Impact of a XENONnT Signal on LHC Dijet Searches

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ABSTRACT: It is well-known that dark matter (DM) direct detection experiments and the LHC are complementary, since they probe physical processes occurring at different energy scales. And yet, there are aspects of this complementarity which are still not fully understood, or exploited. For example, what is the impact that the discovery of DM at XENONnT would have on present and future searches for DM in LHC final states involving a pair of hadronic jets? In this work we investigate the impact of a XENONnT signal on the interpretation of current dijet searches at the LHC, and on the prospects for dijet signal discovery at the High-Luminosity (HL) LHC in the framework of simplified models. Specifically, we focus on a general class of simplified models where DM can have spin 0, 1/2 or 1, and interacts with quarks through the exchange of a scalar, pseudo-scalar, vector, or pseudo-vector mediator. We find that exclusion limits on the mediator’s mass and its coupling to quarks from dijet searches at the LHC are significantly affected by a signal at XENONnT, and that about 100 signal events at XENONnT would drastically narrow the region in the parameter space of simplified models where a dijet signal can be discovered at 5\sigma C.L. at the HL-LHC.
1 Introduction

Observations in a wide range of astronomical and cosmological systems show that the Universe contains about five times as much dark matter (DM) as baryonic matter [1]. While the nature of DM remains unknown, the hypothesis that DM is made of yet unidentified particles is the one explored most extensively [2]. If DM is made of Weakly Interacting Massive Particles (WIMPs), as predicted by many theories beyond the Standard Model (BSM) addressing the hierarchy problem, e.g. supersymmetric extension of the Standard Model (SM), it can potentially be observed in the next stage of direct detection and particle collider experiments [3]. DM direct detection experiments primarily search for nuclear recoils induced by the non-relativistic scattering of Milky Way DM particles in low-background detectors located deep underground [4, 5]. The null results of present DM direct detection experiments place severe constraints on the strength with which DM couples to the fundamental constituents of matter. The most stringent limits from direct detection experiments on the strength of DM-nucleus interactions for WIMPs heavier than about 10 GeV are currently set by the XENON1T collaboration [6], improving on previous results from LUX [7] and PandaX-II [8]. They reported a 90% C.L. exclusion limit on the elastic spin-independent DM-nucleon scattering cross-section with a minimum of $4.1 \times 10^{-47}$ cm$^2$ around a DM mass of 30 GeV. XENONnT, the upgrade of XENON1T, is expected to operate from 2019 onwards using about 7 ton of ultra-pure liquid xenon as target material [9]. Pursuing a complementary approach to direct detection experiments, the Large Hadron Collider (LHC)
at CERN with data analysed by the ATLAS and CMS collaborations is also searching for DM. At the LHC, DM can be produced in the collision of energetic protons, and its production be inferred through the observation of missing transverse momentum in the final state of such collisions. The next run of the LHC (Run 3) will start in 2021 [10]. It will operate at the centre-of-mass energy of $\sqrt{s} = 13$ TeV and reach the expected integrated luminosity of 300 fb$^{-1}$ in 2023. The LHC Run 3 will be followed by the high luminosity run of the LHC (HL-LHC), which is expected to start in 2026, reaching an integrated luminosity of 3000 fb$^{-1}$.

LHC data on processes that might involve DM have been interpreted within different theoretical frameworks. Ultraviolet (UV) complete theories (with a focus on supersymmetric theories) and Effective Field Theories (EFTs) for DM-quark and -gluon interactions have been used extensively in the analyses of the LHC Run 1 results (see [11], and references therein). Limitations in the applicability of an EFT approach to the interpretation of LHC data led to a change of framework for the interpretation of results from the LHC Run 2 [12], which were primarily analysed within the framework of simplified models (see [13], and references therein). By construction, in simplified models for DM the SM is extended by the DM particle and one single mediator particle only. The latter is responsible for the interactions of DM with SM particles, for example, quarks and gluons. When the momentum transferred in a proton-proton collision at the LHC is smaller than the mediator mass (which is not a priori true), the mediator can be “integrated out” and simplified models converge to EFTs. Compared to an EFT approach, the use of simplified models allows for a more complex analysis of the LHC data, especially in the study of processes which involve the mediator explicitly, as in the case of final states including hadronic jets produced by the decay of the mediator into a quark pair. While simplified models are generically not UV complete, and their applicability is subject to constraints from unitarity and anomaly cancellation, e.g. [14, 15], they provide a good compromise between simplicity and completeness.

While processes that directly involve the DM particle are obviously important to reconstruct DM mass and coupling constants at the LHC and have therefore been studied, e.g., in monojet searches [16, 17], events involving the mediator particle alone can be used to obtain important information on the underlying DM model as well [18]. In this context, dijet searches play a special role [19]. Within the framework of simplified models for DM-quark interactions, neutral mediators can be resonantly produced during a proton–proton collision and then decay into a pair of quarks. Due to hadronization these will be seen by a detector as a pair of hadronic jets. The analysis of dijet events at ATLAS and CMS has been one of the main channels in the search for new physics at the LHC. Recently, ATLAS has published results for generic (high-mass) dijet searches from 37 fb$^{-1}$ of data collected during 2015 and 2016 [20], whereas CMS has presented results for 36 fb$^{-1}$ of data from the 2016 dataset [21, 22], as well as preliminary results for 78 fb$^{-1}$ of data from the combined 2016 and 2017 datasets [23]. In these studies, both collaborations have presented results only for a subset of all possible simplified models, mainly focusing on a vector mediator and fermionic DM.
Direct detection experiments and LHC are complementary since they explore physical processes occurring at different energy scales, see e.g. Ref. [24]. Nevertheless, there are aspects of this complementarity which are still not fully understood, or exploited. For example, what is the impact of a signal at XENONnT on present and future searches for DM and new physics in general at the LHC? In this work we investigate the impact of a XENONnT signal on the interpretation of current dijet searches at the LHC, and on the prospects for dijet signal discovery at the HL-LHC. As a theoretical framework, we use a general class of simplified models where DM can have spin 0, 1/2 or 1, and interacts with quarks through the exchange of scalar, pseudo-scalar, vector, or pseudo-vector mediators. This study extends our previous work [25], where we focused on the impact of a XENONnT signal on monojet searches at the LHC. We find that exclusion limits on mediator parameters from dijet searches at the LHC are significantly affected by a signal at XENONnT, and that about 100 signal events at XENONnT dramatically narrow the region in the parameter space of simplified models where a dijet signal can be discovered at 5σ C.L.

The remaining part of this article is organised as follows. In Secs. 2 and 3 we review theoretical framework and statistical methods. Our analysis of the impact of a XENONnT signal on LHC dijet searches is illustrated in Sec. 4. We conclude in Sec. 5 and list useful equations in Appendix A.

2 Theoretical framework

The theoretical framework used in this work consists of a set of simplified models where DM can have spin 0, 1/2 or 1, and interacts with quarks through the exchange of scalar, pseudo-scalar, vector or pseudo-vector mediators [26]. For the purposes of this work, we classify mediators by their interactions to quarks $q$: A spin 0 mediator $\phi$ coupling to the scalar bilinear $\bar{q}q$ is referred to as a scalar mediator, while we refer to the mediator as a pseudo-scalar if $\phi$ couples to $\bar{q}\gamma_5q$. Likewise, we refer to a spin 1 mediator $Z_\mu$ as a vector mediator if $Z_\mu$ couples to the vector bilinear $\bar{q}\gamma_\mu q$, and as a pseudo-vector mediator if $Z_\mu$ couples to $\bar{q}\gamma_\mu\gamma_5q$.

By construction, each simplified model is specified by the spins of the DM particle and mediator, the Lorentz structure of the corresponding interaction vertices, and four free parameters. For each simplified model, the four free parameters are the DM particle and mediator masses, $m_{DM}$ and $m_{med}$, respectively, and two coupling constants: one for the DM-DM-mediator vertex, $g_{DM}$, and the second one for the quark-quark-mediator vertex, $g_q$. This set of simplified models was initially introduced in the context of DM direct detection [26], and later applied to LHC monojet analyses [25] and DM relic density calculations [27]. For completeness, we list the associated Lagrangians in Appendix A.

For each simplified model in Appendix A, we simulate dijet signals at the LHC, and calculate the corresponding dijet invariant mass spectrum, by using the chain of numerical programs
\[ \mathcal{O}_1 = \mathds{1}_X \mathds{1}_N \]
\[ \mathcal{O}_3 = i \hat{S}_N \cdot \left( \frac{\vec{q}_N}{m_N} \times \hat{v}_\perp \right) \mathds{1}_X \]
\[ \mathcal{O}_5 = i \hat{S}_N \cdot \left( \frac{\vec{q}_N}{m_N} \times \hat{v}_\perp \right) \mathds{1}_N \]
\[ \mathcal{O}_6 = \left( \hat{S}_N \cdot \frac{\vec{q}_N}{m_N} \right) \left( \hat{S}_N \cdot \frac{\vec{a}_N}{m_N} \right) \]
\[ \mathcal{O}_7 = \hat{S}_N \cdot \hat{v}_\perp \mathds{1}_X \]
\[ \mathcal{O}_8 = \hat{S}_N \cdot \hat{v}_\perp \mathds{1}_N \]
\[ \mathcal{O}_9 = i \hat{S}_N \cdot \left( \hat{S}_N \times \frac{\vec{a}_N}{m_N} \right) \]
\[ \mathcal{O}_{10} = i \hat{S}_N \cdot \frac{\vec{a}_N}{m_N} \mathds{1}_X \]
\[ \mathcal{O}_{11} = i \hat{S}_N \cdot \frac{\vec{a}_N}{m_N} \mathds{1}_N \]
\[ \mathcal{O}_{12} = \hat{S}_N \cdot \frac{\vec{a}_N}{m_N} \]
\[ \mathcal{O}_{13} = i \left( \hat{S}_N \cdot \hat{v}_\perp \right) \left( \hat{S}_N \cdot \frac{\vec{a}_N}{m_N} \right) \]
\[ \mathcal{O}_{14} = i \left( \hat{S}_N \cdot \frac{\vec{a}_N}{m_N} \right) \left( \hat{S}_N \cdot \hat{v}_\perp \right) \]
\[ \mathcal{O}_{15} = - \left( \hat{S}_N \cdot \frac{\vec{a}_N}{m_N} \right) \left( \hat{S}_N \times \hat{v}_\perp \right) \cdot \frac{\vec{a}_N}{m_N} \]
\[ \mathcal{O}_{17} = i \hat{S}_N \cdot \frac{\vec{a}_N}{m_N} \mathds{1}_N \]
\[ \mathcal{O}_{18} = i \mathds{1}_N \cdot \mathds{1}_N \]

Table 1. Quantum mechanical operators defining the non-relativistic effective theory of DM-nucleon interactions [28]. The operators are expressed in terms of the basic invariants under Galilean transformations: the momentum transfer, \( \hat{q} \), the transverse relative velocity operator \( \hat{v}_\perp \), the nucleon and DM spin operators, denoted by \( \hat{S}_N \) and \( \hat{S}_\chi \), respectively, and the identities in the nucleon and DM spin spaces, \( \mathds{1}_X \) and \( \mathds{1}_N \). All operators have the same mass dimension, and \( m_N \) is the nucleon mass. Standard spin-independent and spin-dependent interactions correspond to the operators \( \mathcal{O}_1 \) and \( \mathcal{O}_4 \), respectively, while \( \mathcal{S} \) is a symmetric combination of spin 1 polarisation vectors [26]. The operators \( \mathcal{O}_{17} \) and \( \mathcal{O}_{18} \) can only arise for spin 1 DM. Following [28], here we do not consider the interaction operators \( \mathcal{O}_2 \) and \( \mathcal{O}_{16} \): the former is quadratic in \( \hat{v}_\perp \) (and the effective theory expansion in [28] is truncated at linear order in \( \hat{v}_\perp \) and second order in \( \hat{q} \)) and the latter is a linear combination of \( \mathcal{O}_{12} \) and \( \mathcal{O}_{15} \).

\textbf{WHIZARD}  \\
\downarrow  \\
\textbf{pythia8}  \\
\downarrow  \\
\textbf{Delphes3 (+FASTJET)}  \\
\downarrow  \\
Custom C++/ROOT code for analysis.

We use WHIZARD [29, 30] with model files implementing the simplified models to generate the hard processes

\[ p, p \to \text{Mediator} \to \eta q + X \]  

(2.1)

where \( X \) stands for additional SM particles and \( \eta q \) can be any pair of quarks with the same flavor. We use parton distribution functions from the CT14lo set as obtained from LHAPDF6 [31]. pythia8 [32] is used for showering and hadronization, Delphes3 [33] for (CMS) detector simulation, and FASTJET [34] for jet reconstruction. We use our own C++ code and ROOT [35] libraries to analyse the signal. We discard events where one (or both) of the leading jets deposit more than 90% of their respective total calorimetric energy in the electromagnetic calorimeter. See [36] for a more detailed discussion of our collider simulations.
| Spin 0 DM | Coeff. | Scalar med. | Vector med. |
|----------|--------|-------------|-------------|
| $c_1$    | $h_N^N g_1 \frac{M_G}{M_Φ}$ | $-2 h_N^N g_1 \frac{M_G}{M_Φ}$ |
| $c_7$    | $4 h_N^N g_4 \frac{M_G}{M_Φ}$ |
| $c_{10}$ | $h_N^N g_1 \frac{M_G}{M_Φ}$ |

| spin 1/2 DM | Coeff. | Scalar med. | Vector med. |
|-------------|--------|-------------|-------------|
| $c_1$       | $h_N^N \lambda_1 \frac{m_N}{m_x}$ | $- h_N^N \lambda_1 \frac{M_G}{M_Φ}$ |
| $c_4$       | $4 h_N^N \lambda_4 \frac{M_G}{M_Φ}$ |
| $c_6$       | $2 h_N^N \lambda_2 \frac{m_N}{m_x}$ |
| $c_7$       | $2 h_N^N \lambda_3 \frac{M_G}{M_Φ}$ |
| $c_8$       | $-2 h_N^N \lambda_4 \frac{M_G}{M_Φ}$ |
| $c_9$       | $-2 h_N^N \lambda_3 \frac{m_N}{m_x} - 2 h_N^N \lambda_4 \frac{M_G}{M_Φ}$ |
| $c_{10}$    | $h_N^N \lambda_1 \frac{M_G}{M_Φ}$ |
| $c_{11}$    | $- h_N^N \lambda_2 \frac{m_N}{m_x}$ |

| Spin 1 DM | Coeff. | Scalar med. | Vector med. |
|-----------|--------|-------------|-------------|
| $c_1$     | $b_1 h_N^N \frac{M_G}{M_Φ}$ | $- q^2 b_1 h_N^N \frac{M_G}{M_Φ}$ |
| $c_4$     | $- 4 h_N^N \text{Re}(\theta_7) \frac{M_G}{M_Φ}$ + $q^2 \frac{M_G}{M_Φ} h_N^N \text{Im}(\theta_6)$ |
| $c_5$     | $- \frac{m_N}{m_x} h_N^N \text{Im}(\theta_6) \frac{M_G}{M_Φ}$ |
| $c_6$     | $- \frac{m_N}{m_x} h_N^N \text{Im}(\theta_6) \frac{M_G}{M_Φ}$ |
| $c_7$     | $4 h_N^N b_6 \frac{M_G}{M_Φ}$ |
| $c_8$     | $2 h_N^N \text{Re}(\theta_7) \frac{M_G}{M_Φ}$ |
| $c_9$     | $- 2 \frac{m_N}{m_x} h_N^N \text{Im}(\theta_6) \frac{M_G}{M_Φ}$ + $2 h_N^N \text{Re}(\theta_7) \frac{M_G}{M_Φ}$ |
| $c_{10}$  | $b_1 h_N^N \frac{M_G}{M_Φ}$ |
| $c_{11}$  | $- \frac{m_N}{m_x} h_N^N \text{Im}(\theta_6) \frac{M_G}{M_Φ}$ |
| $c_{14}$  | $2 \frac{m_N}{m_x} h_N^N \text{Im}(\theta_6) \frac{M_G}{M_Φ}$ |

**Table 2.** Relation between the coupling constants of non-relativistic operators from Tab. 1 (in the proton/neutron basis) and simplified models in this study (see Appendix A and [25, 26] for their Lagrangians). For simplicity, in the second column we omit the index $N$. In the case of spin 1 DM, we do not consider operators that depend on the symmetric combination of polarisation vectors denoted by $S$ in Tab. 1.
For each simplified model in Appendix A, we are also interested in the rate of DM-nucleus scattering events at XENONnT. The expected rate per unit detector mass can be written as

\[
\frac{dR}{dE_R} = \sum_T \xi_T \frac{\rho_X}{m_X m_T} \int_{|v| \geq v_{\text{min}}} d^3v |v| f(v) \frac{d\sigma_T}{dE_R (|v|^2, E_R)},
\]

where \(v_{\text{min}}\) is the minimum DM speed to deposit an energy \(E_R\) in the detector, \(d\sigma_T/dE_R\) is the differential cross section for DM-nucleus scattering, \(\rho_X\) is the local DM density, and \(f(v)\) is the local DM velocity distribution in the detector rest frame. In the sum in Eq. (2.2), we consider the seven most abundant xenon isotopes. Their mass fraction is denoted here by \(\xi_T\).

In order to calculate the expected rate of DM-nucleus scattering events at XENONnT for the simplified models in Appendix A, we proceed as follows. First, we analytically calculate the amplitude for DM scattering on free nucleons as described in detail in \([25, 37–42]\). From these amplitudes, we extract the coupling constants for DM-nucleon interactions, which are related to the ones in the Lagrangians in Appendix A as illustrated in Tab. 2. We then use the coupling constants in Tab. 2 as an input for the package DMFormFactor \([43]\), from which we extract the rate of DM-nucleus scattering events at XENONnT as an output. The result of this calculation depends on the local DM density and velocity distribution. For the DM velocity distribution in the detector rest frame, we assume a Maxwellian velocity distribution with a circular speed of 220 km s\(^{-1}\) for the local standard of rest, and a galactic escape velocity of 544 km s\(^{-1}\). Finally, for the local DM density we adopt the value 0.4 GeV cm\(^{-3}\).

Assuming that the mediator mass is larger than the typical momentum transfer in DM-nucleus scattering events, Tab. 2 relates the Lagrangians in Appendix A to the quantum mechanical operators defining the non-relativistic effective theory of DM-nucleon interactions \([28, 44, 45]\). Interaction operators are denoted by \(\hat{O}_i^{(N)}\), or just \(\hat{O}_i\) for simplicity, and listed in Tab. 1 for completeness. They are expressed in terms of the basic invariants under Galilean transformations and Hermitian conjugation, namely: the momentum transfer operator, \(\hat{q}\), the transverse relative velocity operator \(\hat{v}^\perp\), the nucleon and DM spin operators, \(\hat{S}_N\) and \(\hat{S}_\chi\), respectively, and the identities in the nucleon and DM spin spaces, \(\mathbb{1}_N\) and \(\mathbb{1}_\chi\). In Tab. 1, \(m_N\) is the nucleon mass, and all interaction operators have the same mass dimension. Within this notation, standard spin-independent and spin-dependent interactions correspond to the operators \(\hat{O}_1\) and \(\hat{O}_4\), respectively. The operators \(\hat{O}_2\) and \(\hat{O}_{16}\) do not appear in Tab. 1 for the following reasons: the former is quadratic in \(\hat{v}^\perp\) (while the effective theory expansion in \([28]\) is truncated at second order in \(\hat{q}\) and at linear order in \(\hat{v}^\perp\) ) and the latter is not independent, being a linear combination of the interaction operators \(\hat{O}_{12}\) and \(\hat{O}_{15}\). Finally, the operators \(\hat{O}_{17}\) and \(\hat{O}_{18}\) in Tab. 1 can only arise for spin 1 DM, and \(\mathcal{S}\) is a symmetric combination of spin 1 polarisation vectors \([26]\). In terms of the interaction operators \(\hat{O}_i\), the most general Hamiltonian for non-relativistic DM-nucleon interactions, \(\mathcal{H}\), reads as follows

\[
\mathcal{H} = \sum_{N=p,n} \sum_i c_i^{(N)} \hat{O}_i^{(N)},
\]
where $c_i^{(p)}$ and $c_i^{(n)}$ are the coupling constants for protons and neutrons, respectively. For example, the simplified model characterised by fermionic DM and a vector mediator of mass $M_G$ that couples to DM with coupling constant $g_{DM} = \lambda_3$ and to quarks with (a universal) coupling constant $g_q = h_3$ generates the operator $\hat{O}_1$ in the non-relativistic limit. In this case, $c_i^{(N)} = h_3^{\lambda_3/M_G^2}$, where the nucleon-level and quark-level coupling constants, $h_3^{\lambda_3}$ and $h_3$, are related by $h_3^{\lambda_3} = 3h_3$. For expressions relating $h_i^{(N)}$ to $h_i$, with $i = 1, 2, 4$, we refer to [26]. In the non-relativistic limit, some of the simplified models in Appendix A generate a linear combination of operators in Tab. 1 (see Tab. 2). However, for $m_{DM} = 50$ GeV (the benchmark value used in our calculation), we find that it is always possible to identify a leading operator among those generated from a given simplified model in the non-relativistic limit.

3 Statistical methods

We compute exclusion limits and discovery regions (or sensitivity projections) using the profile likelihood ratio method [46]. In the former case, we compare background plus signal hypothesis, $H_1$, with the background only hypothesis, $H_0$, computing the significance with which a point in parameter space can be excluded. In the latter case, we test the null hypothesis $H_0$ against the alternative $H_1$, computing the significance with which a point in parameter space can be observed. In both cases, we obtain the significance, $Z$, from a profile likelihood ratio $\lambda$ and the test statistic $q = -2\ln \lambda$, using standard asymptotic formulae from [46]

$$Z \simeq \sqrt{q}.$$  

(3.1)

The significance is also related to the $p$-value, i.e. $Z = \Phi^{-1}(1 - p)$, where $\Phi^{-1}$ is the quantile of a Gaussian distribution with mean 0 and variance 1. For example, standard 95% confidence level (C.L.) exclusion limits correspond to a $p$-value of 0.05 and a significance of 1.64.

The exact form of the profile likelihood ratio depends on whether we calculate discovery regions or exclusion limits. For exclusion limits, the profile likelihood ratio takes the following form

$$\lambda = \frac{\mathcal{L}(s, \hat{\theta})}{\mathcal{L}(0, \hat{\theta})},$$

(3.2)

where the likelihood function, $\mathcal{L}$, is defined below. For discovery regions, the profile likelihood ratio is given by

$$\lambda = \frac{\mathcal{L}(0, \hat{\theta})}{\mathcal{L}(s, \theta)}.$$  

(3.3)

Here, the likelihood function for finding a given signal $s = \{s_1, \ldots, s_N\}$ over a background $b = \{b_1, \ldots, b_N\}$ for a dataset $n = (n_1, \ldots, n_N)$ is defined as the product of $N$ Poisson distributions

$$\mathcal{L}(s, \theta) = \prod_{i=1}^{N} \frac{(s_i + b_i(\theta))^{n_i}}{n_i!} e^{-[s_i + b_i(\theta)]},$$

(3.4)
where $N$ is the number of bins in the dijet invariant mass, $s_i$ the number of signal events in the $i$-th bin, $b_i$ the number of background events in the same bin, and $\theta$ a set of nuisance parameters, i.e. the background fit parameters from [47], in our case. In the definition(s) of $\lambda$, $\hat{\theta}$ ($\tilde{\theta}$) is the set of nuisance parameters maximizing the likelihood function for the given signal $s$ (0). Maximising $\mathcal{L}$ with respect to $\theta$ to find $\hat{\theta}$ or $\tilde{\theta}$ at each point in parameter space, we exclude a window around the mediator mass in the dijet invariant mass spectrum. See [36] for further details.

Computing exclusion limits, we evaluate $Z$ for $n_i = n_i^{\text{CMS}}$, where $n_i^{\text{CMS}}$ is the number of observed dijet events at CMS in the $i$-th dijet mass bin. Computing $Z$ for discovery regions, we use the dataset $n_i = b_i(\theta_{\text{bf}}) + s_i$, where $\theta_{\text{bf}}$ is the value of $\theta$ that maximises $L(0, \theta)$ for $n_i = n_i^{\text{CMS}}$.

4 Analysis

In this section we investigate the impact of a XENONnT signal on the interpretation of current dijet searches at the LHC, and on the prospects for dijet signal discovery at the HL-LHC.

Let us start by describing our assumptions about the hypothesised XENONnT signal. We assume that XENONnT with an exposure of 20 ton $\times$ year has detected 150 nuclear recoil events due to DM-nucleus scattering. Roughly, this number of signal events corresponds to DM models lying just below current XENON1T limits. We consider an idealised version of the XENONnT detector with infinite energy resolution, an energy threshold of 5 keV, and 100% detector efficiency. To obtain the number of expected signal events at XENONnT, we follow [26] and integrate the differential rate of nuclear recoil events in Eq. (2.2) in the 5 to 45 keV range using DMFormFactor [43]. The total number of events is then obtained by multiplying the result by a 20 ton $\times$ year exposure. Direct detection experiments are not sensitive to the individual parameters of the simplified models, but only to the DM mass $m_{DM}$ and the effective mediator mass

$$M_{\text{eff}} \equiv \frac{m_{\text{med}}}{(g_q/0.1)(g_{DM}/0.1)}.$$  

(4.1)

For the simplified models in Appendix A, Tab. 3 shows the values of $M_{\text{eff}}$ required to produce 150 signal events at an idealised version of XENONnT. Notice that a signal at XENONnT would constrain $M_{\text{eff}}$ univocally, with an associated uncertainty that would be negligible compared to the uncertainties affecting our dijet signal prediction. Furthermore, experimental errors on the reconstructed value of $m_{DM}$ are also expected to be negligible in this setup, and we therefore set $m_{DM}$ to its benchmark value, i.e. $m_{DM} = 50$ GeV. In addition to the constraints on $M_{\text{eff}}$ and $m_{DM}$ from the detection of 150 signal events at XENONnT, we also require perturbative couplings $|g_{DM}| < \sqrt{4\pi}$ and $|g_q| < \sqrt{4\pi}$. Finally, we assume universal quark couplings $g_u = g_d = g_s = g_c = g_b = g_t \equiv g_q$, and negligible coupling of the mediator to leptons, i.e. $g_\ell \simeq 0$, in agreement with current searches for resonances in dilepton final states at the LHC [48]. Having described our assumptions about the hypothesised XENONnT signal, we now investigate its impact on the interpretation of
Table 3. Benchmark points producing 150 signal events in an idealised version of XENONnT for $m_\chi = 50$ GeV [25]. Consistently with [25], in the case of spin 1 DM we do not consider the contribution to $M_{\text{eff}}$ from effective operators that depend on the symmetric combination of polarisation vectors denoted by $S$.

4.1 Impact of a XENONnT signal on LHC dijet exclusion limits

For the benchmark parameters which would give rise to 150 signal events in XENONnT, we calculate 95% C.L. exclusion limits on the mediator’s coupling to quarks, $g_q$, from current searches for resonances in dijet final states at the LHC. Because we fixed the DM mass to $m_{DM} = 50$ GeV (assuming perfect mass reconstruction), our XENONnT analysis described above yields a surface of parameter points in the space spanned by $\{m_{\text{med}}, g_q, g_{DM}\}$, defined by the respective values of $M_{\text{eff}}$ reported in Tab. 3. On this surface, the mediator’s coupling

| Scalar DM Op. | $g_q$ | $g_{DM}$ | $M_{\text{eff}}$ [GeV] |
|---------------|-------|----------|------------------------|
| 1             | $h_1$ | $g_1$    | 14564.484              |
| 1             | $h_3$ | $g_4$    | 10260.217              |
| 7             | $h_4$ | $g_4$    | 4.509                  |
| 10            | $h_2$ | $g_1$    | 10.706                 |

| Fermionic DM Op. | $g_q$ | $g_{DM}$ | $M_{\text{eff}}$ [GeV] |
|------------------|-------|----------|------------------------|
| 1                | $h_1$ | $\lambda_1$ | 14564.484              |
| 1                | $h_3$ | $\lambda_3$ | 7255.068               |
| 4                | $h_4$ | $\lambda_4$ | 147.354                |
| 6                | $h_2$ | $\lambda_2$ | 0.286                  |
| 7                | $h_4$ | $\lambda_3$ | 3.188                  |
| 8                | $h_3$ | $\lambda_4$ | 225.159                |
| 10               | $h_2$ | $\lambda_1$ | 10.706                 |
| 11               | $h_1$ | $\lambda_2$ | 351.589                |

| Vector DM Op. | $g_q$ | $g_{DM}$ | $M_{\text{eff}}$ [GeV] |
|---------------|-------|----------|------------------------|
| 1             | $h_1$ | $b_1$    | 14564.484              |
| 1             | $h_3$ | $b_5$    | 10260.216              |
| 4             | $h_4$ | $\text{Re}(b_7)$ | 188.302              |
| 5             | $h_3$ | $\text{Im}(b_6)$ | 6.946                   |
| 7             | $h_4$ | $b_5$    | 4.509                  |
| 8             | $h_3$ | $\text{Re}(b_7)$ | 287.728              |
| 9             | $h_4$ | $\text{Im}(b_6)$ | 3.674                   |
| 10            | $h_2$ | $b_1$    | 10.706                 |
| 11            | $h_3$ | $\text{Im}(b_7)$ | 223.794              |
| 14            | $h_4$ | $\text{Im}(b_7)$ | 0.201                   |
Figure 1. 95% C.L. exclusion limits on $g_q$ from the null result of present searches for narrow resonances in dijet final states at the LHC obtained by setting $m_{DM} = 50$ GeV and $g_{DM}$ to the value required by the detection of 150 signal events at XENONnT (coloured lines). Exclusion limits are presented for simplified models with a scalar mediator (a) and for models with a vector mediator (b). In both panels we used data from a CMS search for narrow resonances in final states involving a dijet corresponding to an integrated luminosity of 36 fb$^{-1}$. For models $(h_1, \lambda_1)$ (a) and $(h_3, \lambda_3)$ (b), the figure also shows 95% C.L. exclusion limits obtained by setting $g_{DM}$ to XENONnT-independent values (grey lines).

to DM, $g_{DM}$, is a function of $m_{med}$ and $g_q$. Geometrically, the function $g_{DM} = g_{DM}(m_{med}, g_q)$ can be obtained by projecting the surface defined by $M_{eff}$ to the $m_{med} - g_q$ plane. In practice, for each benchmark point in Tab. 3 we obtain $g_{DM}$ by solving Eq. (4.1) for $g_{DM}$ at each point in the $(m_{med}, g_q)$ plane.

To calculate the 95% C.L. exclusion limits on $g_q(m_{med})$ arising from resonant dijet searches at the LHC and a signal at XENONnT, we use the profile likelihood ratio method outlined in Sec. 3. For each simplified model, we simulate the corresponding dijet invariant mass spectrum using the chain of numerical programmes described in Sec. 2 on a grid in $(m_{med}, g_q)$, setting $g_{DM}$ to the value obtained from the $M_{eff}$ constraint at each point and fixing $m_{DM} = 50$ GeV. We then integrate the simulated dijet mass spectrum to obtain the number of expected dijet events, $s_i$, in bins of dijet invariant mass labeled by the integer $i$ and of variable width, as in the high-mass search for narrow resonances in dijet final states performed by CMS [47]. Following [47], we assume an integrated luminosity of 36 fb$^{-1}$, a centre-of-mass energy of $\sqrt{s} = 13$ TeV, and focus on the 1.6 – 3.9 TeV range for the dijet invariant mass.

Fig. 1 shows the impact that the detection of 150 signal events at XENONnT would have on the 95% C.L. exclusion limits on $g_q$ from the null result of present searches for narrow resonances in dijet final states at the LHC. The left panel refers to simplified models with
scalar or pseudo-scalar mediators, whereas the right panel corresponds to simplified models
with vector or pseudo-vector mediators. In Fig. 1 we report exclusion limits only for a subset
of the simplified models and corresponding benchmark points in Tab. 3. For the models not
shown, the benchmark points would correspond to 95% C.L. exclusion limits extending
to regions in parameter space where coupling constants are non-perturbative. Specifically,
this applies to benchmark points where $M_{\text{eff}} \ll m_{\text{med}}$. The models displayed in Fig. 1
are labeled by the corresponding pair of coupling constants. For the model labeled by $(h_1, \lambda_1)$
in the left panel – fermionic DM and scalar mediator – and the model labeled by $(h_3, \lambda_3)$
in the right panel – fermionic DM and vector mediator – we also show 95% C.L. exclusion
limits obtained using values for $g_{\text{DM}}$ which are not related to constraints on $M_{\text{eff}}$
from the detection of 150 signal events at XENONnT (grey curves). For these two models, 150
signal events at XENONnT require $M_{\text{eff}} \gg m_{\text{med}}$, which implies $g_{\text{DM}} \simeq 0$. This explains
why for models $(h_1, \lambda_1)$ and $(h_3, \lambda_3)$ 95% C.L. exclusion limits computed assuming 150
signal events at XENONnT or setting $g_{\text{DM}} = 0$ are close to each other. On the other hand,
while exclusion limits in Fig. 1 depend only indirectly on $g_{\text{DM}}$ via the total mediator decay
width and branching ratio into quarks, a large coupling to DM can significantly reduce the
branching ratio into quarks, and therefore lead to significantly weaker exclusion limits on $g_q$.

4.2 Impact of a XENONnT signal on LHC dijet 5$\sigma$ discovery contours

In this section, we investigate the impact of a XENONnT signal on the prospects for dijet
signal discovery at the HL-LHC. We use the profile likelihood ratio method outlined in Sec. 3
to identify the contours in the $(m_{\text{med}}, g_q)$ plane where simultaneously: 1) a narrow resonance
in dijet final states at the HL-LHC could be discovered with a statistical significance of
5$\sigma$; 2) 150 signal events are expected at XENONnT. As before, we set $m_{\text{DM}}$ to 50 GeV and
extract $g_{\text{DM}}$ from $M_{\text{eff}}$ using the XENONnT input. Concerning dijet signal and background
calculation, as well as our choice of likelihood function and dataset, we proceed as described
in Sec. 3.

Fig. 2 shows the regions in the $(m_{\text{med}}, g_q)$ plane where a dijet signal could be discovered
at the HL-LHC with a statistical significance larger than or equal to 5$\sigma$, and which are at the
same time compatible with the detection of 150 signal events at XENONnT. Regions with
different colours correspond to distinct simplified models in Appendix A. In each region,
the lower boundary $g_q^{\text{min}}(m_{\text{med}})$ is the smallest coupling $g_q$ for which the corresponding
model could be discovered with a significance of 5$\sigma$ in dijet searches at the HL-LHC. The
upper boundary $g_q^{\text{max}}(m_{\text{med}})$ is given by the 95% C.L. exclusion limits from 36 fb$^{-1}$
of data discussed in Sec 4.1, cf. Fig. 1. The left panel corresponds to simplified models with
a scalar or pseudo-scalar mediator, while the right panel refers to models with vector or
pseudo-vector mediators. As in the previous section, models are labeled by a pair of coupling
constants. Models that do not appear in Fig. 2 are not compatible with the simultaneous
discovery of 150 signal events at XENONnT and the 5$\sigma$ detection of a dijet signal at the
HL-LHC.

Interestingly, we find that only a subset of models would actually be compatible with
the simultaneous detection of a signal at XENONnT and at the HL-LHC. Furthermore, we
Figure 2. Regions in the \((m_{\text{med}}, g_q)\) plane where a narrow resonance could be discovered with \(Z \geq 5\sigma\) C.L. in dijet final states at the HL-LHC that are at the same time compatible with the detection of 150 signal events at XENONnT. The lower boundary of the region for each model is the smallest coupling \(g_q(m_{\text{med}})\) for which we would expect a \(5\sigma\) discovery at the HL-LHC, while the upper boundary is given by the current 95\% C.L. exclusion limits from 36 fb\(^{-1}\) of CMS data. The left panel (a) corresponds to simplified models with a scalar mediator, while the right panel (b) refers to models with a vector mediator.

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find that some of the models in Fig. 2 can potentially be distinguished in dijet searches, since the mediator mass can approximately be reconstructed from an analysis of the dijet invariant mass at the HL-LHC. For example, the discovery of a dijet signal for a mediator mass \(m_{\text{med}} \gtrsim 2.5\) TeV would exclude the models \((h_1, \lambda_2)\) and \((h_3, \lambda_4)\), leaving only models with \(\mathcal{O}_1\) as leading non-relativistic operator for DM-nucleon interactions. On the other hand, the discovery of a dijet signal at lower mediator masses would also be compatible with models \((h_1, \lambda_2)\) and \((h_3, \lambda_4)\), which in the non-relativistic limit generate the interaction operators \(\mathcal{O}_{11}\) and \(\mathcal{O}_8\), respectively. Note, that the regions shown in Fig. 2 do not take into account limits on simplified models which can be obtained from the preliminary results published by the CMS collaborations from 78 fb\(^{-1}\) of data [23] presented in [36]. These results rule out all of the regions for the models \((h_1, \lambda_2)\) and \((h_3, \lambda_4)\) shown in Fig. 2 at 95\% C.L. However, a \(5\sigma\) discovery could still be made at the HL-LHC for such models if one loosens the assumption of 150 signal events being produced at XENONnT. Regarding models \((h_3, \text{Re}(b_7))\) and \((h_3, \text{Im}(b_7))\), we did not compute them explicitly. Indeed, for the benchmark values of \(M_{\text{eff}}\) in Tab. 3 for such models, one would expect similar regions as in the \((h_3, \lambda_4)\) scenario. However, the partial width for a mediator decaying via \(b_7\) is enhanced by a factor \(\sim (m_{\text{med}}/m_{\text{DM}})^2\) with respect to a decay via \(\lambda_3\) [27]. Thus, the branching ratio into quarks is suppressed and we do not expect a significant chance to discover such a scenario at the HL-LHC.
Figure 3. The left panel (a) is the same as Fig. 2, but now focusing on the models \((h_1, g_1)\) and \((h_3, \lambda_3)\). Both models generate \(\hat{O}_1\) as a leading non-relativistic operator for DM-nucleon interactions. Discovery contours are partly non-overlapping. In the right panel (b) we show how the regions where dijets searches could make a 5\(\sigma\) discovery in the \((h_1, \lambda_2)\) model change if instead of \(n = 150\) only \(n = 50\) or \(n = 10\) events would be observed at XENONnT.

As an aside comment, we mention here that \(g_q\) could in principle be inferred from the measurement of the mediator decay width, assuming that \(m_{\text{med}}\) and \(M_{\text{eff}}\) are both known. However, this would require a very accurate measurement of the dijet invariant mass spectrum to extract the mediator decay width from data collected at the HL-LHC. Most likely, the number of signal events recorded in the initial stages of the HL-LHC would not suffice to reconstruct the invariant mass spectrum with the precision required to indirectly infer \(g_q\).

In the left panel of Fig. 3 we compare the two models \((h_1, g_1)\) and \((h_3, \lambda_3)\) in more detail. We find that these models predict partly non-overlapping contours in the \((m_{\text{med}}, g_q)\) plane. This is an interesting result, since it shows that models generating \(\hat{O}_1\) as leading non-relativistic operator can in principle be discriminated if a dijet signal is observed at the HL-LHC. This is in contrast with what we found investigating the impact of a XENONnT signal on LHC monojet searches [25]. In that work, we found that models generating \(\hat{O}_1\) as leading operator in the non-relativistic limit are not observable in monojet searches at the LHC if the model parameters are such that 150 signal events would be observed at XENONnT [25].

Let us now investigate the dependence of our results on the number of signal events observed at XENONnT, \(n\). Fig. 3, right panel, shows how discovery regions change if instead of \(n = 150\) signal events, only \(n = 50\) or \(n = 10\) events are observed at XENONnT. As an example, we show results for the model \((h_1, \lambda_2)\). Note that the effective mass reconstructed
from XENONnT scales with the number of events as $M_{\text{eff}} \sim n^{-4}$. For a fixed combination of parameters $(m_{\text{med}}, g_q)$ this implies the scaling $g_{\text{DM}} \sim \sqrt{n}$. Fewer events at XENONnT therefore imply a smaller partial decay width of the mediator into DM, which makes the discovery of a dijet signal more likely. As a consequence, the exclusion limits as well as the $5\sigma$ discovery contour move towards higher masses and smaller quark couplings. Furthermore, the overall region where a discovery at the HL-LHC is possible becomes larger. We expect a similar behaviour for the model $(h_3, \lambda_4)$. For the remaining models shown in Fig. 2, the mediator’s decay width is dominated by the partial width corresponding to decays into quarks. Therefore, a smaller number of events observed at XENONnT implying a smaller coupling $g_{\text{DM}}$ for fixed $g_q$ and $m_{\text{med}}$ would have virtually no impact on the mediator’s decay width and branching ratios. Thus, the regions in which such models could give rise to a $5\sigma$ discovery at the HL-LHC are nearly independent of the number of events observed at XENONnT.

Let us also qualitatively investigate what impact a DM mass different from $m_{\text{DM}} = 50\text{ GeV}$ would have on our results. We restrict our discussion to the models $(h_1, \lambda_2)$ and $(h_3, \lambda_4)$ for which prospects at the HL-LHC depend strongly on the number of events observed at XENONnT, $n$. Naively, one would assume that the dijet production cross section depends on the DM mass via the branching ratios of the mediator. However, the dependence on the mediator mass, $m_{\text{med}}$, is much stronger, yielding virtually unchanged discovery regions when for example assuming $m_{\text{DM}} = 100\text{ GeV}$ instead of $m_{\text{med}} = 50\text{ GeV}$. On the other hand, DM direct detection experiments using xenon as target material can probe the smallest WIMP-nucleon scattering cross sections for DM masses of $m \sim 50\text{ GeV}$. Thus, observing 150 events at XENONnT corresponds to larger WIMP-nucleus cross sections, and in turn smaller $M_{\text{eff}}$, for larger $m_{\text{DM}}$. Since smaller $M_{\text{eff}}$ correspond to larger $g_{\text{DM}}$ for fixed values of $m_{\text{med}}$ and $g_q$, increasing $m_{\text{DM}}$ to values larger than 50 GeV has the opposite effect as a smaller number of events observed at XENONnT discussed above. Similar to that case, the discovery regions presented for the other models considered here are expected to remain approximately unchanged when for example assuming $m_{\text{DM}} = 100\text{ GeV}$ instead of $m_{\text{med}} = 50\text{ GeV}$.

Finally, we would like to stress that the two models $(h_1, \lambda_2)$ and $(h_3, \lambda_4)$ can simultaneously be observed in monojet [25] and dijet searches at the LHC (see Fig. 2). For model $(h_1, \lambda_2)$, a dijet signal is only observable close to $m_{\text{med}} \approx 2\text{ TeV}$, $g_q \approx 0.3$ and $g_{\text{DM}} \approx 1$, as it can be inferred from Fig. 2. Interestingly, these models also generate non-relativistic operators for DM-nucleon interactions, $\hat{O}_{11}$ and $\hat{O}_8$, respectively, which can statistically be discriminated from an analysis of the associated nuclear recoil energy spectra [25]. These models can therefore be very effectively constrained from a combined analysis of LHC and XENONnT data.

5 Conclusion

In this work we have investigated the impact that a signal at XENONnT would have on the interpretation of current dijet searches at the LHC, and on the prospects for dijet signal discovery at the High-Luminosity LHC in the framework of simplified models. In the
analysis, we have focused on simplified models where DM can have spin 0, 1/2 or 1, and primarily interacts with quarks through the exchange of scalar, pseudo-scalar, vector, or pseudo-vector mediators.

Assessing the impact of a XENONnT signal on the interpretation of current dijet searches at the LHC, we have calculated 95% C.L. exclusion limits on the coupling constant associated with the mediator-quark-quark vertex, $g_q$, as a function of the mediator mass, $m_{\text{med}}$, from the null result of current searches for resonances in dijet final states at the LHC. We have performed this calculation for the simplified models described above (and in greater detail in Appendix A), setting $m_{\text{DM}}$ to the benchmark value of 50 GeV, and taking into account the constraint on the effective mediator mass $M_{\text{eff}}$ (defined in Eq. (4.1)) arising from the detection of 150 signal events at XENONnT. The 95% C.L. exclusion limits on $g_q$ presented here have been calculated using the standard profile likelihood method [46]. We have found that for models for which a XENONnT signal implies $M_{\text{eff}} \ll m_{\text{med}}$ (in the range of $m_{\text{med}}$ values that we have considered), 95% C.L. exclusion limits extend to regions in parameter space where coupling constants are non-perturbative, and therefore become trivial, because of DM detection at XENONnT. At the same time, we have found that models for which 150 signal events at XENONnT require $M_{\text{eff}} \gg m_{\text{med}}$ are characterised by the constraint $g_{\text{DM}} \simeq 0$. In general, we have found that while exclusion limits in the $(m_{\text{med}}, g_q)$ plane depend only indirectly on $g_{\text{DM}}$ via the total mediator decay width and branching ratio into quarks, a large coupling to DM (i.e. $g_{\text{DM}} \sim O(1)$) can significantly reduce the branching ratio into quarks, and therefore lead to significantly weaker exclusion limits on $g_q$.

Assessing the impact of a XENONnT signal on the prospects for dijet signal discovery at the HL-LHC, we have identified the contours in the $(m_{\text{med}}, g_q)$ plane where a narrow resonance could be discovered with a statistical significance of 5$\sigma$ in dijet final states at the HL-LHC, and which are at the same time compatible with the detection of 150 signal events at XENONnT. Interestingly, we have found that only a subset of the simplified models in Appendix A would actually be compatible with the simultaneous detection of a signal at XENONnT and at the HL-LHC. We have also found that some of the models for which the two signals are compatible can potentially be distinguished if the mediator mass is approximately reconstructed from an analysis of the dijet invariant mass at the HL-LHC. Finally, we have found that models generating $\hat{O}_1$ as the leading non-relativistic operator (i.e. canonical spin independent interactions) can in principle be discriminated if a dijet signal is observed at the HL-LHC. Notably, in a previous work [25] we have found that the same models cannot be discriminated by combining a signal at XENONnT with the LHC monojet searches.

Ultimately, our work has explored a new aspect of the well-known complementarity between DM searches at direct detection experiments and at the LHC. The results obtained in this study will be especially useful if DM will be discovered at XENONnT, but the methods illustrated here can in principle be applied to other combinations of DM search experiments.
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A Lagrangians for simplified dark matter models

In this appendix we list the Lagrangians that we considered in the analyses of Sec. 4.1 and Sec. 4.2 [25, 26]. Each Lagrangian listed here describe more than one simplified model. By construction, simplified models are characterised by DM particle mass, \( m_{\text{DM}} \), mediator mass \( m_{\text{med}} \), and just two coupling constants: one for a quark-quark-mediator vertex, \( g_q \), and one for a DM-DM-mediator vertex, \( g_{\text{DM}} \). There are no other interaction vertices in a simplified model. For example, the simplified model associated with fermionic DM of mass \( m_\chi \) and vector mediator of mass \( m_G \), has \( g_q = h_3 \) and \( g_{\text{DM}} = \lambda_3 \) as only coupling constants different from zero. In all numerical applications, we assumed a universal quark-mediator coupling.

A.1 Scalar dark matter \( S \)

Scalar and pseudoscalar mediator \( \phi \):

\[
\mathcal{L}_{S\phi q} = \partial_\mu S^\dagger \partial^\mu S - m_S^2 S^\dagger S - \frac{\lambda_S}{2} (S^\dagger S)^2 \\
+ \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{1}{2} m_\phi^2 \phi^2 - \frac{m_\phi m_{\mu_1}}{3} \phi^3 - \frac{\mu_2}{4} \phi^4 \\
+ i \bar{q} \gamma q - m_\bar{q} \bar{q} q
\]

\[
- g_1 m_S S^\dagger S \phi - \frac{g_2}{2} S^\dagger S \phi^2 - h_1 \bar{q} q \phi - i h_2 \bar{q} \gamma^5 q \phi. \tag{A.1}
\]

Here, \( m_S \) plays the role of \( m_{\text{DM}} \) and \( m_\phi \) that of \( m_{\text{med}} \). \( \lambda_S \) is a dimensionless self-coupling of \( S \) and the \( \mu_i \) are dimensionless self-couplings of the mediator. For the purposes of this work, we set \( \lambda_S = \mu_1 = \mu_2 = 0 \). The \( g_i \) are dimensionless couplings between \( \phi \) and \( S \). The couplings \( h_i \) between the mediators and quarks and the quark mass matrix \( m_q \) should in general be understood as \((6 \times 6)\) matrices and the quarks fields as vectors \( q = (u, d, c, s, t, b) \) in flavor space. Throughout this work we assume universal (diagonal) couplings of the mediators to quarks such that the \( h_i \) can be treated as a single number. The quark mass matrix \( m_q \) can be assumed to be diagonal.
Vector and axial-vector mediator $G_\mu$:  
\[
L_{SG\mu} = \partial_\mu S^\dagger \partial^\mu S - m^2_S S^\dagger S - \frac{\lambda_S}{2} (S^\dagger S)^2 \\
- \frac{1}{4} G^{\mu\nu} G^{\mu\nu} + \frac{1}{2} m^2_G G_\mu G^\mu - \frac{\lambda_G}{4} (G_\mu G^\mu)^2 \\
+ i \bar{q} \gamma_\mu q - m_q \bar{q} q \\
- \frac{g_3}{2} S^\dagger S G_\mu G^\mu - ig_4 (S^\dagger \partial_\mu S - \partial_\mu S^\dagger S) G^\mu \\
- h_3 (\bar{\chi} \gamma_\mu q) G^\mu - h_4 (\bar{\chi} \gamma_\mu \gamma^5 q) G^\mu.
\]  
\[\text{(A.2)}\]

Here, $G_{\mu\nu}$ is the field strength tensor of $G_\mu$, $m_G$ plays the role of $m_{\text{med}}$, and $\lambda_G$ is a dimensionless self-coupling of $G_\mu$ which we set to zero for the purposes of this work. The $g_i$ are the dimensionless couplings of $S$ to $G_\mu$, and the $h_i$ are the couplings of $G_\mu$ to quarks. As before, the $h_i$ are in general $(6 \times 6)$ matrices in flavor space, but can be treated as single numbers for the universal quark coupling assumed here. The remaining parameter are as in Eq. (A.1).

A.2 Fermionic dark matter $\chi$

Scalar and pseudoscalar mediator $\phi$:  
\[
L_{\chi \phi q} = i \bar{\chi} \gamma_\mu \phi - m_\chi \bar{\chi} \phi \\
+ \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{1}{2} m^2_\phi \phi^2 - \frac{m_\phi \mu_1}{3} \phi^3 - \frac{\mu_2}{4} \phi^4 \\
+ i \bar{q} \gamma_\mu q - m_q \bar{q} q \\
- \lambda_1 \phi \bar{\chi} \phi - i \lambda_2 \phi \bar{\chi} \gamma^5 \chi - h_1 \phi \bar{q} q - i h_2 \phi \gamma^5 \bar{q} q.
\]  
\[\text{(A.3)}\]

Here, $m_\chi$ plays the role of $m_{\text{DM}}$. The $\lambda_i$ are the dimensionless couplings between $\phi$ and $\chi$. The remaining parameter are as in Eq. (A.1).

Vector and axial-vector mediator $G_\mu$:  
\[
L_{\chi G \mu} = i \bar{\chi} \gamma_\mu \phi - m_\chi \bar{\chi} \phi \\
- \frac{1}{4} G_{\mu\nu} G^{\mu\nu} + \frac{1}{2} m^2_G G_\mu G^\mu \\
+ i \bar{q} \gamma_\mu q - m_q \bar{q} q \\
- \lambda_3 \bar{\chi} \gamma^\mu \chi G_\mu - \lambda_4 \bar{\chi} \gamma^\mu \gamma^5 \chi G_\mu \\
- h_3 (\bar{\chi} \gamma_\mu q) G^\mu - h_4 (\bar{\chi} \gamma_\mu \gamma^5 q) G^\mu.
\]  
\[\text{(A.4)}\]

Beyond the parameters appearing in Eqs. (A.2) and (A.3), the $\lambda_i$ are the dimensionless couplings between $\chi$ and $G_\mu$. 

– 17 –
A.3 Vector dark matter $X_\mu$

Scalar and pseudoscalar mediator $\phi$:

$$\mathcal{L}_{X\phi q} = -\frac{1}{2} X_{\mu\nu} X^{\mu\nu} + m_X^2 X_\mu^+ X^\mu - \frac{\lambda_X}{2} (X_\mu^+ X^{\mu})^2$$

$$+ \frac{1}{2} (\partial_\mu \phi)^2 - \frac{1}{2} m_\phi^2 \phi^2 - \frac{m_\phi}{3} \phi^3 - \frac{\mu_2}{4} \phi^4$$

$$+ i \bar{q} D_\mu \phi - m_\phi q \phi$$

$$- b_1 m_X \phi X_\mu^+ X^\mu - b_2 \phi^2 X_\mu^+ X^\mu$$

$$- h_1 \phi \bar{q} q - i h_2 \phi \bar{q} \gamma^5 q.$$  \hspace{1cm} (A.5)

Here, $X_{\mu\nu}$ is the field strength tensor of $X_\mu$, $m_X$ plays the role of $m_{DM}$, and $\lambda_X$ a dimensionless self-coupling of $X_\mu$ which we set to zero for the purposes of this work. The $b_i$ are dimensionless couplings of $\phi$ to $X_\mu$. The remaining parameters are as in Eq. (A.1).

Vector and axial-vector mediator $G_\mu$:

$$\mathcal{L}_{XGq} = -\frac{1}{2} X_{\mu\nu} X^{\mu\nu} + m_X^2 X_\mu^+ X^\mu - \frac{\lambda_X}{2} (X_\mu^+ X^{\mu})^2$$

$$- \frac{1}{4} G_{\mu\nu} G^{\mu\nu} + \frac{1}{2} m_G^2 G_\mu^2 - \frac{\lambda_G}{4} (G_\mu G^{\mu})^2$$

$$+ i \bar{q} D_\mu \phi - m_\phi^2 \phi q$$

$$- b_3 \frac{1}{2} G_\mu^2 (X_\mu^+ X^{\nu}) - b_4 \frac{1}{2} (G_\mu^\nu) (X_\mu^+ X^{\nu})$$

$$- \left[ i b_5 X_\mu^\nu \partial_\mu X^\nu G^\mu + b_6 X_\mu^\nu \partial^\mu X_\nu G^\nu + b_7 \varepsilon_{\mu\nu\rho\sigma} (X_\mu^\nu \partial^\rho X^\sigma) G^\sigma + \text{h.c.} \right]$$

$$- h_3 G_\mu \bar{q} \gamma^\mu q - h_4 G_\mu \bar{q} \gamma^\mu \gamma^5 q.$$  \hspace{1cm} (A.6)

Here, the $b_i$ are dimensionless couplings between $X_\mu$ and $G_\mu$. The remaining parameters are as in Eqs. (A.2) and (A.5).

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