraints for a DM mass between 10 GeV and 2 TeV. We introduce the effective Lagrangian for real spin 1/2, 0 and 1 DM particles in section 2. In section 3, we discuss the phenomenology of the dark-matter. We investigate the cosmological constraints on the DM in sub-section 3.1. In sub-section 3.2 we predict and analyse the thermally averaged DM pair annihilation cross-section. Sub-section 3.3 details the computation of DM-Nucleon scattering cross-sections and expected recoiled nucleus event(s) for XENON-1T [6] set up. Section 4 summarises our study and observations.
2 Heavy-Quark and Gluon-philic DM Effective Operators

It is a fact that in a beyond-standard model (BSM) renormalizable gauge theory, as long as the energy range of DM-SM interaction is much below the mediator mass, the study of DM interactions can be restricted by the DM and SM degrees of freedom and their symmetries.

The higher dimensional effective operators are obtained from the BSM Lagrangian by writing the operator product expansion (OPE) of currents in the limit $p^2/m_{\text{Med}}^2$ << 1, where $p^\mu$ represents the four momentum of the virtual mediator of mass $m_{\text{Med}}$. This facilitates the effective contact interaction between DM and any SM third generation heavy quark/ gluon, assuming that the mediator mass scale $m_{\text{Med}}$ is of the order of the effective theory cut-off ($\sim \Lambda_{\text{eff}}$), which is much heavier than the masses of the SM and DM fields in general.

For example, in renormalizable models, the interaction between a Majorana DM $\chi$ and a third generation quark $\psi$ can be written as

$$L^{\text{Dim.}4} = \bar{\psi} (a + b \, \gamma_5) \chi \eta + \bar{\psi} \gamma^\mu (c + d \, \gamma_5) \chi \zeta_\mu + \text{h.c.}$$

(2.1)

where $\eta$ and $\zeta_\mu$ are electrically charged scalar and vector fields respectively. This interaction Lagrangian facilitates the Majorana DM - quark interaction via $t$-channel exchange of the heavy $\eta$ and/ or $\zeta_\mu$.

By expanding the propagator in powers of $p^2/m_{\text{Med}}^2$, the higher dimensional effective four-fermion interaction for spin 1/2 Majorana and heavy-quarks can be constructed. This expansion induces the scalar, pseudo-scalar, axial-vector and twist-2 operators [66]. Such effective operators can also be obtained from other kinds of renormalizable interactions (say $s$-channel interactions).

Since the interactions of the DM particles with heavy-quarks at the next-order in strong coupling constant generate the DM-gluon interactions, it is imperative to include the operators ($\propto \alpha_s/\pi$) involving a pair of Majorana DM and a pair of gluons in the effective Lagrangian. Thus, integrating the heavy degrees of freedom and following the convention of reference [72], the phenomenological effective Lagrangian for the heavy-quark-philic and gluon-philic Majorana DM is written as:

$$L_{\text{eff}}^\chi = \frac{C_q^q}{\Lambda^3} O_{\chi S}^q + \frac{C_q^q}{\Lambda^4} O_{\chi S}^q + \frac{C_p^{\chi T_1}}{\Lambda^4} O_{\chi T_1}^p + \frac{C_p^{\chi T_2}}{\Lambda^5} O_{\chi T_2}^p + \frac{C_q^{\chi A V}}{\Lambda^2} O_{\chi A V}^q$$

(2.2)

where

$$O_{\chi S}^q = m_q \bar{\chi} \chi \bar{q} q$$

$$O_{\chi S}^q = \frac{\alpha_s}{\pi} m_\chi \bar{\chi} \chi \, G^A_{\mu\nu} G^{A\mu\nu}$$

$$O_{\chi T_1}^p = \bar{\chi} i \partial^\mu \gamma^\nu \chi \, O_{\mu\nu}^p$$

$$O_{\chi T_2}^p = \bar{\chi} i \partial^\mu i \partial^\nu \chi \, O_{\mu\nu}^p$$

$$O_{\chi A V}^q = \bar{\chi} \gamma_\mu \gamma_5 \chi \bar{q} \gamma^\mu \gamma_5 q$$

(2.3)
The second rank twist tensor currents \( O_{\mu\nu} \) for the heavy-quarks and gluons are defined as

\[
O_{\mu\nu}^q \equiv i \frac{1}{2} q_L \left( D_{\mu L} \gamma_\nu + D_{\nu L} \gamma_\mu - \frac{1}{2} g_{\mu\nu} D_L \right) q_L + i \frac{1}{2} q_R \left( D_{\nu R} \gamma_\mu + D_{\mu R} \gamma_\nu - \frac{1}{2} g_{\mu\nu} D_R \right) q_R
\]

\[
O_{\mu\nu}^g \equiv \left( (G^A)_\mu^\rho (G^A)^\nu_\rho - \frac{1}{4} g_{\mu\nu} (G^A)^\rho_\sigma (G^A)^\sigma_\rho \right)
\]

(2.4a)

(2.4b)

where \( D_{\mu L} \) and \( D_{\mu R} \) are the covariant derivatives for left and right handed quarks respectively in SM. It should be noted that we did not include severely constrained dimension-8 scalar operators because they are more likely to be suppressed by \( (m_\chi m_q) / \Lambda^4 \) than leading scalar operators. We do not analyse the less competitive heavy-quark-phile and gluon-phile pseudo-scalar operators in our study as their contributions to the DM-nucleon scattering events are found to be spin and velocity suppressed. We exclude the contribution from dimension-9 twist-2 type-2 operators from our analysis because we limit our study to effectiv operators up to mass dimension-8.

We extend the domain of our analysis to include the effective contact interactions of real scalar \( \phi^0 \) and vector \( V^0_\mu \) DM candidates with a SM third generation heavy quarks and gluons. This can be obtained in a similar fashion through OPE from renormalizable BSM interactions of the SM fermion with scalar or vector DM in the presence of a heavy charged fermion(s) as shown in reference \[72\]. The effective scalar and vector DM Lagrangian are given as

\[
L_{eff} = \frac{C^q_0}{\Lambda^2} \mathcal{O}^q_0 + \frac{C^g_0}{\Lambda^2} \mathcal{O}^g_0 + \frac{C^{qV}_0}{\Lambda^4} \mathcal{O}^{qV}_0 + \frac{C^{pV}_0}{\Lambda^2} \mathcal{O}^{pV}_0 \quad \text{where}
\]

\[
\mathcal{O}^q_0 = \phi^0 \phi^0 \left( m_q \bar{q} q \right)
\]

\[
\mathcal{O}^g_0 = \alpha_8 \phi^0 \phi^0 G^A_\mu \gamma_\mu G^A_\nu \gamma_\nu
\]

\[
\mathcal{O}^{qV}_0 = \phi^0 \iota \partial^\mu \iota \partial^\nu \phi^0 \mathcal{O}^{\mu\nu}
\]

(2.5)

and

\[
L_{eff} = \frac{C^{V}_0}{\Lambda^2} \mathcal{O}^{V}_0 + \frac{C^{gV}_0}{\Lambda^2} \mathcal{O}^{gV}_0 + \frac{C^{pV}_0}{\Lambda^2} \mathcal{O}^{pV}_0 \quad \text{where}
\]

\[
\mathcal{O}^{V}_0 = \left( V^0 \right)^\mu_\mu \left( V^0 \right)_{\mu \nu} \left( m_q \bar{q} q \right)
\]

\[
\mathcal{O}^{gV}_0 = \alpha_8 \left( V^0 \right)_{\rho} \left( V^0 \right)^\rho \gamma_\mu G^A_\mu \gamma_\mu G^A_\nu \gamma_\nu
\]

\[
\mathcal{O}^{pV}_0 = \left( V^p \right) \iota \partial^\mu \iota \partial^\nu \left( V^p \right) \mathcal{O}^{\mu\nu}
\]

\[
\mathcal{O}^{gV}_0 = i \epsilon^{\mu\nu\rho\sigma} \left( V^0 \right)^\mu_\mu \left( V^0 \right)_{\rho \sigma} \left( m_q \bar{q} q \right)
\]

(2.6)

where \( \epsilon^{\mu\nu\rho\sigma} \) is the totally antisymmetric tensor with \( \epsilon^{0123} = +1 \).

Operators constructed from scalar, axial-vector and twist-2 currents of heavy-quarks and gluons corresponding to Majorana, real scalar and real vector DM are phenomenologically analysed one at a time.
3 Phenomenological Constraints on Real DM

3.1 Contribution of DM to relic density

The present day relic abundance of the DM species $n_{\text{DM}}(t)$ can be calculated by solving the Boltzmann equation

$$\frac{dn_{\text{DM}}}{dt} + 3H_0 n_{\text{DM}} = -\langle \sigma_{\text{ann}} | \vec{v} \rangle \left( (n_{\text{DM}})^2 - (n_{\text{eq}}^\text{DM})^2 \right);$$  \hspace{1cm} (3.1)

where $n_{\text{eq}}^\text{DM}$ is the DM number density at thermal equilibrium, $H_0$ is the Hubble constant, $\vec{v}$ is relative velocity of the DM pair and $\langle \sigma_{\text{ann}} | \vec{v} \rangle$ is the thermal average of the annihilation cross-section. It is customary to parametrise $\rho_{\text{DM}} \equiv \Omega_{\text{DM}} h^2 \rho_c$, where $\rho_c \equiv 1.05373 \times 10^{-5} h^2 / c^2$ GeV cm$^{-3}$ is the critical density of the universe and dimensionless $h$ is the current Hubble constant in units of 100 km/s/Mpc. Solving the Boltzman equation [74] we then get

$$\Omega_{\text{DM}} h^2 = \frac{\rho_{\text{DM}}}{\rho_{\text{critical}}} h^2 \approx 0.12 \left( \frac{2.2 \times 10^{-26} \frac{\text{cm}^3}{\text{s}}}{{\langle \sigma_{\text{ann}} | \vec{v} \rangle}} \right)^{1/2} \left( \frac{m_{\text{DM}}/T_F}{23} \right) \quad (3.2)$$

where parameter $x_F \equiv m_{\text{DM}}/T_F$ is a function of degrees of freedom of the DM $g$, the effective massless degrees-of-freedom $g_{\text{eff}}$ ($\sim 106.75$ and $86.75$ above $m_{\ell}$ and $m_b$ mass-thresholds respectively) at freeze-out temperature $T_F$:

$$x_F = \ln \left[ a(a+2) \sqrt{\frac{45}{8}} \frac{g M_{pl}}{2\pi^3} \frac{m_{\text{DM}}}{\sqrt{x_F} g_{\text{eff}}(x_F)} \right]; \quad a \sim \mathcal{O}(1) \quad \text{and} \quad M_{pl} = 1.22 \times 10^{19} \text{GeV}.$$  \hspace{1cm} (3.3)

The predicted DM relic density $\Omega_{\text{DM}} h^2$ of $0.1138 \pm 0.0045$ and $0.1199 \pm 0.0022$ by WMAP [3] and Planck [4] collaborations respectively restrict the thermal averaged DM annihilation cross-section $\langle \sigma | \vec{v} \rangle \geq 2.2 \times 10^{-26} \text{ cm}^3 \text{s}^{-1}$ as a smaller thermally averaged cross-section would render large DM abundance which will over-close the universe. We have analytically calculated the DM pair annihilation cross-sections to a pair of third generation heavy quarks and gluons in the Appendix A.

We investigate the DM relic density contributions propelled by the thermally averaged annihilation cross-sections given in the appendix B corresponding to the the scalar, axial-vector and twist-2 currents of heavy quarks for a given Majorana/ scalar/ vector DM mass and the respective Wilson coefficient. The annihilation channels induced by the scalar and twist-2 currents of gluons also contribute to the relic abundance of Majorana, scalar and vector DM. We compute the absolute lower-limits for all twenty-two Wilson coefficients $|C_{\text{DM}, i O_j / \Lambda}^{q, g}|$ in TeV$^{-n}$ ($\text{DM}_i \equiv \chi, \phi^0, V^0$ and $O_j \equiv S, AV, T_k$) by activating one operator at a time and making the sole contribution from the corresponding operator to satisfy the relic density $\Omega_{\text{DM}} h^2 \approx 0.1199 \pm 0.0022$ [4]. The lower bounds on the Wilson coefficients for a given DM mass are translated to the upper bound on $|C_{\text{DM}, i O_j}^{q, g}|^{-1/n} \Lambda$ in units of TeV. The Wilson coefficient larger than its lower limit for a given DM mass, partially satisfy the
Figure 1: The $b$-quark-philic and $t$-quark-philic Majorana DM contours for figures 1a and 1b are drawn in $m_\chi - |C_{q SA}^b|^{-1/3} \Lambda$, $m_\chi - |C_{q SA}^a|^{-1/2} \Lambda$ and $m_\chi - |C_{q TA}^b|^{-1/4} \Lambda$ planes corresponding to scalar, axial-vector and twist-2 type-I operators respectively. The gluon-philic Majorana DM contours in figure 1c are drawn in $m_\chi - |C_{q SA}^g|^{-1/4} \Lambda$ and $m_\chi - |C_{q TA}^g|^{-1/4} \Lambda$ planes for scalar and twist-2 type-1 operators respectively. Cosmologically allowed regions are shaded below the respective contours satisfying $\Omega h^2 = 0.1198$ [4].

relic density and thus may survive in cases where (a) more than one type of DM particle is allowed and/or (b) simultaneous switching of more than one effective operator.

For numerical computation of the DM relic density, we have used MadDM [75, 76]. The input model file required by the MadDM is generated using FeynRules [77, 78], which calculates all the required couplings and Feynman rules by using the full Lagrangian given in equations (2.3), (2.5) and (2.6) corresponding to Majorana, scalar and vector DM particles respectively. We have further verified and validated our numerical results from Mi-
The $b$-quark-philic, $t$-quark-philic and gluon-philic scalar DM contours for figures 2a, 2b and 2c are drawn in $m_{\phi_0} - \left| C_{q,g}^{1n} \Lambda \right|^{-1/2}$ and $m_{\phi_0} - \left| C_{q,g}^{1n} \Lambda \right|^{-1/4}$ planes corresponding to scalar and twist-2 type-2 operators respectively. Cosmologically allowed regions are shaded below the respective contours satisfying $\Omega\phi_0 h^2 = 0.1198$ [4].

In the figures 1, 2 and 3 we plot the relic density contours corresponding to Majorana, real scalar and real vector DM candidates, respectively. The contours of relic density are defined in the planes of $m_{DM} - \left| C_{DM, O_j}^{q,g} \Lambda \right|^{-1/n}$. The left and right panels at the top in all figures depict the contributions by $b$-quark-philic and $t$-quark-philic DM respectively. The lower panel in all figures depicts the contribution by gluon-philic DM.
Figure 3: The b-quark-philic, t-quark-philic and gluon-philic vector DM contours for figures 3a, 3b and 3c are drawn in $m_{\nu^0}$ planes corresponding to scalar, axial-vector and twist-2 type-2 operators respectively. Cosmologically allowed regions are shaded below the respective contours satisfying $\Omega V^0 h^2 = 0.1198$ [4].

3.2 Indirect detection of DM pairs

Today, the WIMP Dark Matter in the universe is expected to be trapped in large gravitational potential wells, which further enhances the number density of DM in the region, resulting in frequent collisions among themselves. This facilitates DM-DM pair annihilation into a pair of SM particles (photons, leptons, hadronic jets, etc) at the Galactic centre, in Dwarf Spheroidal galaxies (dSphs), Galaxy clusters, and Galactic halos. The dwarf spheroidal satellite galaxies of the Milky Way are especially promising targets for
DM indirect detection due to their large dark matter content, low diffuse Galactic γ-ray foregrounds as they travel the galactic distance, and lack of conventional astrophysical γ-ray production mechanisms. Their flux is observed by the satellite-based γ-ray observatory Fermi-LAT [80], PLANCK [4], primary cosmic rays measurements by AMS-02 [81, 82] on International Space Station and the ground-based Cherenkov telescope H.E.S.S. [17, 83–85], HAWC [86–88], MAGIC [16].

![Figure 4](a) and (b) depict the thermally averaged cross-sections for Majorana DM pair annihilation into $b\overline{b}$ and $t\overline{t}$ pairs, respectively. The contributions from axial-vector and twist-2 type-1 operators in both the panels are evaluated using their respective lower bound on $|C_{\chi AV}^{\chi AV, T_1}/\Lambda^n|$ satisfying $\Omega h^2 = 0.1199 \pm 0.0022$ [4] as shown in figure 1 and hence, the unshaded regions above the respective curves are cosmologically allowed. Regions above the the experimental limits obtained from FermiLAT [15] as well as H.E.S.S. [17] are excluded.

$t$-quark-philic DM annihilation yields a pair of $t\overline{t}$, where the $t$ ($\overline{t}$) decays into $b$ ($\overline{b}$) associated with $W^+$ ($W^-$) with 100% branching fraction. The charged gauge Bosons then decays in either leptonic or semi-leptonic or hadronic modes. The gluons and/or pair of $b\overline{b}$ produced by DM pair’s annihilation hadronized partially to neutral pions, which then decayed to photon pairs ($\pi^0 \rightarrow \gamma\gamma$) with a 99% branching fraction, yielding a broad spectrum of photons. We compare the thermally averaged DM annihilation cross-sections to the $b\overline{b}$, $t\overline{t}$ and $gg$ channels with the corresponding upper bounds derived from the observed photon spectrum in both FermiLAT [15] and H.E.S.S. [17] indirect experiments.

The analytic expressions of thermally averaged DM pair annihilation cross-sections $\langle \sigma | v^2 \rangle$ for Majorana, scalar, and vector DM are given in the Appendix B and agree with what is known for other fermions in the literature [38, 44, 89, 90]. These are derived from the cross-sections in the Appendix A in which the center of mass energy squared $s$ is expanded as $\approx 4m_{DM}^2 + m_{DM}^2 v_{DM}^2 + \frac{3}{4} m_{DM}^2 v_{DM}^4 + O(v_{DM}^6)$. It is to be noted that DM velocity $v_{DM}$ is roughly of the order of $10^{-3} c$ at the center of Galaxy and $10^{-5} c$ at dSphs.
Figures 5a, 5b and 5c depict the thermally averaged cross-sections for real scalar DM pair annihilation into $b\bar{b}$, $t\bar{t}$ and $gg$ pairs, respectively. The contributions from scalar and twist-2 type-2 operators in the panels are evaluated using their respective lower bound on $|C_{q,g}^{0,0}\phi_{0}^{0}|^{2}/\Lambda_{n}^{2}$ satisfying $\Omega_{\phi_{0}}h^{2} = 0.1199 \pm 0.0022$ [4] as shown in figure 2 and hence, the unshaded regions above the respective curves are cosmologically allowed. Regions above the experimental limits obtained from FermiLAT [15] as well as H.E.S.S. [17] are excluded.

In contrast to that of $10^{-1}c$ at freeze-out. In comparison to all chiral blind $s$-dominated processes, the leading term contributions from the $p$-wave and $d$-wave channels in $\langle \sigma |v|^{2}\rangle$ are proportional to $v_{pDM}^{2}$ and $v_{dDM}^{4}$, respectively, and thus suppressed.

For varying $m_{DM}$, we investigate the contributions to the thermally averaged DM pair annihilations cross-sections into $b\bar{b}$, $t\bar{t}$ and $gg$ pairs. Using the respective lower bound on the Wilson coefficients for a given $m_{\chi}$ as shown in figure 1, we calculate and depict the variation of cosmological lower bound on the thermally averaged Majorana DM pair...
Figures 6a, 6b and 6c depict the thermally averaged cross-sections for real vector DM pair annihilation into $b\bar{b}$, $t\bar{t}$ and $gg$ pairs, respectively. The contributions from scalar operator in the panels are evaluated using their respective lower bound on $|C_{q,g}^{V_0}/\Lambda^n|$ satisfying $\Omega V^0 h^2 = 0.1199 \pm 0.0022$ [4] as shown in figure 3 and hence, the unshaded regions above the respective curves are cosmologically allowed. Regions above the the experimental limits obtained from FermiLAT [15] as well as H.E.S.S. [17] are excluded.

annihilation cross-sections with Majorana DM mass for $\chi\bar{\chi} \to b\bar{b}$ and $\chi\bar{\chi} \to t\bar{t}$ in figures 4a and 4b respectively. We observe that the variation of chirally suppressed axial-vector operator induced $\langle \sigma |\vec{v} \rangle$ w.r.t. DM mass has a negative slope resulting from its leading term $\propto \left( \frac{m_f^2}{\Lambda^4} \right) \sqrt{1 - \frac{m_f^2}{m^2_\chi}}$ in equation (B.1b) which is in sharp contrast to almost flat-curve (in log scale) variation of that induced by the chirality conserving twist-2 type-1 operator where the leading term is $\propto \frac{m_\chi^6}{\Lambda^8}$ in equation (B.1c). In case of the twist-2 type-1 operator the suppression due to increment in the cosmological upper bound on $\Lambda$ is
compensated with the increment in the DM mass alone. The negligible contributions from the heavy-quark-philic scalar $O_{XXS}^{q}$ operator given in equation (B.1a) is attributed to the chiral suppression along with DM velocity dependence while gluon-philic scalar $O_{XXS}^{g}$ and twist-2 type-1 operators $O_{XXT_{1}}^{g}$ given in equations (B.2a) and (B.2b) respectively are $p$-wave suppressed and hence not shown in the graph.

Similarly, we plot the variation of the cosmological lower bound of the thermally averaged DM annihilation cross-section ($\phi^{0}\phi^{0} \rightarrow b \bar{b}/t \bar{t}/g g$) with scalar DM mass $m_{\phi^{0}}$ in figure 5. The $\langle |\tilde{v}| \rangle$ induced by the heavy-quark-philic twist-2 type-2 operator $O_{\phi^{0} T_{2}}^{q}$ given in (B.3b) is chirally suppressed and therefore falls sharply with an increasing DM mass as shown in figure 2. The sharp fall in the $t$-quark-philic case, on the other hand, is flattened for DM mass ranges of less than 2 TeV. The gluon-philic twist-2 type-2 operator is $d$-wave suppressed as shown in equation (B.4b). Figure 6 show the thermally averaged vector DM pair annihilation cross-sections of the third generation quarks and gluons with varying vector DM mass corresponding to the respective cosmological lower bound on the Wilson coefficient as shown in figure 3. Unlike the chiral $p$-wave suppressed axial-vector and $d$-wave suppressed twist-2 type-2 operators, we observe an appreciable contribution to the $\langle |\tilde{v}| \rangle$ from the heavy-quark-philic and gluon-philic scalar operators $\sim 10^{-26}$ cm$^{3}$s$^{-1}$.

3.3 DM-Nucleon scattering

\[\text{Figure 7: One-loop Majorana DM-gluon scattering diagrams where the blob represents the four fermionic effective interactions induced by the scalar / axial-vector / twist-2 currents of heavy quarks.}\]

In direct-detection experiments, the scattering of DM particles can be broadly classified as (a) DM-electron scattering, (b) DM-atom scattering, and (c) DM-nucleon scattering. In the absence of heavy sea quarks and anti-quarks inside nucleons at the direct-detection energy scale, the $b$-quark-philic and $t$-quark-philic DM interacts with the constituent gluons via a loop, as shown in figure 7.

We compute the dominant DM-gluon one-loop elastic scattering amplitudes induced by the scalar, axial-vector, and twist-2 point interactions among DM and heavy quarks. The mass scale $m_{Q}$ of the heavy quarks running in the loop and the QCD coupling strength
α_s characterise the loop-amplitudes. We derive the phenomenological effective DM-gluon interaction Lagrangians given in (C.1a), (C.1b) and (C.1c) which correspond to Majorana, real scalar, and real vector DM candidates, respectively.

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**Figure 8**: Figures 8a, 8b and 8c depict the spin-independent b-quark-philic, t-quark-philic and gluon-philic Majorana DM-Nucleon scattering cross-sections respectively. The scalar and twist-2 type-1 contributions in all panels are evaluated using their respective lower bounds on the \( |C_{q,g}^{\chi,S,T}/\Lambda^2| \) satisfying \( \Omega \chi h^2 = 0.1199 \pm 0.0022 \) as shown in figures 1 and hence, regions above the solid curves are cosmologically allowed. Regions above the experimental limits obtained from XENON-1T [6] and PandaX-4T [8] are excluded.

Since the non-relativistic DM particles scatter the nucleons and not the free gluons, we perform the non-relativistic reduction of the interaction lagrangian given in the Appendix C. We connect the DM-gluon amplitudes induced by the scalar and twist-2 currents of heavy-quarks at one-loop order with their respective DM-effective nucleon interactions by evaluating the expectation values of the zero momentum scalar and twist-2 gluon-philic
operators between the initial and final nucleons in equations (C.3) and (C.5) respectively. The Majorana, real scalar, and real vector DM-nucleon scattering cross-sections driven by the heavy-quark scalar, axial-vector, and twist-2 currents are given as

$$
\sigma_{AV}^{qN} = \frac{2}{9 \pi} \left( \frac{C_{Vq}^q}{C_{Aq}^q} \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^6 \left[ \frac{m_N}{1 \text{ GeV}} \right]^2 \left[ \frac{\mu_{Nq}}{1 \text{ GeV}} \right]^2 \frac{4}{\alpha_s(\Lambda)}^2 \frac{4}{\alpha_s(\mu)} \frac{10 \text{ MeV}}{} \times |I_{AV}^{gg}|^2 |f_{Tq}^N|^2 (3.9 \times 10^{-46}) \text{ cm}^2
\tag{3.4a}
$$

$$
\sigma_{AV}^{qN} = \frac{16}{3 \pi} \left( C_{AV}^q \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^4 \left[ \frac{\mu_{Nq}}{1 \text{ GeV}} \right]^2 \frac{\alpha_s(\Lambda)}{\alpha_s(\mu)} \left[ \frac{175 \text{ GeV}}{m_Q} \right]^4 |I_{AV}^{gg}|^2 
\times \left[ \frac{\sum_{q=u,d,s} \Delta_{q}^{(N)} m_q}{m_q} \right]^2 \left[ \frac{\mathcal{J} + 1}{\mathcal{J}} \right] S_N S_N' (4.13 \times 10^{-57}) \text{ cm}^2
\tag{3.4b}
$$

$$
\sigma_{Tq}^{qN} = \frac{2}{\pi} \left( C_{Tq}^q \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^8 \frac{\alpha_s(\Lambda)}{4\pi} \frac{\mu_{Nq}}{1 \text{ GeV}} \left[ \frac{M_{Nq}}{100 \text{ GeV}} \right]^2 \frac{m_N}{1 \text{ GeV}} \times \left[ \ln \left( \frac{\Lambda^2}{m_Q} \right) \right]^2 |g(2; \Lambda)|^2 (1.98 \times 10^{-48}) \text{ cm}^2
\tag{3.4c}
$$

$$
\sigma_{Sq}^{qN} = \frac{4}{81 \pi} \left( C_{Sq}^q \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^4 \left[ \frac{100 \text{ GeV}}{m_{q^0}} \right]^2 \frac{m_N}{1 \text{ GeV}} \frac{\alpha_s(\Lambda)}{\alpha_s(\mu)} \left[ \frac{m_{q^0}}{100 \text{ GeV}} \right]^2 \frac{m_N}{1 \text{ GeV}} \times |I_{Sq}^{gg}|^2 |f_{Tq}^N|^2 (3.9 \times 10^{-44}) \text{ cm}^2
\tag{3.5a}
$$

$$
\sigma_{Tq}^{qN} = \frac{1}{2\pi} \left( C_{Tq}^q \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^8 \frac{\alpha_s(\Lambda)}{4\pi} \left[ \frac{\mu_{q^0}}{1 \text{ GeV}} \right]^2 \frac{M_{q^0}}{100 \text{ GeV}} \frac{m_N}{1 \text{ GeV}} \times \left[ \ln \left( \frac{\Lambda^2}{m_Q} \right) \right]^2 |g(2; \Lambda)|^2 (1.98 \times 10^{-48}) \text{ cm}^2
\tag{3.5b}
$$

$$
\sigma_{Sq}^{qN} = \frac{4}{243 \pi} \left( C_{Sq}^q \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^4 \left[ \frac{100 \text{ GeV}}{m_{q^0}} \right]^2 \frac{m_N}{1 \text{ GeV}} \frac{\alpha_s(\Lambda)}{\alpha_s(\mu)} \left[ \frac{m_{q^0}}{100 \text{ GeV}} \right]^2 \frac{m_N}{1 \text{ GeV}} \times |I_{Sq}^{gg}|^2 |f_{Tq}^N|^2 (3.9 \times 10^{-44}) \text{ cm}^2
\tag{3.6a}
$$

$$
\sigma_{AV}^{qN} = \frac{32}{9 \pi} \left( C_{AV}^q \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^4 \left[ \frac{\mu_{Nq}}{1 \text{ GeV}} \right]^2 \frac{\alpha_s(\Lambda)}{\alpha_s(\mu)} \left[ \frac{175 \text{ GeV}}{m_Q} \right]^4 |I_{AV}^{gg}|^2 
\times \left[ \frac{\sum_{q=u,d,s} \Delta_{q}^{(N)} m_q}{m_q} \right]^2 \left[ \frac{\mathcal{J} + 1}{\mathcal{J}} \right] S_N S_N' (4.13 \times 10^{-57}) \text{ cm}^2
\tag{3.6b}
$$
\[
\sigma_{T_2}^{\phi_{0N}} = \frac{1}{6\pi} \left( C_{T_2}^{\phi_{0N}} \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^8 \left[ \frac{\alpha_s(\Lambda)}{4\pi} \right]^2 \left[ \frac{\mu_{\phi_{0N}}}{1 \text{ GeV}} \right]^2 \left[ \frac{M_{\phi_{0N}}}{100 \text{ GeV}} \right]^2 \left[ \frac{m_N}{1 \text{ GeV}} \right]^2 \\
\times \left[ \ln \left( \frac{\Lambda^2}{m^2_{\phi_{0N}}} \right) \right]^2 |g(2; \Lambda)|^2 \times \left( 1.98 \times 10^{-48} \right) \text{ cm}^2
\]  
(3.6c)

where \( \mu_{\text{DM}} \equiv (m_N m_{\text{DM}}) / (m_N + m_{\text{DM}}) \) is the reduced mass for the respective DM-nucleon system. For the computation of scale dependent \( \alpha_s(\mu_R) \) and \( g(x; \mu_R) \), we access CTEQ6l1 [91] PDF data set from LHAPDF6 [92] library. However, the scale dependence of the Wilson coefficients is found to be smaller than that of \( \alpha_s \), as noted by the authors of the reference [93]. It is to be noted that the observed logarithmically enhanced one-loop induced scattering cross-sections in equations (3.4c), (3.5b) and (3.6c) result from the explicit momentum dependence in the twist-2 operator interaction Lagrangians (2.3), (2.5) and (2.6) respectively. The real vector DM-nucleon scattering cross-sections driven by the scalar and twist-2 type-2 currents of heavy quarks in equations (3.6a) and (3.6c) respectively are found to be 1/3 of the scalar DM-nucleon scattering cross-sections corresponding to scalar and twist-2 type-2 currents of heavy-quarks in equations (3.5a) and (3.5b) respectively.

For completion, we compute and display the tree-level spin-independent Majorana DM-nucleon, scalar DM-nucleon, and vector DM-nucleon scattering cross-sections induced by the scalar and twist-2 current of gluons in equations (3.7a)-(3.7b), (3.8a)-(3.8b) and (3.9a)-(3.9b) respectively as

\[
\sigma_{S}^{\phi_{0N}} = \frac{128}{81 \pi} \left( C_{S}^{\phi_{0N}} \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^8 \left[ \frac{m_X}{100 \text{ GeV}} \right]^2 \left[ \frac{m_N}{1 \text{ GeV}} \right]^2 \left[ \frac{\mu_{N_X}}{1 \text{ GeV}} \right]^2 \\
\times \left[ \frac{\alpha_s(\Lambda)}{\alpha_s(\mu)} \right]^2 |f_{T_1}^{\phi_{0N}}|^2 \times \left( 3.9 \times 10^{-48} \right) \text{ cm}^2
\]  
(3.7a)

\[
\sigma_{T_1}^{\phi_{0N}} = \frac{9}{8 \pi} \left( C_{T_1}^{\phi_{0N}} \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^8 \left[ \frac{m_X}{100 \text{ GeV}} \right]^2 \left[ \frac{m_N}{1 \text{ GeV}} \right]^2 \left[ \frac{\mu_{N_X}}{1 \text{ GeV}} \right]^2 \\
\times |g(2; \mu_F)|^2 \times \left( 3.9 \times 10^{-48} \right) \text{ cm}^2
\]  
(3.7b)

\[
\sigma_{S}^{\phi_{0N}} = \frac{64}{81 \pi} \left( C_{S}^{\phi_{0N}} \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^4 \left[ \frac{100 \text{ GeV}}{m_{\phi_{0N}}} \right]^2 \left[ \frac{m_N}{1 \text{ GeV}} \right]^2 \left[ \frac{\mu_{N_{\phi_{0N}}}}{1 \text{ GeV}} \right]^2 \left[ \frac{\alpha_s(\Lambda)}{\alpha_s(\mu)} \right]^2 \\
\times |f_{T_1}^{\phi_{0N}}|^2 \left( 3.9 \times 10^{-44} \right) \text{ cm}^2
\]  
(3.8a)

\[
\sigma_{T_2}^{\phi_{0N}} = \frac{9}{32 \pi} \left( C_{T_2}^{\phi_{0N}} \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^8 \left[ \frac{\alpha_s(\Lambda)}{4\pi} \right]^2 \left[ \frac{\mu_{N_{\phi_{0N}}}}{1 \text{ GeV}} \right]^2 \left[ \frac{m_{\phi_{0N}}}{100 \text{ GeV}} \right]^2 \left[ \frac{m_N}{1 \text{ GeV}} \right]^2 \\
\times |g(2; \mu_F)|^2 \left( 3.9 \times 10^{-48} \right) \text{ cm}^2
\]  
(3.8b)
Figure 9: Figures 9a, 9b and 9c depict the spin-independent $b$-quark-philic, $t$-quark-philic and gluon-philic scalar DM-Nucleon scattering cross-sections respectively. The scalar and twist-2 type-2 contributions in all the panels are evaluated using their respective lower bounds on $|C_{q,g}^{bT_2}|/\Lambda^n$ satisfying $\Omega^{bT_2}h^2 = 0.1199 \pm 0.0022$ \cite{4} as shown in figures 2 and hence, regions above the solid curves are cosmologically allowed. Regions above the experimental limits obtained from XENON-1T \cite{6} and PandaX-4T\cite{8} are excluded.

\[
\sigma_{S}^{g^{1-0N}} = \frac{64}{243} \frac{1}{\pi} \left( C_{V_0}^{qT_2} \right)^2 \left[ \frac{1 \text{ TeV}}{\Lambda} \right]^4 \left[ \frac{100 \text{ GeV}}{m_{V^0}} \right]^2 \left[ \frac{m_N}{1 \text{ GeV}} \right]^2 \left[ \frac{\mu_{N^0V^0}}{1 \text{ GeV}} \right]^2 \left[ \frac{\alpha_s(\Lambda)}{\alpha_s(\mu)} \right]^2 \times |f_N^V|^2 \left( 3.9 \times 10^{-44} \right) \text{ cm}^2
\]  

(3.9a)
Figures 10a, 10a and 10c depict the spin-independent $b$-quark-philic, $t$-quark-philic and gluon-philic vector DM-Nucleon scattering cross-sections respectively. The scalar and twist-2 type-2 contributions in all the panels are evaluated using their respective lower bounds on $\left| |C_{q,g}^{V_0 S,T_2}| \right|$ satisfying $\Omega V_0 h^2 = 0.1199 \pm 0.0022$ [4] as shown in figures 3 and hence, regions above the solid curves are cosmologically allowed. Regions above the experimental limits obtained from XENON-1T [6] and PandaX-4T [8] are excluded.

The analytical expressions for the gluon-philic DM-nucleon scattering cross-sections are in agreement with those given in reference [72].

We display the Majorana, real scalar and real vector DM-Nucleon scattering cross-sections w.r.t. DM mass in figures 8, 9 and 10 respectively. The scalar and twist-2 currents
induced by $b$-quark-philic and $t$-quark-philic DM interactions are depicted in the left and right panels of all the three figures. Since, the cross-sections are evaluated using the lower bound on the Wilson coefficients obtained from the relic density contours satisfying $\Omega_{DM}h^2$ in figures 1, 2 and 3 for the Majorana, scalar and vector DM respectively, the solid curves represent the cosmological lower-limits of the scattering cross-sections and therefore the regions above the curves shall be considered for the validation from the XENON-1T [6] and PandaX-4T [8].

The spin-dependent DM-nucleon cross-section is suppressed by the fourth power of the momentum exchanged $\propto v_{DM}^4$ as a result of the DM-gluon effective pseudo-scalar Lagrangian generated by the axial-vector coupling of Majorana/Vector DM to heavy quarks. The scattering cross-sections given in equations (3.4b) and (3.6b) however, are compared with the upper-limits on the spin-dependent scattering cross-sections [10]. The contribution of velocity independent spin dependent DM-nucleon scattering cross-sections can be validated from IceCube and PICO experiments using the $b \bar{b}$ channel annihilation data as depicted in figure 2 of reference [94]. We find that the cosmologically constrained Majorana, real scalar and real vector DM candidates interacting through axial-vector currents of third generation heavy quarks are allowed by the available data for spin dependent scattering cross-sections.

### 3.3.1 Nuclear Recoil Spectrum

For a fixed target detector exposure $\epsilon_T$ and target nucleus mass $m_{nuc.}$, the differential nuclear recoil event rate w.r.t. nuclear recoil energy $dR_{nuc.}/dE_r$ corresponding to the DM-Nucleon scattering cross-section $\sigma_N$ is given as

$$\frac{dR_{nuc.}}{dE_r} = \frac{\epsilon_T \rho_0}{m_{nuc.} m_{DM}} \int_{v_{\text{min}}}^{v_{\text{esc}}} v f(v) \frac{d\sigma}{dE_r} dv = \frac{\epsilon_T \rho_{DM}}{2 m_{DM}} \frac{\sigma_N A^2}{\mu_N^2} |F(|q| r_n)|^2 \int_{v_{\text{min}}}^{v_{\text{max}}} \frac{1}{v} f(v) d^3 v$$

(3.10)

where $F(|q| r_n)$ is the Helm form factor taking account of the non-vanishing finite size of the nucleus [95]. Here $r_n = 1.2 \times A^{1/3}$ is the effective radius of the nucleus with atomic mass $A$ and $|q|$ is the momentum transfer corresponding to the recoil energy $E_r$. Following reference [96], the velocity integral for normalized Maxwellian DM velocity distribution is solved as

$$\int_{v_{\text{min}}}^{v_{\text{max}}} \frac{f(v)}{v^2} d^3 v = \int_{v_{\text{min}}}^{v_{\text{esc}}} \sqrt{m_{nuc.} E_r/(2 \mu_{nuc.}^2)} \frac{1}{v} e^{-(v+v_E)/v_0} d^3 v = \frac{1}{2 v_E} \text{erf}(\eta_-, \eta_+) - \frac{1}{\pi v_E} (\eta_+ - \eta_-) e^{-[v_{\text{esc}}/v_0]^2}$$

where

$$\eta_{\pm} = \min \left[ \sqrt{m_{nuc.} E_r/(2 \mu_{nuc.}^2)} \pm v_E, \frac{v_{\text{esc}}}{v_0} \right]$$

and $\mu_{nuc.} \equiv \frac{m_{DM} m_{nuc.}}{m_{DM} + m_{nuc.}}$.

(3.11a)

For numerical computation we have taken Earth’s velocity relative to galactic frame to be $v_E = 232 \text{ km/sec}$, $v_0 = 220 \text{ km/s}$ and the escape velocity $v_{\text{esc}} = 544 \text{ km/s}$. Further, in order to incorporate the detector based effects, the nucleus recoil event rate $dR_{nuc.}/dE_r$
is convoluted with the detector efficiency [6]. Integrating over the recoil energy from $E_{\text{th}} \sim 4.9$ KeV to the maximum $E_{\text{max}}$ for a fixed duration and size of the detector, we can estimate the expected total number of recoiled nucleus events.

Figure 11: Figures 11a, 11b and 11c depict the spin-independent recoiled nucleus event due to Majorana DM - Xe nucleus, scalar DM - Xe nucleus and vector DM - Xe nucleus scatterings respectively. The scalar and twist-2 interactions of $b$-quark, $t$-quark and gluon in all the three panels are evaluated using their respective lower bounds on $\left| C_{q,g}^{S,T} / \Lambda_n \right|$, $\left| C_{\phi, S,T}^{q,g} / \Lambda_n \right|$ and $\left| C_{V,S,T}^{q,g} / \Lambda_n \right|$ satisfying $\Omega_{\text{DM}} h^2 = 0.1199 \pm 0.0022$ [4] as shown in figures 1, 2 and 3 respectively and hence, regions above the solid curves are cosmologically allowed. Shaded region is excluded due to non-observation of at least one-event in the XENON-1T experiment.

As an illustration, we plot the probable number of Xe nuclear recoil events w.r.t. varying DM mass in a XENON-1T [6] setup where 1.3 tonnes of Xe target are exposed for
a duration of 278.8 days, which is equivalent to one Ton-year of net target exposure. The expected number of recoiled nucleus events due to Majorana DM-Xe nucleus scattering, scalar DM-Xe nucleus scattering and vector DM-Xe nucleus scattering, respectively, are depicted in figures 11a, 11b and 11c, which are induced by scalar and twist-2 currents of heavy quarks at one loop level and of gluons at tree level. Since the solid curves in both figures correspond to the lower-limits of the expected number of recoiled nucleus events for a given DM mass, the region above the respective curves is cosmologically allowed.

The absence of at least one event in the XENON-1T experiment rules out the contributions of $b$-quark-philic and gluon-philic Majorana DM scalar interactions for $m_\chi \leq 1.8$ and 1.3 TeV, respectively, as shown in figure 11a. Barring the said two operators, we find that contributions from all other operators pertaining to Majorana, scalar and vector DM for $m_{DM} \geq 200$ GeV may be probed in the ongoing and future direct-detection experiments with enhanced target exposure.

4 Summary and conclusions

Figure 12: Figures 12a and 12b corresponding to scalar and twist-2 type-1 operators induced interactions depict the allowed range of the Wilson coefficients in red, blue and green for the $b$-quark-philic, $t$-quark-philic and gluon-philic Majorana DM respectively. The shaded region enclosed by the lower-limit and upper-limit contours of Wilson coefficients $|C_{\chi S,T,1}^{n,g}/\Lambda^n|$ satisfy the relic density constraint $\Omega h^2 = 0.1199 \pm 0.0022$ [4] and non-observation of at least one-event in XENON-1T experiment [6] respectively.

In this article we assess the viability of the $b$-quark-philic, $t$-quark-philic, and gluon-philic self-conjugated spin 1/2, 0 and 1 DM candidates in the EFT approach. We have formulated the generalised effective interaction Lagrangian induced by the scalar, axial-vector, and twist-2 operators for the real particles in section 2. The contribution from the vector operator vanishes for the real particles.
Figure 13: Figures 13a and 13b corresponding to scalar and twist-2 type-2 operators induced interactions depict the allowed range of the Wilson coefficients in red, blue and green for the $b$-quark-philic, $t$-quark-philic and gluon-philic scalar DM respectively. The shaded region enclosed by the lower-limit and upper-limit contours of Wilson coefficients $\left| \frac{C_{q,g}}{\Lambda^n} \right|$ satisfy the relic density constraint $\Omega_\phi^0 h^2 = 0.1199 \pm 0.0022$ [4] and the non-observation of at least one-event in XENON-1T experiment [6] respectively.

For a given DM mass, the relic abundance of Majorana, scalar, and vector DM is computed in section 3.1 using the thermally averaged cross-sections in appendix B. Figures 1, 2 and 3 show twenty-two interaction strengths in the form of Wilson coefficients $\left| \frac{C_{q,g}}{\Lambda^n} \right|$ in TeV$^{-n}$ (eight for Majorana DM, six for scalar DM and eight for vector DM) are then constrained from the PLANCK’s observation $\Omega_{\text{DM}} h^2 \approx 0.1198$ [4].

Using the constrained Wilson coefficients, we study the thermal averaged DM pair annihilation cross-sections for the indirect detection of varying DM masses ($10^{-1} - 2$ TeV) in figures 4, 5 and 6. The contributions to the lower limits on the annihilation cross-sections to $b\bar{b}$, $t\bar{t}$ and $gg$ channels driven by the effective couplings are found to be consistent when compared with the upper limit on the $b\bar{b}$ annihilation cross-section obtained from FermiLAT [15] and H.E.S.S. [17] indirect experiments. The predicted annihilation cross-sections for all the twenty-two operators are found to be consistent with the available data from indirect-detection experiments.

The scattering of the incident heavy-quark-philic DM particle off the static nucleon induced by the effective one loop-interactions of the Majorana / scalar / vector DM with gluons in the direct-detection experiments are studied. The contributions of gluon-philic Majorana, scalar and vector DM are revisited and found to be in agreement with results in the literature [72]. The lower limits of the scattering cross-sections for the dominant spin-independent processes due to scalar and twist currents for Majorana, scalar DM and vector DM are shown in figures 8, 9 and 10 respectively and are compared with the available results.
Figures 14a and 14b corresponding to scalar and twist-2 type-2 operators induced interactions depict the allowed range of the Wilson coefficients in red, blue and green for the $b$-quark-philic, $t$-quark-philic and gluon-philic vector DM respectively. The shaded region enclosed by the lower-limit and upper-limit contours of Wilson coefficients $|C^{q,g}_{V_S,T_2}|/\Lambda$ satisfy the relic density constraint $\Omega^{V_0}h^2 = 0.1199 \pm 0.0022$ [4] and the non-observation of at least one-event in XENON-1T experiment [6] respectively.

from XENON-1T [6] and PandaX-4T [8]. Further, using the lower limits on the Wilson coefficients, the expected number of events scattered Majorana, scalar and vector DM in the XENON-1T experiment for an equivalent target exposure of 1 Ton-year are shown in figures 11a, 11b and 11c, respectively.

Finally, figures 12, 13 and 14 corresponding to real spin 1/2, 0 and 1 DM candidates respectively encapsulate the predicted range for the eighteen Wilson coefficients associated with scalar and twist-2 operators where the cosmological relic density [4] and the non-observation of at least one-event in XENON-1T direct-detection experiment [6] give the lower and upper bounds for varying DM mass $\sim 0.01 \leq m_{DM} \leq 2$ TeV respectively. Findings of our analysis are summarised as follows:

- Contributions from $b$-quark-philic, $t$-quark-philic, and gluon-philic Majorana DM scalar operators are found to be viable solutions for $m_\chi \geq 1.8$ TeV, 200 GeV, and 1.25 TeV, respectively, as shown in figure 12a. The contributions of $b$-quark-philic, $t$-quark-philic, and gluon-philic twist-2 type-1 Majorana DM operators shown in figure 12b have been validated for a much wider spectrum of $m_\chi \geq 60$ GeV, 200 GeV, and 200 GeV, respectively.

- In figure 13a contributions from $b$-quark-philic, $t$-quark-philic and gluon-philic scalar DM scalar operators are found to be consistent with current experimental data for $m_\phi \geq 350$ GeV, 200 GeV and 350 GeV, respectively. The corresponding contributions
induced by twist-2 type-2 DM operators are validated for $m_{\phi^0} \geq 150$ GeV, 200 GeV and 280 GeV, respectively as shown in the figure 13b.

- In figure 14a contributions from $b$-quark-philic, $t$-quark-philic and gluon-philic vector DM scalar operators are validated for $m_{V^0} \geq 350$ GeV, 180 GeV and 350 GeV, respectively. The corresponding contributions induced by twist-2 type-2 DM operators are validated for $m_{V^0} \geq 100$ GeV, 180 GeV and 200 GeV, respectively as shown in figure 14b.

This analysis shows that the third-generation heavy-quark-philic and gluon-philic real DM are promising for $m_{\text{DM}} \geq 200$ GeV and are likely to be probed and falsified otherwise in the ongoing and future upgraded direct detection experiments. This study also opens up the scope for future collider based DM search analysis through additional channels induced by the cosmologically constrained operators.

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Appendix

A DM Pair Annihilation Cross-section

The annihilation cross-sections of a pair of Majorana DM $\chi$ to a pair of third generation heavy quarks induced by the scalar, axial-vector and twist-2 type-1 operators are given as:

$$\sigma_S(\chi \bar{\chi} \rightarrow ff) = \left[ \frac{C_f^S m_f}{\Lambda^3} \right]^2 \frac{C_a}{16 \pi} \beta_f^3 \beta_f s$$ (A.1a)

$$\sigma_{AV}(\chi \bar{\chi} \rightarrow ff) = \left[ \frac{C_f^{AV}}{\Lambda^2} \right]^2 \frac{C_a}{12 \pi} \beta_f \beta_f s (1 - 4(x_f + x_f) + 28x_f x_f)$$ (A.1b)

$$\sigma_{T_1}(\chi \bar{\chi} \rightarrow ff) = \left[ \frac{C_f^{T_1}}{\Lambda^4} \right]^2 \frac{C_a}{960 \pi} \beta_f^3 s^3 \left[ 4x_f^2 (92x_f^2 + 9x_f - 8) + 3x_f (12x_f^2 + 9x_f + 2) - 32x_f^2 + 6x_f + 8 \right]$$ (A.1c)

where $C_a = 3$, $x_i = \frac{m_i^2}{s}$ and $\beta_i = \sqrt{1 - 4x_i}$.

As mentioned in the text, we include the study of gluon-philic Majorana DM interactions induced by the scalar and twist-2 type-1 operators. The annihilation cross-sections of
the Majorana DM pair to a pair of gluons are given as

\[
\sigma_S (\chi \bar{\chi} \to gg) = \left[ \frac{C_{gS}^2}{\Lambda^4} \right] \frac{C_a C_f}{4 \pi} \left( \frac{\alpha_s}{\pi} \right)^2 \beta_\chi s^2 M^2 \tag{A.2a}
\]

\[
\sigma_{T_1} (\chi \bar{\chi} \to gg) = \left[ \frac{C_{gS T_1}^2}{\Lambda^4} \right] \frac{C_a C_f}{80 \pi} \beta_\chi s^3 \left( 1 + \frac{8}{3} x_\chi \right) \tag{A.2b}
\]

where \( C_f = 4/3 \).

The annihilation cross-sections of a pair of real spin 0 DM \( \phi^0 \) to a pair of third generation heavy quarks induced by the scalar and twist-2 type-2 operators are given as:

\[
\sigma_S (\phi^0 \phi^0 \to f \bar{f}) = \left[ \frac{C_{\phi_0 S}^2}{\Lambda^4} \right] \frac{C_a \beta^3}{2 \pi \beta_\phi} \tag{A.3a}
\]

\[
\sigma_{T_2} (\phi^0 \phi^0 \to f \bar{f}) = \left[ \frac{C_{\phi_0 T_2}^2}{\Lambda^4} \right] \frac{C_a \beta^3}{240 \pi \beta_\phi} s^3 \left[ 2x_\phi^2 (23x_f + 8) - 4x_\phi (7x_f + 2) + 6x_f + 1 \right] \tag{A.3b}
\]

The cross-sections for the pair of scalar DM annihilation into a pair of gluons induced by scalar and twist-2 type-2 operators are given as

\[
\sigma_S (\phi^0 \phi^0 \to gg) = \left[ \frac{C_{\phi_0 S}^2}{\Lambda^4} \right] \frac{C_a \beta^3}{2 \pi \beta_\phi} \frac{s}{\beta_\phi} \tag{A.4a}
\]

\[
\sigma_{T_2} (\phi^0 \phi^0 \to gg) = \left[ \frac{C_{\phi_0 T_2}^2}{\Lambda^4} \right] \frac{C_a C_f}{60 \pi} s^3 \beta_\phi^3 \tag{A.4b}
\]

Similarly, the production of a pair of real vectors DM \( V^0 \) as a result of the annihilation of a pair of third generation heavy quarks induced by the scalar, axial-vector, and twist-2 type-2 operators is given as

\[
\sigma_S (V^0 V^0 \to f \bar{f}) = \left[ \frac{C_{V_0 S}^2}{\Lambda^4} \right] \frac{C_a \beta^3}{72 \pi \beta_{V^0}} \frac{1 - 4x_{V^0} + 12x_{V^0}^2}{x_{V^0}^2} \tag{A.5a}
\]

\[
\sigma_{AV} (V^0 V^0 \to f \bar{f}) = \left[ \frac{C_{V_0 AV}^2}{\Lambda^4} \right] \frac{C_a \beta_\phi \beta_{V^0}}{27 \pi} \frac{s}{\beta_{V^0}} \left( 1 - 4 (x_{V^0} + x_f) + 28 x_{V^0} x_f \right) \tag{A.5b}
\]

\[
\sigma_{T_2} (V^0 V^0 \to f \bar{f}) = \left[ \frac{C_{V_0 T_2}^2}{\Lambda^4} \right] \frac{C_a \beta^3}{8640 \pi} \frac{\beta_\phi^3 \beta_{V^0}^3}{s^3} \frac{1 - 4x_{V^0} + 12x_{V^0}^2}{x_{V^0}^2} \tag{A.5c}
\]

The annihilation cross-sections of a pair of vector DM induced by the scalar and twist-2
type-2 gluon currents are given as

\[
\sigma_S(V^0 V^0 \rightarrow g g) = \left[ \frac{C_{tg}^g}{\Lambda^2} \right]^2 \frac{C_a C_f}{18 \pi} \left( \frac{\alpha_s}{\pi} \right)^2 \frac{s}{\beta_{V^0}^2} \frac{1 - 4x_{V^0} + 12x_{V^0}^2}{x_{V^0}^2} \quad (A.6a)
\]

\[
\sigma_{T_2}(V^0 V^0 \rightarrow g g) = \left[ \frac{C_{tg}^g}{\Lambda^2} \right]^2 \frac{C_a C_f}{2160 \pi} s \beta_{V^0}^3 \frac{(1 - 4x_{V^0} + 12x_{V^0}^2)}{x_{V^0}^2} \quad (A.6b)
\]

**B Thermally-averaged annihilation cross-section**

In order to compute the probability of a DM particle being annihilated by another one, the aforementioned annihilation cross-sections are re-written in terms of the magnitude of relative velocity \( |\vec{v}_{DM_1} - \vec{v}_{DM_2}| \equiv |\vec{v}| \) and the dimensionless ratio \( \xi_f \equiv m_f^2 / m_{DM}^2 \).

The thermal average of the heavy-quark-philic Majorana DM pair annihilation cross-sections given in equations (A.1a), (A.1b) and (A.1a) corresponding to scalar, axial-vector and twist-2 operators respectively are given as

\[
\langle \sigma_S(\chi \chi \rightarrow f \bar{f}) | | \vec{v} \rangle = \left[ \frac{C_{\chi|s}^f}{\Lambda^3} \right]^2 \frac{C_a}{8 \pi} \frac{m_f^2}{m_{\chi}} \frac{m_f^2}{m_{\chi}} (1 - \xi_f)^{3/2} \quad (B.1a)
\]

\[
\langle \sigma_{AV}(\chi \chi \rightarrow f \bar{f}) | | \vec{v} \rangle = \left[ \frac{C_{\chi AV}^f}{\Lambda^2} \right]^2 \frac{C_a}{2 \pi} \sqrt{1 - \xi_f} \left[ 1 + \frac{|\vec{v}_\chi|^2 (8 \xi_f^{-1} - 28 + 23 \xi_f)}{1 - \xi_f} \right] \quad (B.1b)
\]

\[
\langle \sigma_{T_1}(\chi \chi \rightarrow f \bar{f}) | | \vec{v} \rangle = \left[ \frac{C_{\chi T_1}^f}{\Lambda^4} \right]^2 \frac{C_a}{2 \pi} \frac{m_f^6}{m_{\chi}} \sqrt{1 - \xi_f} \left[ 1 + \frac{|\vec{v}_\chi|^2 (56 - 41 \xi_f - 8 \xi_f^2 + 11 \xi_f^3)}{48 (1 - \xi_f) (2 + \xi_f)} \right] \quad (B.1c)
\]

Using annihilation cross-sections in (A.2a) and (A.2b) the thermal averaged annihilation cross-sections for the gluon-philic Majorana DM are given as

\[
\langle \sigma_S(\chi \chi \rightarrow g g) | | \vec{v} \rangle = \left[ \frac{C_{\chi g}^g}{\Lambda^2} \right]^2 \frac{C_a}{\pi} \frac{2 C_a}{\pi} \frac{m_f^6}{m_{\chi}} \quad (B.2a)
\]

\[
\langle \sigma_{T_2}(\chi \chi \rightarrow g g) | | \vec{v} \rangle = \left[ \frac{C_{\chi T_2}^g}{\Lambda^4} \right]^2 \frac{2 C_a}{3 \pi} \frac{m_f^6}{m_{\chi}} \quad (B.2b)
\]

The thermal average of heavy-quark-philic scalar DM pair annihilation cross-sections displayd in equations (A.3a) and (A.3b) are given as

\[
\langle \sigma_S(\phi^0 \phi^0 \rightarrow f \bar{f}) | | \vec{v} \rangle = \left[ \frac{C_{\phi_2}^f}{\Lambda^2} \right]^2 \frac{C_a}{\pi} \frac{(1 - \xi_f)^{3/2}}{8} \left[ 1 + \frac{|\vec{v}_{\phi}|^2 (-2 + 5 \xi_f)}{(1 - \xi_f)} \right] \quad (B.3a)
\]

\[
\langle \sigma_{T_2}(\phi^0 \phi^0 \rightarrow f \bar{f}) | | \vec{v} \rangle = \left[ \frac{C_{\phi_2}^f}{\Lambda^4} \right]^2 \frac{C_a}{4 \pi} \frac{m_f^2 m_{\phi}^2}{m_{\phi}} \frac{m_f^6}{m_{\chi}} \left[ 1 + \frac{|\vec{v}_{\phi}|^2 (10 - \xi_f)}{24 (1 - \xi_f)} \right] \quad (B.3b)
\]
Similarly, the thermal average of heavy-quark-philic scalar DM pair annihilation cross-sections displayed in equations (A.4a) and (A.4b) are given as

\[
\langle \sigma_S (\phi^0 \phi^0 \rightarrow gg) | \vec{v}| \rangle = \left[ \frac{C_g^{\phi^0}}{\Lambda^2} \right]^2 \left( \frac{\alpha_s}{\pi} \right)^2 \frac{16 C_a C_f m_{\phi^0}^2}{\pi} \quad (B.4a)
\]

\[
\langle \sigma_{T_2} (\phi^0 \phi^0 \rightarrow gg) | \vec{v}| \rangle = \left[ \frac{C_g^{\phi^0}}{\Lambda^4} \right]^2 \frac{2 C_a C_f m_{\phi^0}^6 |\vec{v}_{\phi^0}|^4}{15 \pi} \quad (B.4b)
\]

The thermal average of the vector DM pair annihilation cross-sections given in equations (A.5a), (A.5b) and (A.5c) corresponding to the scalar, axial-vector and twist-2 type-2 operators, respectively, are given as

\[
\langle \sigma_S (V^0 V^0 \rightarrow f \bar{f}) | \vec{v}| \rangle = \left[ \frac{C_f V^0}{\Lambda^2} \right]^2 \frac{C_a}{3 \pi} (1 - \xi_f)^{3/2} \left[ 1 + \left| \vec{v}_{V^0} \right|^2 \frac{(2 + 7 \xi_f)}{24} \right] \quad (B.5a)
\]

\[
\langle \sigma_{AV} (V^0 V^0 \rightarrow f \bar{f}) | \vec{v}| \rangle = \left[ \frac{C_f V^0}{\Lambda^2} \right]^2 \frac{2 C_a m_f^2 \left| \vec{v}_{V^0} \right|^2}{9 \pi} \sqrt{1 - \xi_f} \quad (B.5b)
\]

\[
\langle \sigma_{T_2} (V^0 V^0 \rightarrow f \bar{f}) | \vec{v}| \rangle = \left[ \frac{C_f V^0}{\Lambda^4} \right]^2 \frac{C_a \left| \vec{v}_{V^0} \right|^4 m_{V^0}^6}{90 \pi} \left( 1 - \xi_f \right)^{3/2} \left[ 1 + \frac{3}{2} \xi_f \right] \quad (B.5c)
\]

Similarly, the thermal average of the vector DM pair annihilation to gluon channels as displayed in equations (A.6a) and (A.6b) corresponding to the scalar and twist-2 currents respectively are given as

\[
\langle \sigma_S (V^0 V^0 \rightarrow gg) | \vec{v}| \rangle = \left[ \frac{C_g V^0}{\Lambda^2} \right]^2 \left( \frac{\alpha_s}{\pi} \right)^2 \frac{16 m_{V^0}^2 C_a C_f}{3 \pi} \left[ 1 + \frac{\left| \vec{v}_{V^0} \right|^2}{3} \right] \quad (B.6a)
\]

\[
\langle \sigma_{T_2} (V^0 V^0 \rightarrow gg) | \vec{v}| \rangle = \left[ \frac{C_g V^0}{\Lambda^4} \right]^2 \frac{2 C_a C_f m_{V^0}^6 \left| \vec{v}_{V^0} \right|^4}{45 \pi} \quad (B.6b)
\]

C Effective DM-Nucleon Interactions from 1-Loop DM-Gluon amplitudes

The DM-gluon scattering arising due to the one loop Feynman-diagrams in figure 7 is induced by the scalar, axial-vector and twist-2 currents of the heavy quarks. These one-loop amplitudes characterise the effective point interaction Lagrangian for the gluon with
Majorana, real scalar and real vector DM candidates respectively and are given as

\[
\mathcal{L}_{\text{eff.}}^{\chi gg} = \frac{C_g^{\chi g}}{\Lambda^3} \frac{\alpha_s}{4\pi} (\bar{\chi} \chi) (G^a_{\alpha \beta}) (G^a_{\alpha \beta}) I^g_S + \frac{C_q^{\chi q}}{\Lambda^2} \frac{\alpha_s}{4\pi} m_{\chi} (\bar{\chi} i \gamma^5 \chi) (\bar{G}^a_{\alpha \beta}) (G^a_{\alpha \beta}) 4 I^g_A
\]

\[
\mathcal{L}_{\text{eff.}}^{\phi^0 \phi^0 gg} = \frac{C_g^{\phi^0 g}}{\Lambda^2} \frac{\alpha_s}{4\pi} (\phi^0 \phi^0) (G^a_{\alpha \beta}) (G^a_{\alpha \beta}) I^g_S
\]

\[
\mathcal{L}_{\text{eff.}}^{V^0 V^0 gg} = \frac{C_g^{V^0 g}}{\Lambda^2} \frac{\alpha_s}{4\pi} (V^0)^{\mu \nu} (V^0)^{\mu \nu} (G^a_{\alpha \beta}) (G^a_{\alpha \beta}) I^g_S
\]

The dimensionless one-loop integrals \( I^g_S \), \( I^g_A \), and \( I^g_T \) are defined as

\[
I^g_S = \int_0^1 dx \int_0^{1-x} dy \frac{1 - 4 x y}{1 - \frac{q^2}{m_Q^2} x y}
\]

\[
I^g_A = \int_0^1 dx \int_0^{1-x} dy \frac{x^2 - x - x y}{1 - \frac{q^2}{m_Q^2} x y}
\]

where the square of the four momentum transferred \( q^2 \approx 2 m_N E_r \) (\( E_r \lesssim 100 \text{ KeV} \) is the recoil energy of the nucleon). \( m_Q \) and \( m_N \) are the concerned heavy quark mass (\( m_b/m_t \)) running in the loop and nucleon mass respectively.

We observe that the one-loop effective contact interactions generated from quark scalar, twist and axial vector currents contain the scalar, pseudoscalar and twist gluon operators respectively. We perform the non-relativistic reduction of these operators for zero momentum partonic gluons and evaluate the gluonic operators between the nucleon states. The zero momentum gluon contribution to the hadronic matrix element \( f_{T G}^N \approx .923 \) is extracted in terms of light quark as [97–99]

\[
f_{T G}^N \equiv - \frac{1}{m_N} \left\langle N \left| \frac{9}{8} \frac{\alpha_s}{\pi} G_{\alpha \beta} G_{\alpha \beta} \right| N \right\rangle = 1 - \sum_{q=u, d, s} f_{T q}^N \equiv 1 - \sum_{q=u, d, s} \left\langle N \left| \frac{m_Q}{m_N} (\bar{q} q) \right| N \right\rangle
\]

According to [99] and [100], the pseudoscalar gluon operator between nucleon states is computed as

\[
\left\langle N \left| G_{\alpha \beta} \bar{G}_a^{\alpha \beta} \right| N \right\rangle = m_N \bar{m} \sum_{q=u, d, s} \frac{1}{m_q} \Delta^{(N)}_q \quad \text{where } \bar{m} = \left[ \frac{1}{m_u} + \frac{1}{m_d} + \frac{1}{m_s} \right]^{-1}
\]
The axial vector current \( \langle N \mid \bar{q} \gamma^\mu \gamma^5 q \mid N \rangle = 2\Delta_q^{(N)} s^\mu \) specifies the coefficient \( \Delta_q^{(N)} \) as the spin content of the nucleon’s quark \( q \). The coefficients for light quarks satisfy \( \Delta_u^{(p)} = \Delta_d^{(n)} \), \( \Delta_u^{(n)} = \Delta_d^{(p)} \) and \( \Delta_s^{(p)} = \Delta_s^{(n)} \), while the contribution from heavy quarks is found to be vanishingly small [101]. It is important to observe that a quark axial-vector current is related to the gluonic pseudoscalar current by PCAC [100]. The zero momentum nucleonic matrix element for gluon twist-2 operators is defined as

\[
\langle N \mid O_{\mu\nu}^g \mid N \rangle = -\frac{1}{m_N} \left[ k^\mu k^\nu - \frac{1}{4} g^{\mu\nu} m_N^2 \right] g(2; \mu_R) ;
\]

where \( g(2; \mu_R) = \int_0^1 x g(x; \mu_R) \, dx \) \( (C.5) \)

The \( \alpha_s(\mu_R) \) and gluon PDF \( g(x; \mu_R) \) are evaluated at the renormalisation scale \( \mu_R = \Lambda \) [68].

References

[1] Vera C. Rubin, W. Kent Ford, Jr., Astrophysical Journal, vol. 159, p.379 (1970), doi:10.1086/150317
[2] D. Clowe, M. Bradac, A. H. Gonzalez, M. Markevitch, S. W. Randall, C. Jones and D. Zaritsky, Astrophys. J. Lett. 648, L109-L113 (2006) doi:10.1086/508162 [arXiv:astro-ph/0608407 [astro-ph]].
[3] E. Komatsu et al. [WMAP Science Team], PTEP 2014, 06B102 (2014) doi:10.1093/ptep/ptu083 [arXiv:1404.5415 [astro-ph.CO]].
[4] N. Aghanim et al. [Planck], Astron. Astrophys. 641 (2020), A6 doi:10.1051/0004-6361/201833910 [arXiv:1807.06209 [astro-ph.CO]].
[5] D. S. Akerib et al. [LUX], Nucl. Instrum. Meth. A 704 (2013), 111-126 doi:10.1016/j.nima.2012.11.135 [arXiv:1211.3788 [physics.ins-det]].
[6] E. Aprile et al. [XENON], Phys. Rev. Lett. 121, no.11, 111302 (2018) doi:10.1103/PhysRevLett.121.111302 [arXiv:1805.12562 [astro-ph.CO]].
[7] P. Agnes et al. [DarkSide], Phys. Rev. Lett. 121, no.8, 081307 (2018) doi:10.1103/PhysRevLett.121.081307 [arXiv:1802.06994 [astro-ph.HE]].
[8] Y. Meng et al. [PandaX-4T], [arXiv:2107.13438 [hep-ex]].
[9] A. H. Abdelhameed et al. [CRESST], Phys. Rev. D 100, no.10, 102002 (2019) doi:10.1103/PhysRevD.100.102002 [arXiv:1904.00498 [astro-ph.CO]].
[10] C. Amole et al. [PICO], Phys. Rev. D 100, no.2, 022001 (2019) doi:10.1103/PhysRevD.100.022001 [arXiv:1902.04031 [astro-ph.CO]].
[11] E. Behnke, M. Besnier, P. Bhattacharjee, X. Dai, M. Das, A. Davour, F. Debris, N. Dhungana, J. Farine and M. Fines-Neuschild, et al. Astropart. Phys. 90, 85-92 (2017) doi:10.1016/j.astropartphys.2017.02.005 [arXiv:1611.01499 [hep-ex]].
[12] M. Schumann, J. Phys. G 46, no.10, 103003 (2019) doi:10.1088/1361-6471/ab2ea5 [arXiv:1903.03026 [astro-ph.CO]].
[13] M. Bauer and T. Plehn, Lect. Notes Phys. 959, pp. (2019) doi:10.1007/978-3-030-16234-4 [arXiv:1705.01987 [hep-ph]].
[14] J. Billard, M. Boulay, S. Cebrián, L. Covi, G. Fiorillo, A. Green, J. Kopp, B. Majorovits, K. Palladino and F. Petricca, et al. [arXiv:2104.07634 [hep-ex]].
[15] M. Ackermann et al. [Fermi-LAT], Phys. Rev. Lett. 115 (2015) no.23, 231301 doi:10.1103/PhysRevLett.115.231301 [arXiv:1503.02641 [astro-ph.HE]].
[16] V. A. Acciari et al. [MAGIC], Phys. Dark Univ. 28 (2020), 100529 doi:10.1016/j.dark.2020.100529 [arXiv:2003.05260 [astro-ph.HE]].
[17] H. Abdallah et al. [H.E.S.S.], Phys. Rev. Lett. 117 (2016) no.11, 111301 doi:10.1103/PhysRevLett.117.111301 [arXiv:1607.08142 [astro-ph.HE]].
[18] R. Abbasi et al. [IceCube], PoS ICRC2019, 541 (2020) doi:10.22323/1.358.0541 [arXiv:1908.07255 [astro-ph.HE]].
[19] A. Albert et al. [ANTARES], Phys. Lett. B 805, 135439 (2020) doi:10.1016/j.physletb.2020.135439 [arXiv:1912.05296 [astro-ph.HE]].
[20] K. Bays et al. [Super-Kamiokande], Phys. Rev. D 85, 052007 (2012) doi:10.1103/PhysRevD.85.052007 [arXiv:1111.5031 [hep-ex]].
[21] N. F. Bell, G. Busoni and I. W. Sanderson, JCAP 03 (2017), 015 doi:10.1088/1475-7516/2017/03/015 [arXiv:1612.03475 [hep-ph]].
[22] A. Albert, M. Bauer, J. Brooke, O. Buchmueller, D. G. Cerdeño, M. Citron, G. Davies, A. De Costa, A. De Roeck and A. De Simone, et al. Phys. Dark Univ. 16 (2017), 49-70 doi:10.1016/j.dark.2017.02.002 [arXiv:1607.06680 [hep-ex]].
[23] P. Ko, A. Natale, M. Park and H. Yokoya, JHEP 01 (2017), 086 doi:10.1007/JHEP01(2017)086 [arXiv:1605.07058 [hep-ph]].
[24] C. Englert, M. McCullough and M. Spannowsky, Phys. Dark Univ. 14 (2016), 48-56 doi:10.1016/j.dark.2016.09.002 [arXiv:1604.07975 [hep-ph]].
[25] T. Abe et al. [LHC Dark Matter Working Group], Phys. Dark Univ. 27 (2020), 100351 doi:10.1016/j.dark.2019.100351 [arXiv:1810.09420 [hep-ex]].
[26] S. Dutta, A. Goyal and L. K. Saini, JHEP 02 (2018), 023 doi:10.1007/JHEP02(2018)023 [arXiv:1709.00720 [hep-ph]].
[27] M. Bauer, M. Klassen and V. Tenorth, JHEP 07 (2018), 107 doi:10.1007/JHEP07(2018)107 [arXiv:1712.06597 [hep-ph]].
[28] M. Bauer, U. Haisch and F. Kahlhoefer, JHEP 05 (2017), 138 doi:10.1007/JHEP05(2017)138 [arXiv:1701.07422 [hep-ph]].
[29] S. Baek, P. Ko and J. Li, Phys. Rev. D 95 (2017) no.7, 075011 doi:10.1103/PhysRevD.95.075011 [arXiv:1701.04131 [hep-ph]].
[30] A. Carrillo-Monterverde, Y. J. Kang, H. M. Lee, M. Park and V. Sanz, JHEP 06 (2018), 037 doi:10.1007/JHEP06(2018)037 [arXiv:1803.02144 [hep-ph]].
[31] S. Kraml, U. Laa, K. Mawatari and K. Yamashita, Eur. Phys. J. C 77 (2017) no.5, 326 doi:10.1140/epjc/s10052-017-4871-0 [arXiv:1701.07008 [hep-ph]].
[32] G. Aad et al. [ATLAS], [arXiv:2109.00925 [hep-ex]].
[33] R. C. Cotta, J. L. Hewett, M. P. Le and T. G. Rizzo, Phys. Rev. D 88 (2013) 116009
doi:10.1103/PhysRevD.88.116009 [arXiv:1210.0525 [hep-ph]].

[34] J. Y. Chen, E. W. Kolb and L. T. Wang, Phys. Dark Univ. 2 (2013) 200
doi:10.1016/j.dark.2013.11.002 [arXiv:1305.0021 [hep-ph]].

[35] A. Crivellin, U. Haisch and A. Hibbs, Phys. Rev. D 91 (2015) 074028
doi:10.1103/PhysRevD.91.074028 [arXiv:1501.00907 [hep-ph]].

[36] P. Athron et al. [GAMBIT], [arXiv:2106.02056 [hep-ph]].

[37] H. Bharadwaj and S. Dutta, [arXiv:1901.00192 [hep-ph]].

[38] H. Bharadwaj and A. Goyal, Chin. Phys. C 45 (2021) no.2, 023114
doi:10.1088/1674-1137/abe50 [arXiv:2008.13621 [hep-ph]].

[39] N. Chen, J. Wang and X. P. Wang, [arXiv:1501.04486 [hep-ph]].

[40] H. Dreiner, M. Huck, M. Krämer, D. Schmeier and J. Tattersall, Phys. Rev. D 87
    no. 7, 075015 (2013) doi:10.1103/PhysRevD.87.075015 [arXiv:1211.2254 [hep-ph]].

[41] Y. J. Chae and M. Perelstein, “Dark Matter Search at a Linear Collider: Effective Operator
    Approach”, JHEP 1305 (2013) 138 [arXiv:1211.4008 [hep-ph]].

[42] P. J. Fox, R. Harnik, J. Kopp and Y. Tsai, Phys. Rev. D 85 (2012) 056011
doi:10.1103/PhysRevD.85.056011 [arXiv:1109.4398 [hep-ph]].

[43] N. F. Bell, J. B. Dent, A. J. Galea, T. D. Jacques, L. M. Krauss and T. J. Weiler, Phys. Rev.
    D 86 (2012) 096011 doi:10.1103/PhysRevD.86.096011 [arXiv:1209.0231 [hep-ph]].

[44] S. Dutta, D. Sachdeva and B. Rawat, Eur. Phys. J. C 77, no. 9, 639 (2017)
doi:10.1140/epjc/s10052-017-5188-8 [arXiv:1704.03994 [hep-ph]].

[45] F. Kahlhoefer, Int. J. Mod. Phys. A 32, no. 13, 1730006 (2017)
doi:10.1142/S0217751X1730006X [arXiv:1702.02430 [hep-ph]].

[46] V. A. Mitson, J. Phys. Conf. Ser. 651 (2015) no.1, 012023
doi:10.1088/1742-6596/651/1/012023 [arXiv:1402.3673 [hep-ex]].

[47] A. Boveia and C. Doglioni, Ann. Rev. Nucl. Part. Sci. 68 (2018), 429-459
doi:10.1146/annurev-nucl-101917-021008 [arXiv:1810.12238 [hep-ex]].

[48] A. De Simone and T. Jacques, Eur. Phys. J. C 76 (2016) no.7, 367
doi:10.1140/epjc/s10052-016-4208-4 [arXiv:1603.08002 [hep-ph]].

[49] B. Bhattacharjee, D. Choudhury, K. Harigaya, S. Matsumoto and M. M. Nojiri, JHEP 04
    (2013), 031 doi:10.1007/JHEP04(2013)031 [arXiv:1212.5013 [hep-ph]].

[50] S. Baek, P. Ko and P. Wu, JCAP 07 (2018), 008 doi:10.1088/1475-7516/2018/07/008
    [arXiv:1709.00697 [hep-ph]].

[51] C. Kilic, M. D. Klimek and J. H. Yu, Phys. Rev. D 91 (2015) no.5, 054036
doi:10.1103/PhysRevD.91.054036 [arXiv:1501.02202 [hep-ph]].

[52] M. A. Gomez, C. B. Jackson and G. Shaughnessy, JCAP 12 (2014), 025
doi:10.1088/1475-7516/2014/12/025 [arXiv:1404.1918 [hep-ph]].

[53] B. Fuks, J. Phys. Conf. Ser. 1271 (2019) no.1, 012017 doi:10.1088/1742-6596/1271/1/012017
    [arXiv:1903.02477 [hep-ph]].
[54] S. Colucci, B. Fuks, F. Giacchino, L. Lopez Honorez, M. H. G. Tytgat and J. Vandecasteele, Phys. Rev. D 98 (2018), 035002 doi:10.1103/PhysRevD.98.035002 [arXiv:1804.05068 [hep-ph]].

[55] M. Garny, J. Heisig, M. Hufnagel and B. Lülf, Phys. Rev. D 97 (2018) no.7, 075002 doi:10.1103/PhysRevD.97.075002 [arXiv:1802.00814 [hep-ph]].

[56] Y. Zhang, Phys. Lett. B 720 (2013), 137-141 doi:10.1016/j.physletb.2013.01.063 [arXiv:1212.2750 [hep-ph]].

[57] A. M. Sirunyan et al. [CMS], JHEP 03, 141 (2019) doi:10.1007/JHEP03(2019)141 [arXiv:1901.01553 [hep-ex]].

[58] A. M. Sirunyan et al. [CMS], Phys. Rev. Lett. 122 (2019) no.1, 011803 doi:10.1103/PhysRevLett.122.011803

[59] G. Aad et al. [ATLAS], Eur. Phys. J. C 81 (2021), 860 doi:10.1140/epjc/s10052-021-09566-y [arXiv:2011.09308 [hep-ex]].

[60] [CMS], CMS-PAS-B2G-15-007.

[61] T. Lin, E. W. Kolb and L. T. Wang, Phys. Rev. D 88 (2013) no.6, 063510 doi:10.1103/PhysRevD.88.063510 [arXiv:1303.6638 [hep-ph]].

[62] U. Haisch, F. Kahlhoefer and J. Unwin, JHEP 07 (2013), 125 doi:10.1007/JHEP07(2013)125 [arXiv:1208.4605 [hep-ph]].

[63] J. Goodman, M. Ibe, A. Rajaraman, W. Shepherd, T. M. P. Tait and H. B. Yu, Phys. Rev. D 82, 116010 (2010) doi:10.1103/PhysRevD.82.116010 [arXiv:1008.1783 [hep-ph]].

[64] C. Arina, M. Backović, E. Conte, B. Fuks, J. Guo, J. Heisig, B. Hespel, M. Krämer, F. Maltoni and A. Martini, et al. JHEP 11 (2016), 111 doi:10.1007/JHEP11(2016)111 [arXiv:1605.09242 [hep-ph]].

[65] O. Mattelaer and E. Vryonidou, Eur. Phys. J. C 75 (2015) no.9, 436 doi:10.1140/epjc/s10052-015-3665-5 [arXiv:1508.00564 [hep-ph]].

[66] Manuel Drees and Mihoko M. Nojiri, Phys. Rev. D 48, 3483, doi.org/10.1103/PhysRevD.48.3483

[67] J. Hisano, K. Ishiwata and N. Nagata, Phys. Rev. D 82 (2010) 115007 doi:10.1103/PhysRevD.82.115007 [arXiv:1007.2601 [hep-ph]].

[68] J. Hisano, K. Ishiwata, N. Nagata and M. Yamanaka, Prog. Theor. Phys. 126 (2011) 435 doi:10.1143/PTP.126.435 [arXiv:1012.5455 [hep-ph]].

[69] J. Hisano, K. Ishiwata, N. Nagata and T. Takesako, JHEP 1107 (2011) 005 doi:10.1007/JHEP07(2011)005 [arXiv:1104.0228 [hep-ph]].

[70] J. Hisano, K. Ishiwata and N. Nagata, Phys. Lett. B 706 (2011) 208 doi:10.1016/j.physletb.2011.11.017 [arXiv:1110.3719 [hep-ph]].

[71] J. Hisano, K. Ishiwata and N. Nagata, Phys. Rev. D 87 (2013) 035020 doi:10.1103/PhysRevD.87.035020 [arXiv:1210.5985 [hep-ph]].

[72] J. Hisano, R. Nagai and N. Nagata, JHEP 1505 (2015) 037 doi:10.1007/JHEP05(2015)037 [arXiv:1502.02244 [hep-ph]].

[73] J. Hisano, K. Ishiwata and N. Nagata, JHEP 1506 (2015) 097 doi:10.1007/JHEP06(2015)097 [arXiv:1504.00915 [hep-ph]].
[94] M. G. Aartsen et al. [IceCube and PICO], Eur. Phys. J. C 80 (2020) no.9, 819
doi:10.1140/epjc/s10052-020-8069-5 [arXiv:1907.12509 [astro-ph.HE]].

[95] R. H. Helm, Phys. Rev. 104 (1956), 1466-1475 doi:10.1103/PhysRev.104.1466.

[96] J. D. Lewin and P. F. Smith, Astropart. Phys. 6 (1996), 87-112
doi:10.1016/S0927-6505(96)00047-3

[97] M. Cirelli, E. Del Nobile and P. Panci, JCAP 10 (2013), 019
doi:10.1088/1475-7516/2013/10/019 [arXiv:1307.5955 [hep-ph]].

[98] G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, Comput. Phys. Commun. 185
(2014), 960-985 doi:10.1016/j.cpc.2013.10.016 [arXiv:1305.0237 [hep-ph]].

[99] H. Y. Cheng and C. W. Chiang, JHEP 07 (2012), 009 doi:10.1007/JHEP07(2012)009
[arXiv:1202.1292 [hep-ph]].

[100] H. Y. Cheng, Phys. Lett. B 219 (1989), 347-353 doi:10.1016/0370-2693(89)90402-4

[101] M. V. Polyakov, A. Schafer and O. V. Teryaev, Phys. Rev. D 60 (1999), 051502
doi:10.1103/PhysRevD.60.051502 [arXiv:hep-ph/9812393 [hep-ph]].