Top partners tackling vector dark matter

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The WIMP-nucleon scattering cross section in a simple dark matter model and its constraints from the latest direct detection experiment are treated here at the loop level. We consider a scenario with an emerging vector dark matter field interacting with the Standard Model quarks, via loop contributions that are sourced by a scalar mediator. The involved parameter space for the dark matter-mediator masses is constrained by the Xenon1T limit and the neutrino floor. The current direct detection results are eluded by invoking the top partners in a Composite Higgs model, whose mass will help us in properly suppressing the WIMP-nucleon cross section.

The dark matter (DM) stands as one of the most tantalising components in our Universe, becoming nowadays one of the most mysterious conundrums unknown so far. The particle physics and astrophysics community are currently seeking for DM signals in order to unveil its nature and involved properties. A weak interacting massive particle (WIMP) [1] has been a long-standing DM candidate, having resurfaced since the successful experimental confirmation of the Standard Model (SM) Higgs sector. Singlet elementary scalars have been additionally considered in the last years [2–6], extended to the SM quark mass [7, 8] and further to the strong sector at the scale $\Theta$ [9]. Nonetheless, such coupling might appear from different frameworks, as in the case of CHMs. These scenarios represent an appealing option to supersymmetry, rendering the DM candidates much lighter than the SM particles via the scalar mediator $\eta$. We assume in this work a framework consisting of composite scalars in the light of composite Higgs models (CHMs). These scenarios represent an appealing option to supersymmetry, rendering the DM candidates much lighter than the new physics compositeness scale, as they enter onto the stage as pseudo-Nambu-Goldstone bosons (pNGBs) of a new strongly interacting sector. Specific models have been explored for instance in [7–10]. In other contexts, there have been interesting claims for vector dark matter (VDM), arising from an additional $U(1)_X$ gauge symmetry [11–16], recently analysed and compared against scalar dark matter (SDM) models [17].

On the other hand, different WIMP-ordinary matter interactions exit, dividing thus the WIMP searches into three methods: direct detection (DD), indirect detection (ID) and collider search. Among them, DD experiments pursue the scattering of dark matter particles off atomic nuclei. No clear signals of dark matter from DD experiments have been detected till now, leaving us instead with upper limits on the WIMP-nuclei scattering cross section. Enhanced by the squared total nucleon number with upper limits on the WIMP-nuclei scattering cross section, the DD constraint will be compared with those from indirect detection and collider search in a future work. This approach has been thoroughly used for analysing dark matter searches in DD, ID and collider experiments [20–30].

In this simple scenario, the dark matter annihilation occurs through the s-channel exchange of either a spin-0 or spin-1 mediator. The latter case has been constrained by $Z' \to \text{dijet}$ searches at the Large Hadron Collider (LHC) [31–34]. We assume henceforth a scalar mediator scenario\textsuperscript{1}, generically written as [37]

$$\mathcal{L}_X = -g_{\eta X} M_X X^\mu X^\mu \eta - i \eta \sum_q g_{\eta q} \bar{q} \gamma^\mu q.$$  \hspace{1cm} (1)

VDM model, exploring as well the constraints imposed by the latest direct detection experiments [18, 19]. The DD constraint will be compared with those from indirect detection and collider search in a future work. This approach has been thoroughly used for analysing dark matter searches in DD, ID and collider experiments [20–30].

The analysis of dark matter searches in the literature commonly assumes $g_{\eta X} = g_{\eta q} = 1$. Under this assumption, these dark matter models are described by two parameters, the dark matter and the mediator masses $M_X$ and $M_\eta$ respectively. Nonetheless, such coupling $g_{\eta q}$ might appear from different frameworks, as in the case of CHM scenarios. Invoking the minimal global symmetry $\mathcal{G} = SO(5)$ [38], spontaneously broken to $SO(4)$ by the strong sector at the scale $\theta$, and implementing the elementary sector fields together with the top partners $\Psi$ transforming in the unbroken $SO(4)$, we can have the following mass mixing terms

$$\mathcal{L}_{\text{mix}} = y_L f \left( \bar{q}_L U \right)_i \left( \Psi_{4R} \right)_i^\dagger + y_R f \left( \bar{q}_R U \right)_i \left( \Psi_{4L} \right)_i^\dagger + \text{h.c.},$$

$$+ \tilde{y}_L f \left( \bar{q}_L U \right)_5 \left( \Psi_{1R} \right) + \tilde{y}_R f \left( \bar{q}_R U \right)_5 \left( \Psi_{1L} \right) + \text{h.c.}$$

\hspace{1cm} (2)

\textsuperscript{1} The spin-0 mediator case is generically suppressed by the Yukawa couplings via the minimal flavor violation prescription, the which suppresses the scalar bilinear $\eta \bar{q} q$ and pseudo-scalar one $\eta \bar{q} \gamma_5 q$ by the SM quark mass $m_q$ [35, 36]. In here the scalar mediator will couple to the SM fermions proportionally to their mass as it can be seen later on.

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and
\[ L'_{\text{mix}} = y_L f \left( U^T q_1^{14} U \right)_{55} \tilde{\Psi}_{4R}^i + y_R f \left( U^T q_1^{14} U \right)_{55} U^1_R + \]
\[ + y_R f \left( U^T q_1^{14} U \right)_{55} u^1_R + \text{h.c.} \]
with the pNGB fields encoded in the 5 × 5 Goldstone matrix \( U \), defined as
\[ U = \exp \left[ i \frac{\sqrt{2}}{f} \Pi^i T^i \right], \]
with \( T^i \) the coset \( SO(5)/SO(4) \)-generators (App. A), where \( \Pi^i \) are the pNGB fields and \( f \) the decay constant. Small mixing couplings \( y_{L(R)} \) and \( \tilde{y}_{L(R)} \) trigger the Goldstone Boson (GB) symmetry breaking. The top partners fourplet \( \tilde{\Psi}_4 \) and singlet \( \tilde{\Psi}_1 \) are embedded in the unbroken \( SO(4) \) as
\[ \tilde{\Psi}_4 = \frac{1}{\sqrt{2}} \left( \begin{array}{cc} iB - iX_{5/3} \\ B + X_{5/3} \\ i\tilde{T} + iX_{2/3} \\ -\tilde{T} + X_{2/3} \end{array} \right), \quad \tilde{\Psi}_1 = \tilde{T}. \]

One SM-like quark doublet \((T, B)\), plus an exotic 7/6-hypercharge \((X_{5/3}, X_{2/3})\) are contained in \( \tilde{\Psi}_4 \), while \( \tilde{\Psi}_1 \) contains only one exotic top-like state \( \tilde{T} \). The elementary sector is prescribed by the partial compositeness mechanism through the GB symmetry breaking terms \( y \tilde{q} \tilde{O}_q \), with the strong sector operator \( \tilde{O}_q \) transforming in one of the \( SO(5) \)-representations. We assume here two elementary sector embeddings: either as a fundamental 5
\[ q^5_L = \frac{1}{\sqrt{2}} \left( \begin{array}{cccc} id_L, d_L, iu_L, -u_L, 0 \end{array} \right)^T, \]
\[ u^5_R = \left( \begin{array}{c} 0 \end{array} \right)^T, \]
or the 14 representation
\[ q^{14}_L = \frac{1}{\sqrt{2}} \left( \begin{array}{ccc} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \\ id_L & d_L & iu_L \end{array} \right), \quad u^1_R. \]

In the former scenario, both fermion chiralities have elementary representatives coupled to the strong sector through 5-plets, whereas in the latter the right-handed quark enters as a totally composite state. All these matter content shape four models, here denoted as \( M_{\Psi + q} = \{ M_{4+5}, M_{4+14}, M_{1+5}, M_{1+14} \} \).

The aforementioned fermionic matter is coupled to the scalar resonance \( \eta \), described here as a singlet of \( SU(2)_L \times SU(2)_R \), and previously considered in CHMs [39]. In fact, in the 5-plets scenarios coupled to \( \eta \), the pseudoscalar couplings can be retrieved from [40]
\[ L_{q\tilde{q} \tilde{\Psi}} = \eta \left[ y_{q\psi} \left( \gamma^5 U \right)_{55} \tilde{\Psi}_{4R}^i + y_{u\psi} \left( \gamma^5 U \right)_{55} \tilde{\Psi}_{4L}^i + \right. \]
\[ + \tilde{y}_{q\psi} \left( \gamma^5 U \right)_{55} \tilde{\Psi}_{1R} + \tilde{y}_{u\psi} \left( \gamma^5 U \right)_{55} \tilde{\Psi}_{1L} \left. \right] + \text{h.c.}, \]
with the couplings \( y_{q(w)\psi} \) and \( \tilde{y}_{q(w)\psi} \) controlling the scalar-fermion interactions, and responsible for the decays of \( \eta \) into the pure SM final states and single-double partner final modes. For the 14-plets models we have similar Lagrangians
\[ L_{q\tilde{q} \tilde{\Psi}} = \eta \left[ y_{q\psi} \left( U^T q_1^{14} U \right)_{55} \tilde{\Psi}_{4R}^i + y_{u\psi} \left( U^T q_1^{14} U \right)_{55} \tilde{\Psi}_{4L}^i + \right. \]
\[ + \tilde{y}_{q\psi} \left( U^T q_1^{14} U \right)_{55} \tilde{\Psi}_{1R} + \tilde{y}_{u\psi} \left( U^T q_1^{14} U \right)_{55} \tilde{\Psi}_{1L} \left. \right] + \text{h.c.}, \]

Via top partner’s equations of motion (App. B), the Lagrangians in (9) and (10) translates into a set of interactions involving SM top quark content coupled to the spin-0 mediator. Their associated couplings are appropriately mapped onto the effective couplings correspondingly at each model as
\[ M_{4+5} : \quad g_{q\tilde{q} \tilde{\Psi}} = \sqrt{2} \left( \frac{\eta_{\tilde{R}} \text{Im}(q_{\tilde{q} \psi}) - \eta_{\tilde{L}} (\eta_{\tilde{R}}^2 + 1) \text{Im}(y_{u\psi})}{(\eta_{\tilde{R}}^2 + 1)^{3/2}} \right), \]
\[ M_{4+14} : \quad g_{q\tilde{q} \tilde{\Psi}} = \sqrt{2} \left( \frac{i \eta_{\tilde{L}} \eta_{\tilde{R}} \text{Re}(q_{\tilde{q} \psi}) - (\eta_{\tilde{R}}^2 + 1) \text{Im}(y_{u\psi})}{(\eta_{\tilde{R}}^2 + 1)^{3/2}} \right), \]
\[ M_{1+5} : \quad g_{q\tilde{q} \tilde{\Psi}} = \sqrt{2} \left( \frac{\eta_{\tilde{L}} \text{Im}(q_{\tilde{q} \psi}) - \tilde{\eta}_{\tilde{R}} (\tilde{\eta}_{\tilde{R}}^2 + 1) \text{Im}(y_{u\psi})}{(\tilde{\eta}_{\tilde{R}}^2 + 1)^{3/2}} \right), \]
\[ M_{1+14} : \quad g_{q\tilde{q} \tilde{\Psi}} = -\sqrt{2} \eta_{\tilde{L}} + 1 \text{Im}(y_{u\psi}) \]

with \( q = c, b, t \), where \( \xi = \sqrt{v^2/f^2} \) with \( v = 246 \text{GeV} \), and the coefficients \( \eta_{L(R)} \equiv y_{L(R)} f / M_4 \) and \( \tilde{\eta}_{L(R)} \equiv \tilde{y}_{L(R)} f / M_4 \) (App. B). The mass scales \( M_{4(1)} \) are associated to the fermionic resonances \( \Psi_{4(1)} \). They are commonly parametrised through the generic coupling \( g_{\Psi} \) as \( M_4 = M_1 = M_{\Psi} \equiv g_{\Psi} f \). Notice how the top partner mass dependence is directly exhibited through the parameters \( \eta_{L(R)} \) and \( \tilde{\eta}_{L(R)} \), while it is also manifested via the parameter \( \xi \) and the relation \( M_{\Psi} = y_{\Psi} f \). When
exploring the cross sections and their sensitivity to the top partner mass variations, two options appear:

- either to vary the top partner mass scale while properly changing the coupling $g_\Psi$, conserving thus $\xi$ at a fixed value, or
- just to vary $M_\Psi$ while keeping fixed $g_\Psi$, implying therefore a varying $\xi$.

The second option effectively suppresses much more the couplings in (11), rather than the first one. This feature will be reflected later on in the computation of the WIMP-nucleon cross section. From the effective Lagrangian in (1), the WIMP-quark scattering amplitudes are originated via loop diagrams as it is shown in Fig. 1, and schematically written as

$$\sum_q g_{3q}^2 g_{qq}^2 g_{\text{loop}}^2 X^\mu_X X^\nu_q \Rightarrow \sigma_{\text{SI}}^{\text{DMN}} = \frac{4 \mu_{\text{DMN}}^2}{\alpha} \left| C(M_X, M_\eta) \right|^2$$

(12)

where we have transformed the WIMP-quark amplitude to the WIMP-nucleon amplitude, in terms of the nucleon form factor $f_{\text{TG}} \approx 0.89$ for the gluon condensate [41–43], with $C_{\text{loop}}^2 = C(M_X, M_\eta, m_q)$ and

$$C(M_X, M_\eta) \approx \sum_q g_{3q}^2 g_{qq}^2 m_N^2 \frac{2}{27} f_{\text{TG}} C(M_X, M_\eta, m_q),$$

(13)

with $\mu_{\text{DMN}} = M_X m_N/(M_X + m_N)$ the reduced mass and $m_N$ the nucleon mass. The loop coefficient $C(M_X, M_\eta, m_q)$ is reported in App. C, and it has been computed using FeynCalc [44, 45] with the resulting Passarino-Veltman functions [46] being further reduced via Package-X [47]. In Fig. 2, we display the result of the 1-loop SI DM-nucleon scattering cross section versus the DM mass, assuming $g_{X_\eta} = 1$ and for a generic situation with $g_{qq} = 1$ and mediator masses $M_\eta = 10, 50, 100$GeV (thick, dashed, dotted curves). Notice how the heavier mediator reduces the WIMP-nucleon cross section, while the Xenon1T upper limit (thick black) is partially evaded for a mediator of 50GeV, while the neutrino floor (brown) can be eluded in a tiny DM mass region with a heavier mediator of 100GeV.

When implementing the effective couplings in Eq. (11), the Xenon1T limit and the neutrino floor can be partially evaded in some of the models for a fixed mediator mass $M_\eta = 10, 50, 100$GeV (upper-centred-lower plots) and two different top partner masses $M_\Psi = 500, 1500$GeV (thick-dashed). See the text for details.
tially eluding the Xenon1T limits excepting in \( M_{1+14} \) (orange). We have set \( \text{Im}(y_{qL}) = \text{Im}(y_{uL}) = \text{Re}(y_{qL}) = \text{Im}(y_{uL}) = 1/2 \) and \( \text{Im}(\tilde{y}_{qL}) = \text{Im}(\tilde{y}_{uL}) = 1/2 \). Additionally, the Yukawa couplings are chosen as \( y_R = \tilde{y}_R = 1 \), while \( y_L \) and \( \tilde{y}_L \) are properly fixed in order to maintain the top mass at its experimental value through the corresponding top mass formula. This is obtained for the 5-plet and 14-plet scenarios in [40, 48].

It is remarkable to observe how for a mediator mass of \( M_\eta = 50\text{GeV} \) (Fig. 3 centred plot), the Xenon1T limit, and even the neutrino floor, are evaded. Indeed, for \( M_\Psi = 500\text{GeV} \) all the models are partially below Xenon1T, with \( M_{4+5}, M_{4+14} \) and \( M_{1+5} \) (red-blue-green) in a broader DM mass region. Increasing the scale up to \( M_\Psi = 1500\text{GeV} \), the latter models completely elude Xenon1T in the DM mass range \( 10-10000\text{GeV} \). Such models simultaneously evade the neutrino floor in a smaller mass regions of \( 10-600\text{GeV} \), \( 10-400\text{GeV} \) and \( 10-300\text{GeV} \) correspondingly, whilst \( M_{1+14} \) in the reduced range \( 10-50\text{GeV} \).

The previous suppression is further enhanced for a higher mediator mass \( M_\eta = 100\text{GeV} \) (Fig. 3 lower plot). Indeed, for \( M_\Psi = 500\text{GeV} \), the Xenon1T limit is notoriously evaded in all the scanned DM mass range for the scenarios \( M_{4+5}, M_{4+14} \) and \( M_{1+5} \). In the same DM mass region, such models are remarkably below the neutrino floor for \( M_\Psi = 1500\text{GeV} \). In this case, \( M_{1+14} \) becomes completely below Xenon1T, partially eluding the neutrino floor in the range \( 10-400\text{GeV} \).

The cross sections displayed in Fig. 3 were computed for a varying \( \xi \), i.e. a varying \( M_\Psi \) with fixed \( g_\Psi \). From (12)-(13) and the \( \xi \)-dependence in (11), we roughly have

\[
\sigma_{\text{DMN}}^\text{SI} \approx g_{qq}^4 \sim \xi^2 \sim \frac{g_\Psi^4 v^4}{M_\Psi^4}.
\]

Augmenting \( M_\Psi \) and fixing \( g_\Psi \), clearly introduces a suppressing factor in addition to the one induced by the \( \eta \) and \( \tilde{\eta} \)-dependence in (11). This situation\(^2\) is therefore more favoured in eluding the DD experiments, as well as the neutrino floor in some DM mass regions, compared with the scenario of \( M_\Psi \) and \( g_\Psi \) simultaneously changing (fixed \( \xi \)). The latter situation is explicitly depicted in Fig. 5 App. D.

The parameter space for the DM-mediator masses permitted by Xenon1T (gray), and the regions with the cross section below the neutrino floor (orange) are shown in Fig. 4, for both the fourplet and singlet models (top-bottom plots). Notice that mediator masses in the range \( 30-90\text{GeV} \) are allowed by Xenon1T in the fourplet models, while \( 35-100\text{GeV} \) in \( M_{1+14} \) and \( 65-180\text{GeV} \) in \( M_{1+14} \) for \( M_\Psi = 500\text{GeV} \) (darker areas). In this case, the neutrino floor becomes eluded for a mediator mass of \( 10-150\text{GeV} \) in the DM mass regions of \( 10-5000\text{GeV} \) (\( M_{4+5} \)), \( 10-3000\text{GeV} \) (\( M_{4+14} \)), \( 10-1000\text{GeV} \) (\( M_{1+5} \)) and \( 10-125\text{GeV} \) (\( M_{1+14} \)). As soon as the top partner is heavier as \( M_\Psi = 1500\text{GeV} \) (lighter areas), smaller mediator masses in the approximate range \( 15-35\text{GeV} \) are constrained in the fourplet scenarios by Xenon1T, while \( 15-45\text{GeV} \) in \( M_{1+5} \) and \( 30-90\text{GeV} \) for \( M_{1+14} \). The neutrino floor becomes eluded for a mediator mass of \( 10-150\text{GeV} \) in all the explored DM mass region. A smaller DM mass range of \( 10-5000\text{GeV} \) in \( M_{1+14} \) permits to be below the neutrino floor.

All in all, the scenarios \( M_{4+5}, M_{4+14} \) and \( M_{1+5} \) better suppress the WIMP-nucleon cross sections rather than \( M_{1+14} \). In fact, in the limit of a large \( M_\Psi \), its coupling \( g_{qq} \) in (11) involves a suppression only from the \( \xi \)-dependence, unlike to the others whose suppression is directly enhanced by the contribution of the \( \eta \) and \( \tilde{\eta} \)-parameters. This feature disfavours \( M_{1+14} \) in effectively eluding the latest DD Xenon1T bounds. Future observations will help us in discriminating the best framework among \( M_{4+5}, M_{4+14} \) and \( M_{1+5} \) for the explanation of the DD experiments.

I. SUMMARY

The WIMP-nucleon scattering cross section in a simple vector dark matter model is analysed at the loop-level. A scalar mediator is coupled only to the vector dark field and to the SM quarks via pseudo-bilinear in-

\(^2\) In Fig. 3 we have explored two situations \( M_\Psi = 500, 1500\text{GeV} \) for a fixed \( g_\Psi = 1 \), which corresponds to \( \xi \sim 0.26, 0.03 \). The former value is compatible with the latter EWPT bounds, as well as the vector resonance direct production bounds at LHC and the expected single Higgs production at the LHC. The latter value is favoured by the 95\% combined limit from direct production of vector resonances at the LHC [49] and for a mass of \( \sim 2 \text{ TeV} \).
teractions. The latest direct detection experiments are implemented to bound the involved parameter space for the dark matter-mediator masses. We provide effective mediator-quark couplings whose effect suppress the total WIMP-nuclei cross section, evading thus the recent direct detection Xenon1T limits. We motivate such suppression from New Physics scales reachable at the LHC and inspired by Composite Higgs scenarios. They naturally source the top partners, whose implied masses suppress the effective pseudo-bilinear couplings. In summary, we will be able to evade the upper current experimental bounds, furthermore predicting reasonable regions for the dark matter-mediator masses, and experimentally consistent in a coherent theoretical framework.

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Appendix A: CCWZ formalism

The \(SO(4) \cong SU(2)_L \times SU(2)_R\) unbroken generators and the broken ones parametrising the coset \(SO(5)/SO(4)\) in the fundamental representation are

\[
(T^{\alpha}_{\chi})_{IJ} = -\frac{i}{2} \left[ \frac{1}{2} \epsilon^{\alpha bc} \left( \delta^c_j \delta^a_i - \delta^c_i \delta^a_j \right) \pm \left( \delta^c_j \delta^a_i - \delta^c_i \delta^a_j \right) \right],
\]

and

\[
T^\chi_{IJ} = -\frac{i}{\sqrt{2}} \left( \delta^c_j \delta^a_i - \delta^c_i \delta^a_j \right),
\]

with \(\chi = L, R\) and \(a = 1, 2, 3\), while \(i = 1, \ldots, 4\). The normalization of \(T^{\alpha}\)'s is chosen as \(\text{Tr}[T^A, T^B] = \delta^{AB}\).

Appendix B: Top partners EOM

For a heavy top partners mass scenario the corresponding fields may be integrated out from the physical spectrum, via their associated equations of motion. Concerning the models \(M_{4+5}\) and \(M_{1+4+5}\) altogether, and after diagonalising the mass terms at \(\mathcal{L}_{\text{mix}}\) in (2), plus the terms associated to their kinetic composite sector, we obtain the field redefinitions for the left handed components as

\[
T_L \rightarrow \frac{-\xi}{4} \frac{\eta_L}{\eta_L^2 + 1} t_L, \quad \tilde{T}_L \rightarrow \sqrt{\frac{\xi}{2}} \frac{\tilde{\eta}_L}{(\eta_L^2 + 1) \sqrt{\eta_L^2 + 1}} t_L,
\]

\[
X_{2/3,L} \rightarrow \frac{\xi}{4} \frac{\eta_L}{\sqrt{\eta_L^2 + 1}} t_L, \quad \mathcal{B}_L \rightarrow 0,
\]

while for the right handed components one obtains

\[
T_R \rightarrow \sqrt{\frac{\xi}{2}} \frac{\eta_R}{(\eta_R^2 + 1) \sqrt{\eta_R^2 + 1}} t_R, \quad \tilde{T}_R \rightarrow \frac{-\xi}{2} \frac{\tilde{\eta}_R}{\eta_R^2 + 1} t_R,
\]

\[
X_{2/3,R} \rightarrow -\sqrt{\frac{\xi}{2}} \frac{\eta_R}{\sqrt{\eta_R^2 + 1}} t_R, \quad \mathcal{B}_R \rightarrow -\tilde{\eta}_R b_R,
\]

where the parameters \(\eta_{L(R)}\) are defined through

\[
\eta_{L(R)} \equiv \frac{\eta_{L(R)} f}{M_4}, \quad \tilde{\eta}_{L(R)} \equiv \frac{\tilde{\eta}_{L(R)} f}{M_4}. \quad (B1)
\]

Likewise, from \(\mathcal{L}_{\text{mix}}\) the in (3) and the associated kinetic terms for \(M_{4+14}\) and \(M_{4+14}\), we obtain

\[
T_L \rightarrow -\frac{5\xi}{4} \frac{\eta_L}{\eta_L^2 + 1} t_L, \quad \tilde{T}_L \rightarrow -\frac{\sqrt{T_L}}{\sqrt{\eta_L^2 + 1}} \tilde{\eta}_L t_L,
\]

\[
X_{2/3,L} \rightarrow \frac{3\xi}{4} \frac{\eta_L}{\sqrt{\eta_L^2 + 1}} t_L, \quad \mathcal{B}_L \rightarrow \frac{-\xi}{2} \frac{\eta_L}{\eta_L^2 + 1} b_L,
\]

while for the right handed components one obtains

\[
T_R \rightarrow -\sqrt{\frac{2\xi}{5}} \frac{\sqrt{\eta_R^2 + 1}}{\eta_R^2 + 1} \eta_L \eta_R t_R, \quad \tilde{T}_R \rightarrow -\tilde{\eta}_R t_R,
\]

\[
\mathcal{B}_R \rightarrow -\tilde{\eta}_R b_R, \quad X_{2/3,R} \rightarrow 0 \quad (B4)
\]

with the coefficients \(\eta\) similarly defined as in (B3), but with the Yukawa couplings \(y\) and \(\tilde{y}\) corresponding to those for the \(M_{4+14}\) and \(M_{1+4+14}\) models. Both the left and right handed components for the field \(X_{5/3}\) are zero in all the models.

Appendix C: Loop function

The loop function \(C(M_X, M_B, m_q)\), obtained from the box diagrams in Fig. 1, is reported here in terms of the Passarino-Veltman functions and the Mandelstam variables \(s = (p_1 + p_2)^2\), \(t = (p_1 - k_1)^2\) and \(u = (p_1 - k_2)^2\) for the process \(q(p_1) X(p_2) \rightarrow q(k_1) X(k_2)\). We obtain

\[
C(M_X, M_B, m_q) = \frac{m_q M_X^2}{16\pi^2} \quad (C1)
\]

\[
[D_0 \left( m_q^2, M_X^2, M_X^2, M_X^2, m_q^2, s, t, u \right) + D_0 \left( m_q^2, M_X^2, M_X^2, m_q^2, t, u, m_q^2, M_X^2, M_X^2 \right) + \ldots]
\]
Appendix D: WIMP-nucleon cross sections at fixed $\xi$

For a fixed $\xi$ value, i.e. $M_0$ and $g_\phi$ properly varying, the effective pseudo-bilinear couplings in Eq. (11) will depend on the top partner mass $M_t$ only via the parameters $\eta_{L(R)}$ and $\eta_{L(L)}$. In this case, the induced suppression in the cross sections is less notorious as in the case of a decreasing $\xi$, i.e. an increasing $M_0$ with fixed $g_\phi$.

Fig. 5 shows the SI DM-nucleon scattering cross section for mediator masses $M_0 = 50, 100\text{GeV}$ (left-right plots) and $M_0 = 500, 1500\text{GeV}$ (thick-dashed curves), with the corresponding couplings $g_\phi = 1, 3$, such that $\xi$ remains fixed at a rough value of 0.2.

Notice how the cross sections suffer of a less suppression when augmenting the top partner mass, compared with those in Fig. 3. The evasion of Xenon1T mainly depends on the mediator mass for this case.

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