Neutron EDM constrains direct dark matter detection prospects

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A non-relativistic effective field theory (NREFT) offers a bottom-up framework to classify Dark Matter (DM) – nucleon interactions relevant for scattering at direct detection experiments by organizing the interactions in powers of the momentum transfer $\vec{q}$ and DM velocity $\vec{v}$. This approach generates a number of operators including P-odd and T-odd operators; these can only be generated from a relativistic theory with CP violating interactions. We consider the leading order P-odd, T-odd operators viz. $O_{10}$, $O_{11}$ and $O_{12}$ and compare the constraints on these operators from leading direct detection searches and from the bound on the neutron EDM (nEDM). We perform our analysis using simplified models with charged mediators and compute the loop diagrams contributing to the nEDM. We find that constraints on the DM scattering cross section from the bound on the nEDM are several orders of magnitude stronger than the limits from direct searches, and even well below the neutrino floor for such NREFT operators, for the entire sub-GeV to TeV DM mass range. This indicates that these operators need not be considered when analyzing data from present or future direct dark matter detection experiments.

I. INTRODUCTION

The particle identity of Dark Matter (DM) has not yet been understood despite overwhelming evidence for its gravitational interactions. Dedicated direct searches exploit the possible non-gravitational interactions of DM by looking for the recoil of target nuclei due to elastic scattering and have already probed a large fraction of DM mass and cross-section parameter space \cite{1–10}.

Much of the direct detection search strategy for roughly three decades has been focused on a category of DM candidates known as Weakly Interacting Massive Particles (WIMPs). WIMPs with weak-scale masses, produced in the early universe via freeze-out from the thermal plasma \cite{11}, can easily match the measured DM relic abundance via roughly electroweak-strength couplings to Standard Model (SM) particles. Most analyses assume DM-nucleon interactions to be dominated by only two operators, which describe Spin-Independent (SI) or Spin-Dependent (SD) interactions in the limit of vanishing WIMP velocity. The leading bounds on SI DM-nucleon cross section and SD DM-neutron cross section are currently set by XENON1T \cite{1, 3}; for a 30 GeV WIMP, the bounds are $4 \times 10^{-47}$ cm$^2$ and $6 \times 10^{-42}$ cm$^2$, respectively. The best direct detection bound on the SD DM-proton cross section comes from the PICO-60 experiment \cite{7}; for 25 GeV DM mass cross sections above $2 \times 10^{-41}$ cm$^2$ are excluded.

The standard SI and SD operators are actually the leading (zeroth order) terms of an EFT organized in powers of small expansion parameters such as DM velocity $v$ and three-momentum transfer normalized to the nucleon mass $\frac{|\vec{q}|}{m_N}$. Since the DM velocity is relatively small in the solar neighborhood, $v/c \sim O(10^{-3})$, the momentum exchange is limited to $|\vec{q}| \lesssim \min[v m_\chi, v m_A] \lesssim O(100 \text{ MeV})$; here $m_\chi$ is the WIMP mass and $m_A$ is the mass of the target nucleus. This is sizable on nuclear physics scales but certainly well below the electroweak scale as well as the range of WIMP masses to which these experiments are sensitive. Keeping only the leading order terms therefore a priori seems justified.

Nevertheless a non-relativistic effective field theory (NREFT) was developed which retains NLO and NNLO terms by modeling the nucleus after taking into account the finite spatial extent of its charge and spin densities \cite{12–15}. Working up to second order then generates a total of 14 operators, some of which lead to novel nuclear responses motivating the inclusion of angular-momentum dependent (LD) as well as spin and angular-momentum dependent (LSD) interactions, along with SI and SD. Moreover, these operators can interfere with each other and can lead to appreciable differences from the usual exponentially falling recoil energy spectra. It has recently been pointed out that this affects the optimization of the experimental search strategy, in particular by allowing events with larger recoil energy \cite{16}. Considerable effort has also been devoted to analyzing published data in this framework \cite{17–20}.

The NREFT contains a total of 28 free parameters if operators involving neutrons and protons are treated separately. This large parameter space can be difficult to probe in its entirety. There have been a number of global analyses of
the full multi-dimensional parameter space defining the NREFT using constraints from ongoing and planned direct
detection experiments [21–28]. One conclusion is that direct searches constrain momentum or velocity suppressed
operators that are odd under P (parity) and T (time reversal) transformations about as strongly as zeroth order SD
interactions, but the former have been less explored by both theorists and experimenters. This motivates analyses of
particle physics aspects of models that can give rise to these operators in the non-relativistic limit.

From a particle physics perspective, the CPT theorem implies that all P- and T-odd NREFT operators must
necessarily arise from a CP violating theory, where C refers to charge conjugation. However, CP violation (CPV)
is highly constrained experimentally [29] and therefore bounds on CPV observables can be used to constrain such
NREFT operators. The neutron EDM (nEDM) is one such observable which probes flavor diagonal CP violation and
provides one of the most stringent bounds on CPV in extensions of the SM. Ongoing nEDM experiments are sensitive
to a possible signal which is many orders of magnitude larger than the SM prediction [29, 30]. The current constraint
is
$$|d_n| < 2.9 \times 10^{-26} \text{ e} \cdot \text{cm} \quad (90\% \ C.L.) \, . \tag{1.1}$$

Various works have linked well-motivated extensions of the SM to low energy NREFT operators [31–35]. In par-
cular, Ref. [36] lists a set of simplified models for scalar, spinor and vector DM candidates and derives the full set
of NREFT operators in terms of the parameters for each simplified model. In this letter, we investigate the P- and
T-odd operators in the NREFT with the help of simplified models that respect $SU(3) \times U(1)_Y$ invariance and are
renormalizable, as in [36]. The models in this work extend the SM by a WIMP candidate and a mediator particle.
We show that in these models the Wilson coefficients of the CPV NREFT operators and of the neutron EDM scale
with the same combination of Yukawa couplings. As a result, we find that the constraints from the nEDM on the
scattering cross section are many orders of magnitude stronger for both CPV NREFT operators than the current
XENON1T bound and even lie below the irreducible background from coherent neutrino scattering (“neutrino floor”)
[37–40]. This casts doubt on the importance of these operators for the analysis of future experiments [41] and for the
formalism of NREFT itself.

The rest of this letter is organized as follows. In Section II we introduce the NREFT formalism, the matching
procedure of a relativistic theory to NREFT operators and review the computation of the differential event rate for
elastic WIMP-nucleus scattering. In Section III we consider two simplified models yielding these P- and T-odd CPV
NREFT operators and compute the nEDM at 1-loop order for each case. Section IV contains our main numerical
results, comparing the nEDM and the leading direct detection bound. We conclude in Section V.

II. NREFT FORMALISM

A non-relativistic effective field theory (NREFT) of elastic scattering between DM and nuclei which goes beyond the
strict velocity $v \to 0$ limit [12–15] offers a parametrically richer description than the traditional relativistic four-field
effective operator approach. Since an incoming DM particle striking the target nucleus is fairly slow in the target rest
frame, $v/c \sim O(10^{-3})$, an NREFT is well suited to describe the scattering process. We briefly review this formalism for
completeness.

We begin by discussing the basis of NREFT operators that describe the interactions of a WIMP $\chi$ with a nucleon
$N$ at low relative velocities. Galilean symmetry restricts the operator basis to consist of quantities which are invariant
under Galilean boosts. The three-momentum transfer $\vec{q}$ and relative initial velocity $\vec{v} = \vec{v}_\chi - \vec{v}_N$ satisfy this
criterion, as does any quantity that is computed only from these three-vectors. Since any two relative velocities are
sufficient to describe the scattering, it is helpful to consider orthogonal combinations such that the scalar product of
the two invariants vanishes. The transverse velocity $\vec{v}^\perp \equiv \vec{v} + \frac{\vec{q}}{2m_N}$ is by construction orthogonal to the momentum
transfer $\vec{q}$, where $\mu_N = m_\chi m_N/(m_\chi + m_N)$ is the reduced mass of the WIMP-nucleon system; $\vec{v} \cdot \vec{q} = 0$ follows from
energy and momentum conservation.

Along with the two three-vectors $\vec{q}$ and $\vec{v}^\perp$, the scattering can be fully described by the spins of the target nucleon $\vec{S}_N$ and of the DM particle $\vec{S}_\chi$ (when non-zero), which are also invariant under a Galilean transformation. The
operator basis constructed so far is not so convenient for the construction of an interaction Hamiltonian since it is
not entirely Hermitean; for instance the momentum transfer $\vec{q}$ is anti-Hermitean. Therefore, the most general
non-relativistic interactions for elastic DM–nucleus scattering can be written down as functions of the following four
Hermitean, Galilean invariant quantities:
$$i \vec{q}, \quad \vec{S}_N, \quad \vec{S}_\chi \, . \tag{2.1}$$

A set of linearly independent operators $\mathcal{O}_i$ can be constructed using the building blocks in Eq. (2.1). Table I lists
the set of operators which are obtained when terms up to second order in momentum exchange $\vec{q}$ are retained. The
operators $\mathcal{O}_1$ and $\mathcal{O}_4$ correspond to the traditional SI and SD interaction, respectively.
\[ O_1 = 1 \chi_1 N; \quad O_6 = \left( \frac{q}{m_N} \cdot \bar{S}_N \right) \left( \frac{q}{m_N} \cdot \bar{S}_X \right); \quad O_{10} = i \frac{q}{m_N} \cdot \bar{S}_N; \]
\[ O_3 = i \bar{S}_N \cdot \left( \frac{q}{m_N} \times \vec{v}^\perp \right); \quad O_7 = \bar{S}_N \cdot \vec{v}^\perp; \quad O_{11} = i \frac{q}{m_N} \cdot \bar{S}_N; \quad O_{14} = i (\bar{S}_N \cdot \vec{v}^\perp) \left( \frac{q}{m_N} \cdot \bar{S}_X \right); \]
\[ O_4 = \bar{S}_X \cdot \bar{S}_N; \quad O_8 = \bar{S}_X \cdot \vec{v}^\perp; \quad O_{12} = \bar{S}_X \cdot (\bar{S}_N \times \vec{v}^\perp); \]
\[ O_5 = i \bar{S}_X \cdot \left( \frac{q}{m_N} \times \vec{v}^\perp \right); \quad O_9 = i \bar{S}_X \cdot \left( \frac{q}{m_N} \times \bar{S}_N \right); \quad O_{13} = i (\bar{S}_X \cdot \vec{v}^\perp) \left( \frac{q}{m_N} \cdot \bar{S}_N \right); \]

**TABLE I:** List of operators in the NREFT for elastic WIMP-nucleon scattering. We adopt the conventions of [14] by defining the operators normalized by the nucleon mass \( m_N \) in order to have a dimensionless basis. We omit the invariant \( O_2 = v^1_\perp \) because it is a second order correction to the SI operator \( O_1 \), as well as
\[ O_{15} = - \left( \bar{S}_X \cdot \frac{q}{m_N} \right) \left( (\bar{S}_N \times \vec{v}^\perp) \cdot \frac{q}{m_N} \right) \] since it generates a cross section of order \( v_\perp^0 \), which is \( \mathcal{N}^3 \text{LO} \).

In deriving the NREFT operators, no requirement to obey discrete symmetries such as invariance under spatial parity P and time reversal T transformations has been imposed.\(^1\) Therefore, the operators in Table I can be classified according to their P and T quantum numbers. Since velocities and three momenta are odd under both P and T whereas spin is odd under T but even under P, the following operators are both P-odd and T-odd:
\[ O_{10}, O_{11}, O_{12} \quad \text{P-odd and T-odd}; \quad (2.2) \]

note that the anti-linear transformation T maps a complex number to its complex conjugate.

In order to match some UV-complete or simplified model for DM onto the NREFT, a two-step procedure is used. In the first step one integrates out the heavy mediator field(s), leading to a still relativistic but non-renormalizable description in terms of four-field DM-quark operators. If WIMPs are spin-1/2 fermions the nucleon and DM bilinears can be formed using a basis of 16 linear hermitean matrices \( \Gamma_{\chi,N} \in \{ I_4, i\gamma_5, \gamma^\mu, \gamma^\mu\gamma^5, \sigma^{\mu\nu} \} \). For scalar DM the Lorentz structure appearing in the scalar bilinear can be\(^2\) \( \{ I_4, i\gamma_5, \gamma^\mu, \gamma^\mu\gamma_5 \} \). In either case only combinations are allowed where all Lorentz indices are contracted between the two bilinears. By convention the scalar DM bilinears are multiplied with a factor of the scalar mass \( m_S \) i.e. \( m_S S^I \Gamma_S S \) such that the Wilson coefficients obtained have the same mass dimension as those of the fermionic WIMPs.

In the non-relativistic limit these dimension-6 four field operators composed of a nucleon bilinear and a DM bilinear reduce to \( O_N \) and \( O_X \) respectively; their products can be matched on to NREFT operators \( O_{i,\text{NR}} \) of Table I with Wilson coefficients \( c^N_{\chi} \) (in the isospin basis) of mass dimension GeV\(^{-2} \), with \( O_{i,\text{NR}} \equiv O_i \cdot O_N \).

Any relativistic theory with a DM candidate can thus be matched on to a particular linear combination of the NREFT operators. The rate, differential in the recoil energy \( E_R \), of elastic WIMP–nucleus scattering events per unit time and unit detector mass, \( dR/dE_R \), is then given by
\[ \frac{dR}{dE_R} = N_T \rho_X m_A \int_{v_{\text{min}}}^{v_{\text{max}}} \frac{f(v)}{v} \frac{1}{2j_X + 1} \frac{1}{2j_A + 1} \sum_{\text{spins}} |\mathcal{M}_{\text{sc}}|^2 \, d^3v. \quad (2.3) \]

Here \( N_T \) is the number of target nuclei in the detector, \( \rho_X \) is the local DM density, \( m_X \) is the DM mass, \( m_A \) is the mass of the target nucleus\(^3\), \( f(v) \) is the DM halo velocity distribution in the laboratory frame, \( v_{\text{min}} \) is the minimum velocity required to cause a nuclear recoil \( E_R \), \( v_{\text{max}} \) is related to the galactic DM escape velocity, and \( j_X \) and \( j_A \) are the total spin of the DM and nucleus, respectively. Finally, the matrix element \( \mathcal{M}_{\text{sc}} \) is given by [14]:
\[ \frac{1}{2j_X + 1} \frac{1}{2j_A + 1} \sum_{\text{spins}} |\mathcal{M}_{\text{sc}}|^2 = \sum_{k', \Delta_S, \Sigma' \chi''} R^{NN'}_k (v^2, q^2) W^{NN'}_k (q^2b^2) + \sum_{k', \Delta_S, \Sigma' \chi''} \frac{q^2}{m_N} R^{NN'}_k (v^2, q^2) W^{NN'}_k (q^2b^2). \quad (2.4) \]

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1 Since antiparticles do not occur in the non-relativistic limit, C transformations do not play a role.
2 In general a bilinear with two derivatives can also contribute; however, this does not happen for the models we discuss below.
3 If the detector contains different types of nuclei, a sum over all isotopes is understood in eq.(2.3).
The result (in the isospin basis) has been factorized into WIMP response functions $R_k$ and nuclear response functions $W_k$. Both response functions depend on the momentum transfer. The former also depend on the relative WIMP-nucleon velocity and encode the particle physics of the scattering process, and thus explicitly depend on the product of Wilson coefficients $c^2 \gamma^2$ and can be computed perturbatively; see Ref. [14] for a list of explicit expressions. The latter encode the target-dependent nuclear physics; they are functions of $y = (q\nu/2)^2$, where $b$ is the nuclear size. The sum over $k$ in Eq. (2.4) involves the enlarged set of five independent and two interfering responses: the standard SI $M$, the transverse SD $\Sigma'$, the longitudinal SD $\Sigma''$, the angular momentum dependent (LD) $\Delta$ and the angular momentum and spin dependent (LSD) $\Phi''$ whereas interference of the angular momentum and spin dependent response with the standard SI leads to $\Phi''M$ and the transverse SD and the longitudinal angular momentum response interfere to give rise to $\Delta \Sigma'$. Five of the seven responses are accompanied by a factor of $q^2/m_N^2$, a parameter related to the relative velocities of nucleons bound in the nucleus, further reflecting that the NREFT operators associated with the responses are all suppressed by at least $q^2/m_N^2$.

If the Wilson coefficients in the final non-relativistic effective action are all of the same order of magnitude, the NREFT operators listed in Table I are not equally relevant for scattering. The traditional SI operator $O_1$ then dominates over the rest since it is enhanced by a factor of $A^2$, where $A$ is the nucleon number, and has no velocity suppression. Naively, one would conclude that the other remaining momentum/velocity independent operator $O_4$ would be the dominant contribution to scattering in the absence of $O_1$; however, that is not always the case. Recall that the momentum transfer involved in elastic DM-nucleon scattering can be of the order a hundred MeV, $|q| \sim 100$ MeV. Once the target nucleus can be resolved into its constituent nucleons, there will be contributions which are only suppressed by the normalized momentum transfer $\bar{q}^2/m_N^2$, which can be $O(10^{-2})$, rather than by the squared velocity $v_\perp^2 \sim O(10^{-6})$. Therefore the contribution from momentum transfer suppressed NREFT operators frequently dominates over those from velocity suppressed ones, at least for relatively large recoil energies. As a result, in the absence of $O_1$ the SI P- and T-odd operator $O_{11}$ often dominates scattering instead of $O_1$, if the corresponding Wilson coefficients are of similar size. This explains why numerical scans [21, 22] found fairly stringent constraints on the Wilson coefficient of these operators.

### III. DM SIMPLIFIED MODELS

Simplified models of DM provide a minimal framework to explore its phenomenology across direct, indirect and LHC searches. In these models the SM is augmented with a WIMP candidate and a mediator via which quarks and/or leptons can interact with WIMPs. Although renormalizable and invariant under $SU(3)_C \times SU(1)_{\text{em}}$ gauge transformations, they usually ignore the $SU(2)$ part of the SM gauge group; they are therefore not UV complete theories. However, since simplified models contain a relatively small number of parameters they can be powerful tools to interpret experimental results and explore the complementarity of different searches.

Ref. [36] has constructed a set of simplified models by exhaustively listing different WIMP and mediator spins and matched the resulting relativistic Lagrangians to the set of non-relativistic operators listed in Table I. In these models the WIMP is an $SU(3)_C \times SU(1)_{\text{em}}$ singlet; the interactions are further constrained by the requirement that the WIMP cannot decay.

Here we focus on two simplified models with color-triplet mediators which in the NR limit give rise to the leading order P- and T-odd operators $O_{11}, O_{10}$ and $O_{12}$; $O_{11}$ is spin-independent while $O_{10}$ and $O_{12}$ are spin-dependent. In the following subsections we describe these models and use the experimental upper bound on the neutron EDM to constrain their direct detection prospects.

#### A. Model I

Model I contains a complex spin-0 WIMP $S$ and one or more heavy quark-like mediator(s) $Q$, both of which are odd under a discrete symmetry $Z_2$ to forbid dark matter decay. The SM quarks $q$ can have two new Yukawa-type interactions with the WIMP and the mediator(s), via scalar couplings $y^q_1$ and pseudoscalar couplings $y^q_2$, where the superscript denotes the quark flavor index $q$. The Lagrangian thus reads:

$$
\mathcal{L}_{\text{Model I}} = \mathcal{L}_\text{SM} + \partial_\mu S^\dagger \partial_\mu S - m^2 S^\dagger S - \lambda _S (S^\dagger S)^2 + i \bar{Q}_k \Gamma_Q Q_k - m_{Q_k} \bar{Q}_{Q_k} Q_k - \bar{S} \gamma^\dagger Q_k (y^q_1 + y^q_2 \gamma^5) Q_k - \bar{S} \gamma^\dagger (y^q_1 + y^q_2 \gamma^5) Q_k .
$$

(3.1)

$SU(1)_{\text{em}}$ invariance implies that at least two mediators are required if the WIMP $S$ is to couple to both up- and down-type quarks. Moreover, a mediator $Q_k$ with nonvanishing couplings to quarks of different generations will give
rise to flavor changing neutral current (FCNC) processes at one-loop order. For example, a $Q_k$ coupling to both $d$ and $s$ quarks will contribute to $K^0-\overline{K^0}$ mixing via box diagrams with $Q_k$ and $S$ in the loop. We will therefore assume that each mediator $Q_q$ couples only to a single SM quark $q$ in the mass basis, so that the Yukawa coupling matrices appearing in eq. (3.1) reduce to diagonal matrices.

The matrix element for $s$-channel scattering $S(p_1) + q(p_2) \rightarrow S(k_1) + q(k_2)$ is then given by:

$$
\mathcal{M}_{Sq \rightarrow Sq} = \frac{m_{Q_q}}{m_{Q_q} - m_S} \left[ (|y_1|^2 - |y_2|^2) \, \bar{u}(k_2) \, u(p_2) - 2 \text{Im}(y_1^q y_2^\dagger_q) \, \bar{u}(k_2) \, i \gamma^5 \, u(p_2) \right] + \frac{1}{2} \left( \frac{m_{Q_q}}{m_{Q_q} - m_S} \right) \left[ m_q \, \bar{u}(k_2) \, u(p_2) + \bar{u}(k_2) \, \frac{p_1^\mu + k_1^\mu}{2} \, u(p_2) \right] + \frac{1}{2} \left( \frac{m_{Q_q}}{m_{Q_q} - m_S} \right) 2 \text{Re}(y_1^q y_2^\dagger_q) \, \left[ \bar{u}(k_2) \, \frac{p_1^\mu + k_1^\mu}{2} \, \gamma^5 \, u(p_2) \right].
$$

(3.2)

Here we have neglected the mass of the incoming quark as well as terms of order $|\bar{q}|^2$ in the denominator of the $Q_q$ propagator, but we have kept these terms in the numerator; note that at tree-level only scattering on $u$, $d$ and $s$ quarks contribute to WIMP-nucleon scattering. Moreover, we have used the Dirac equation and 4-momentum conservation to write the matrix element in a form that is symmetric in WIMP momenta. This facilitates the construction of the corresponding relativistic effective Lagrangian, which results when the mediator $Q_q$ is integrated out. Defining the Hermitian derivative on the complex scalars as $iS^\dagger \frac{\partial S}{\partial S} \equiv \frac{i}{2} (S^\dagger \frac{\partial}{\partial S} - S \frac{\partial}{\partial S^\dagger})$, we have:

$$
\mathcal{L}_{\text{eff}}^{\text{Model I}} \supset c_1^{q, \text{dS}} (S^\dagger S) \bar{q} \, \gamma^\mu q + c_0^{q, \text{dS}} (S^\dagger S) \bar{q} \, i \gamma^5 \, \gamma^\mu q + c_1^{q, \text{dS}} (S^\dagger \frac{\partial S}{\partial S} S) \bar{q} \, \gamma^\mu q + c_7^{q, \text{dS}} (S^\dagger \frac{\partial S}{\partial S} S) \bar{q} \, i \gamma^5 \, \gamma^\mu q.
$$

(3.3)

The subscripts on the Wilson coefficients denote the NREFT operator that the corresponding relativistic effective operator reduces to and the superscripts denote the mass dimension of the corresponding relativistic effective operator.

| $S^\dagger \Gamma_S S \bar{N} \Gamma_N N$ | $c_i \, \mathcal{O}_i$ |
|---------------------------------|------------------|
| $c_1^{q, \text{dS}} S^\dagger S \bar{q} \, \gamma^\mu q$ | $\frac{m_{Q_q} \, |y_1|^2 - |y_2|^2|}{m_S \, m_{Q_q} - m_S^2} \, \mathcal{J}^N \, \mathcal{O}_1$ |
| $c_0^{q, \text{dS}} S^\dagger S \bar{q} \, i \gamma^5 \, \gamma^\mu q$ | $\frac{m_{Q_q} \, \text{Im}(y_1^q y_2^\dagger_q)}{m_S \, m_{Q_q} - m_S^2} \, 2 \Delta^N \, \mathcal{O}_{10}$ |
| $c_1^{q, \text{dS}} \left( S^\dagger \frac{\partial S}{\partial S} S \right) \bar{q} \, \gamma^\mu q$ | $\frac{|y_1|^2 + |y_2|^2}{m_{Q_q} - m_S^2} \, N^N \, \mathcal{O}_1$ |
| $c_7^{q, \text{dS}} \left( S^\dagger \frac{\partial S}{\partial S} S \right) \bar{q} \, i \gamma^5 \, \gamma^\mu q$ | $- \frac{\text{Re}(y_1^q y_2^\dagger_q)}{m_{Q_q} - m_S^2} \, 2 \Delta^N \, \mathcal{O}_7$ |

TABLE II: Non-relativistic reduction of effective operators in Model I. $\mathcal{J}^N$, $N^N$, $\Delta^N$ and $\Delta^N$ are coefficients arising due to promoting quark bilinear to nucleon bilinears [42, 43]. We use the values as given in the Appendix of Ref. [36].

Table II contains the matching of the relativistic effective operators to the NREFT operators in terms of the parameters of Model I. The Wilson coefficients for the dimension-5 operators $S^\dagger S \bar{q} \, \gamma^\mu q$ and $S^\dagger S \bar{q} \, i \gamma^5 \, \gamma^\mu q$ have been divided by a factor of $m_S$ such that the expressions for all DM-nucleon cross sections contain the same factor $\frac{m_S}{m_{Q_q}} (c_i^N)^2$ irrespective of the mass dimension of the relativistic operator involved. The operators $(S^\dagger S) (\bar{q} \gamma^\mu q)$ and $i \left( S^\dagger \frac{\partial S}{\partial S} S \right) \bar{q} \, \gamma^\mu q$ both reduce to the leading SI $\mathcal{O}_1$ operator but with different coupling combinations. Ignoring the sub-leading $\mathcal{O}(m_{Q_q}/m_S^2)$ contributions, the contribution $\propto (S^\dagger S) (\bar{q} \gamma^\mu q)$ scales as the difference of the absolute value squared of the Yukawa couplings $|y_1|^2 - |y_2|^2$ whereas $\propto i \left( S^\dagger \frac{\partial S}{\partial S} S \right) \bar{q} \, \gamma^\mu q$ scales as the sum of the absolute value squared of the Yukawa couplings $|y_1|^2 + |y_2|^2$; the latter contribution is suppressed by a relative factor $m_S/m_{Q_q}$. The scalar-pseudoscalar operator $(S^\dagger S) (\bar{q} \, i \gamma^5 \gamma^\mu q)$ reduces to the $\bar{q}$ suppressed P- and T-odd SD operator $\mathcal{O}_{10}$ and $i \left( S^\dagger \frac{\partial S}{\partial S} S \right) \bar{q} \, \gamma^\mu i \gamma^5 q$ reduces to the $\bar{v}_L$ suppressed P-even, T-odd SD operator $\mathcal{O}_7$.

$\mathcal{O}_1$ is the leading order operator, because it does not suffer from any velocity suppression. It therefore contributes dominantly when compared to $\mathcal{O}_7$ and $\mathcal{O}_{10}$, unless its Wilson coefficient is suppressed by a factor of $10^{-3}$ or less.


should be emphasized that this suppression should occur for scattering on both neutrons and protons.\footnote{4} Since the relevant hadronic matrix elements are different for neutrons and protons, one \textit{cannot} arrange both cancellations with only a single mediator. In other words, the P- and T-odd operators can only be significant in this model if one introduces at least two mediators. Suppressing all contributions from $O_1$ would then require two relations between the twelve couplings $y_1^q, y_2^q$ and six masses $m_{Q_i}$ to hold to better than one part in $10^3$.

On the other hand, if one introduces mediators with charge $1/3$ and with charge $2/3$, one can require the Lagrangian (3.1) to respect strong isospin invariance. In this case the cancellations for protons and neutrons are in fact (almost) the same. In the most symmetric set-up where each SM quark $q$ is assigned a mediator $Q_q$, the couplings $y_1^q \equiv y_1$ and $y_2^q \equiv y_2$ are flavor-universal, and all mediators have the common mass $m_Q$ the contributions from $O_1$ vanish if

$$|y_1|^2 = \frac{1 - \frac{N^N}{f_T} \frac{m_S}{m_Q}}{1 + \frac{N^N}{f_T} \frac{m_S}{m_Q}} |y_2|^2. \quad (3.4)$$

Here $f_T^N \equiv \sum_q \langle \bar{N}|q q|N\rangle$ is the contribution of light and heavy quarks to the nucleon mass (scalar nucleon bilinear) and $N^N \equiv \sum_q \langle \bar{N}|q\gamma^\mu q|N\rangle$ is the number of valence quarks in the nucleon (vector nucleon bilinear). For our assumption of flavor-universal couplings, $f_T^N$ is given by

$$f_T^N = \sum_{q=u,d,s} \frac{m_N}{m_q} f_{Tq}^N + \frac{2}{27} \left(1 - \sum_{q'=u,d,s} f_{Tq'}^N\right) \sum_{q=c,b,t} \frac{m_N}{m_q}. \quad (3.5)$$

where the heavy quark contribution is due to the trace anomaly of the energy momentum tensor [4].\footnote{5} The vector nucleon bilinear coefficient $N^N$ is 3 for both neutrons and protons reflecting the number of valence quarks. Using the values for $f_{Tq}^N$ given in Ref. [45], the ratio of the nucleon bilinears appearing in Eq.(3.4) is

$$\frac{N^N}{f_T^N} = \begin{cases} 0.212^{+0.043}_{-0.038}, & N = n \\ 0.219^{+0.051}_{-0.044}, & N = p \end{cases}. \quad (3.6)$$

\footnote{4}{For a detector containing a single isotope one would only need to suppress the coefficient $c_1$ for the specific linear combination of neutrons and protons determined by the target. However, the Xenon detectors, which currently give the best bounds, contain several isotopes, and unsuppressed $O_1$ constraints from Germanium are still significantly stronger than $O_{10}$ constraints from Xenon.}

\footnote{5}{The heavy quark contributions to $f_T^N$ in Eq.(3.5) assumes $m_Q \gg m_q$ so that the mediators can be integrated out consistently, leaving only SM quarks behind. This may not be a good approximation for the top quark and its mediator. For flavor-universal couplings the contributions from heavy SM quarks is in any case negligible. However, one might also entertain the possibility that the $|y_1|^2$ scale $\propto m_q$, so that the $Q-S$ two-point function corrections to the SM quark masses scale like $m_Q$; in this case the contributions from heavy SM quarks would be comparable to that in models where the WIMP interacts with nucleons via Higgs exchange.}
Note that the stability of the WIMP $S$ requires $m_S \leq m_Q$, hence eq.(3.4) relates the two Yukawa couplings via an $O(1)$ factor. From the simplified model point of view there is a priori no reason why Eq.(3.4) (or a suitable modification thereof) should hold, i.e. in almost all of the parameter space the contribution from the traditional operator $O_1$ will in fact dominate.\footnote{The contribution to the Wilson coefficient of $O_1$ which is suppressed by $m_S/m_Q$ has not been included in \cite{36}. If this contribution is neglected, the Wilson coefficient vanishes for $|y_1| = |y_2|$, which could be motivated via chiral symmetry. However, an otherwise unsuppressed contribution from $(m_S/m_Q)O_1$ would still dominate over contributions from $O_7$ and $O_{10}$ unless $m_Q \gtrsim 10^3 m_S$ in which case the scattering cross section is anyway very small.}

Imposing Eq.(3.4), $O_{10}$ remains as dominant operator since its contribution to the scattering matrix element is only suppressed by $\frac{m_q^2}{m_S^2}$ as compared to a suppression $\frac{m_q^2}{m_Q^2}$, as discussed in Section II.

The quark EDM, $d_q$, can be calculated as the coefficient of a dimension-5 P- and T-odd interaction term $(-i/2) \bar{q} \sigma_{\mu \nu} \gamma_q F^{\mu \nu}$ at vanishing momentum transfer. Since the mediator $Q$ is charged under $SU(3)_{C}$, it can couple to gluons, hence non-vanishing chromo-quark EDMs might also be generated. These are calculated similar to quark EDMs, by finding the coefficient of $(-i/2) \bar{q} \sigma_{\mu \nu} k_{a} \gamma_q C^{\mu \nu}_{a}$ at vanishing momentum transfer. From the left diagram in Fig. 1 we calculate the quark EDM to be

$$d_q|_{\text{Model I}} = \frac{1}{(4\pi)^2} e Q_q m_{Q_q} \Im(g_{1q}^3 q_{1q}^d) F(m_{q}^2, m_{S}^2, m_{Q_q}^2),$$

(3.7)

where $Q_q$ is the electric charge of the quark $q$, and hence of the fermionic mediator (in units of the proton charge), and the loop function $F(m_{q}^2, m_{S}^2, m_{Q_q}^2)$ is given by

$$F(m_{q}^2, m_{S}^2, m_{Q_q}^2) = \int_0^1 dz m_{q}^2 + z(m_{S}^2 - m_{Q_q}^2 - m_{q}^2) + m_{Q_q}^2.$$

(3.8)

The chromo-EDM, $\tilde{d}_q$ can be obtained by replacing the external photon with a gluon as in Fig. 1a. The chromo-quark EDM is thus also given by Eq. (3.7) after replacing $e Q_q$ with the strong coupling $g_s$. In the limit $m_q \to 0$ the loop function (3.8) varies between $1/(2m_{Q_q}^2)$ for $m_S^2 \ll m_Q^2$ and $1/(3m_{Q_q}^2)$ for $m_S^2 = m_Q^2$.

B. Model II

Model II introduces a fermionic WIMP $\chi$ and one or more complex spin-0 mediator(s) $\Phi$, both odd under a $Z_2$ symmetry which stabilizes $\chi$. This model is similar to Model I, but the spins of the WIMP and the mediator have been exchanged, i.e. Model II resembles a supersymmetric model where $\chi$ is the lightest neutralino and the mediators are squarks. The renormalizable $SU(3)_{C} \times U(1)_{em}$ invariant Lagrangian is given by

$$L_{\text{Model II}} = L_{\text{SM}} + i \bar{\chi} \not{D} \chi - m_{\chi} \bar{\chi} \chi$$

$$+ (\partial_{\mu} \Phi_{q}^d)(\partial^{\mu} \Phi_{q}) - m_{\Phi_{q}}^2 \Phi_{q}^d \Phi_{q} - \frac{\lambda_{\Phi_{q}}}{2}(\Phi_{q}^d \Phi_{q})^2$$

$$- (\partial_{\mu} \Phi_{q}^d)(\partial^{\mu} \Phi_{q}) + \frac{\lambda_{\Phi_{q}}}{2}(\Phi_{q}^d \Phi_{q})^2 + h.c.).$$

(3.9)

We have again assumed the two new types of Yukawa couplings, $l_1$ and $l_2$, to be flavor-diagonal in order to avoid new one-loop contributions to FCNC processes due to mediator-WIMP loops. Of course, the electric charge of the mediator $\Phi_{q}$ must again be the same as that of the corresponding quark $q$. Moreover, complex phases in the new Yukawa couplings will again lead to CP violation.

The matrix element for WIMP-quark scattering via $\Phi_q$ exchange in the s-channel is now given by

$$M_{\chi q \to \chi q} = \frac{1}{m_{\Phi_q}^2 - m_{\chi}^2} \left( l_1^d l_1^d \bar{u}(k_2) v(k_1) \right) \left( [\bar{v}(p_1) v(p_2)] - l_1^d l_2^d \bar{u}(k_2) v(k_1) \right) \left( [\bar{v}(p_1) v(p_2)] - l_2^d l_2^d \bar{u}(k_2) v(k_1) \right) \left( [\bar{v}(p_1) v(p_2)] - l_2^d l_2^d \bar{u}(k_2) v(k_1) \right) \left( [\bar{v}(p_1) v(p_2)] - l_2^d l_2^d \bar{u}(k_2) v(k_1) \right).$$

(3.10)

Applying Fierz rearrangement identities from Appendix B to the spinor quadrilinears appearing in eq.(3.10), a number of relativistic effective operators are generated. They are listed in Table III, together with their matching to the NREFT operators in terms of the parameters of this model.
We see that Model II generates the following NREFT operators: \( O_1, O_4, O_6, O_7, O_8, O_9, O_{10}, O_{11} \) and \( O_{12} \). Referring back to Table 1, \( O_6, O_7, O_9, O_{10} \) and \( O_{12} \) are all spin-dependent but are suppressed by at least one factor of \( \bar{q}^+ \sim 10^{-3} \) or \( \bar{q}/m_N \). They can therefore be neglected relative to \( O_4 \), which did not contribute in Model I. Moreover, \( O_8 \) is suppressed by \( \bar{q}^+ \sim 10^{-3} \) and can thus also be neglected. As in Model I the dominant term will come from \( O_1 \) unless its coefficient is tuned to be tiny. The contributions from \( O_4 \) and \( O_{11} \) might be roughly comparable: the latter is suppressed by a single power of \( \bar{q}/m_N \), but enhanced by the coherence factor \( A \) since it does not depend on the spin of the nucleus.

As in Model I, there are two contributions that reduce to \( O_1 \): the scalar-scalar bilinear \( \bar{\chi} \bar{q} q \) contributes proportional to the difference of the absolute squares of the Yukawa couplings \( |q_1|^2 - |q_2|^2 \) whereas the vector-vector bilinear \( \bar{\chi} \gamma^\mu \bar{q} q \gamma^\mu q \) contributes \( \left( |q_2|^2 + |q_1|^2 \right) \). The total coefficient of \( O_1 \) can therefore again be made to vanish by an explicit cancellation. Assuming universal couplings and masses for all mediators, the cancellation condition is

\[
|l_1|^2 = \left( 1 - \frac{N^N}{1 + \frac{N^N}{r}} \right) |l_2|^2.
\]

This relation depends on the same ratio of nuclear matrix elements as the analogous relation (3.4) for Model I. We note that these universality assumptions forbid to also tune the coefficient of \( O_4 \) to zero. This could be arranged for general couplings; however, if the new interactions violate strong isospin, the coefficients of \( O_1^p \) and \( O_1^n \) would have to be tuned to zero separately.

We see from Table III that the coefficient of the leading new operator \( O_{11} \) is proportional to the relative phase between the two new Yukawa couplings; this is true also for the coefficients of the other two P- and T-odd operators, \( O_{10} \) and \( O_{12} \). Since this phase signals CP violation, we expect it to show up in the the quark EDM and CEDMs. The
relevant diagrams in this setup are depicted in Fig. 1b. The resulting quark EDM is given by

$$d^q_{\text{Model II}} = \frac{1}{(4\pi)^2} e Q_q \text{Im}(l^l_l) G(m_q^2, m_{\Phi}^2, m_\chi^2). \tag{3.12}$$

Here $Q_q$ is again the electric charge of $q$, which is equal to the charge of $\Phi_q$, and the loop function $G(m_q^2, m_{\Phi}^2, m_\chi^2)$ is given by

$$G(m_q^2, m_{\Phi}^2, m_\chi^2) = \int_0^1 dz \frac{z(1-z)}{z^2 m_q^2 + z(m_\chi^2 - m_{\Phi}^2 - m_q^2) + m_{\Phi}^2}. \tag{3.13}$$

Again, the quark CEDM can be obtained by replacing $e Q_q$ by the strong coupling $g_s$ in Eq. (3.12). The EDMs in Eq. (3.7) and (3.12) are proportional to the mass of the fermion running in the loop, since the dipole operators violate chirality. Moreover, we see the same CP-violating combination of couplings as in the coefficient of the P- and T-odd NREFT operators. The loop function $G$ varies from $1/(2m_{\Phi}^2)$ for $m_\chi^2 \ll m_{\Phi}^2$ to $1/(6m_{\Phi}^2)$ for $m_\chi^2 = m_{\Phi}$. We note that our loop calculations in (3.7) and (3.12) are in complete agreement with earlier results in the literature, for instance the model independent calculation of EDMs of ref. [47].

**IV. RESULTS AND DISCUSSION**

We had seen in the previous Section that the same combination of Yukawa couplings appears in the expression of the nEDM and in the coefficients of the P- and T-odd operators contributing to WIMP-nucleon interactions. The
upper bound on the nEDM [30] therefore leads to upper bounds on the WIMP-nucleon scattering cross sections due to these operators. In this section we discuss these constraints quantitatively, focusing on $\mathcal{O}_{10}$ for Model I and $\mathcal{O}_1$ for Model II which are the leading new operators (beyond $\mathcal{O}_1$ and $\mathcal{O}_4$) in these models. We compare the resulting constraints with the most recent XENON1T results [1] as well as with the irreducible “neutrino floor” background in Xenon experiments.

We compute the nEDM induced from quark EDM and CEDM dimension-5 contributions which were given in the previous Section. We assume that these are the only contributions to the (color) EDMs of the quarks. In order to calculate the value of the nEDM from quark EDM $\delta q$ and CEDM $\delta q$, we use the following formula:

$$d_n = g^2_T d_u + g^2_T d_d + g^2_T d_s + 1.1 \, e \left(0.5 \, d_u + d_d\right). \quad (4.1)$$

Here the tensor charges $g^2_T = -0.233(28)$, $g^2_T = 0.774(66)$ and $g^2_T = 0.009(8)$ have been calculated using lattice QCD [49, 50] (see also Refs.[51–53]) at a renormalization scale of 2 GeV. We are not aware of a reliable lattice computation of the contribution of the chromo-EDMs to $d_n$; we therefore employ a computation using QCD sum rules evaluated at a renormalization scale of 2 GeV [54], although there is an $\mathcal{O}(50\%)$ uncertainty in these results [55, 56]. We will again assume that the new Yukawa couplings are universal. In this case the uncertainty in the coefficients in Eq.(4.1) might shift the boundary of the excluded region slightly, without affecting our results qualitatively. We also stress again assume that the new Yukawa couplings are universal. In this case the uncertainty in the coefficients in Eq.(4.1)

$$\rho \sim o$$

of the new operators is to be significant. We therefore employ a computation using QCD sum rules evaluated at a renormalization scale of 2 GeV [54], although there is an $\mathcal{O}(50\%)$ uncertainty in these results [55, 56]. We will again assume that the new Yukawa couplings are universal. In this case the uncertainty in the coefficients in Eq.(4.1) might shift the boundary of the excluded region slightly, without affecting our results qualitatively. We also stress again assume that the new Yukawa couplings are universal. In this case the uncertainty in the coefficients in Eq.(4.1) might shift the boundary of the excluded region slightly, without affecting our results qualitatively. We also stress

Using the computed expressions for quark EDM and CEDM in Eq.(3.7) and (3.12), we find limits on $|\text{Im}(y_1 y_2^*)|$ and $|\text{Im}(l_{11}^L)|$ respectively according to the stringent experimental bound [30] given in Eq.(1.1). We then convert these limits on the imaginary part of the product of Yukawa couplings to limits on the cross section, using the relations given in Tables II and III as well as the following expression for the “WIMP-nucleon scattering cross section” [57]:

$$\sigma_{\mathcal{O}_{11}} = \frac{\mu_{XN}^2}{\pi} (\epsilon_{11}^N)^2 \quad \text{and} \quad \sigma_{\mathcal{O}_{16}} = 3\frac{\mu_{XN}^2}{\pi} (\epsilon_{10}^N)^2. \quad (4.2)$$

We note that due to the chosen normalization of the operators, these expressions look very similar to those for the traditional operators $\mathcal{O}_1$ and $\mathcal{O}_4$.

In order to compare the nEDM-derived 90\% c.l. limit with the XENON1T sensitivity, we assume a standard isothermal DM halo with $\rho_X = 0.3$ GeV cm$^{-3}$, $v_0 = 220$ km/s, $v_r = 232$ km/s and $v_{esc} = 544$ km/s. We use the Mathematica code dmformfactor [14] for computing the Xenon nuclear response functions $W_{NN'}$ but use our own routines to calculate the exclusion limits using the procedure outlined in Appendix A. This leads to the current exclusion limits on the WIMP-nucleon cross section for $\mathcal{O}_{10}$ and $\mathcal{O}_{11}$ at 90\% C.L. from the most recent XENON1T results [1]. We also show the irreducible background level from coherent neutrino-nucleus scattering (“neutrino floor”) as estimated in Ref. [48].

Figure 2 depicts the $\sigma_{\mathcal{O}_{10}} - m_S$ exclusion involving the operator $\mathcal{O}_{10}$ for Model I with four values of the common mediator mass $m_Q$. We see that the current XENON1T exclusion contour does not go below $10^{-5}$ pb. Since $\mathcal{O}_{10}$ is a momentum-suppressed SD operator, the limits plotted are roughly $q^2/(m_{Q}^2 A^3) \sim \mathcal{O}(10^{-5} - 10^{-8})$ suppressed in comparison to the SI results quoted by the XENON1T collaboration.

In order to derive the constraint on the scattering cross section that results from the bound on the nEDM, one can replace the combination $\text{Im}(y_1 y_2^*) m_Q$ by a constant (proportional to the nEDM) divided by the loop function $F$ of eq.(3.8). For $m_S^2 \ll m_Q^2$ the result then scales like $m_{Q}^2/(m_S + m_N)^2$, i.e. it approaches a constant for $m_S \ll m_N \approx 1$ GeV but scales like $1/m_S^2$ for $m_S > 1$ GeV. Note that this constraint is independent of $m_Q$; as long as $m_S^2 \ll m_Q^2$, the four frames of Fig. 2 only differ in the end point of the $x$ axis, which is given by $m_Q$. For $m_S \approx m_Q$ the constraint becomes somewhat weaker, partly because the loop function $F$ becomes smaller, as we saw above, but mostly due to the $1/(m_{Q}^2 - m_S^2)$ factor from the squared $Q$ propagator in the scattering cross section. We see that this indirect constraint is more than twelve orders of magnitude below the present XENON1T sensitivity, and at least eight orders of magnitude below the neutrino floor. In other words, in this scenario the nEDM constraint implies that the contribution from $\mathcal{O}_{10}$ to WIMP-nucleon scattering is totally negligible.

Figure 3 shows analogous results for the $\mathcal{O}_{11}$ operator in Model II. $\mathcal{O}_{11}$ is independent of the nuclear spin but momentum suppressed; in fact, for small $q = |\vec{q}|$ the suppression must be of order $q^2/m_A^2$ rather than $q^2/m_N^2$ since a

$^7$ These quantities can be obtained by setting momentum transfer and WIMP velocity to unity and are used to facilitate comparison of the Wilson coefficients for elastic scattering on different target nuclei across all NREFT operators. However, the physical cross sections for sub-leading operators contain appropriate factors of momentum transfer and/or WIMP velocity.
pointlike (unresolved) nucleus does not “know” about its nucleonic constituents. The XENON1T limit on the cross section is therefore weaker by three to four orders of magnitude compared to the case of the traditional SI operator $O_1$.

Regarding the indirect constraint on the scattering cross section from the nEDM bound, the situation is slightly more complicated here than in Model I because according to Table III there are two contributions to the Wilson coefficient $c_{11}$, with different dependence on the WIMP mass. The constraint can be derived by replacing $\text{Im}(l_1^1 l_2^2)$ by a constant (again involving $d_n$) divided by the WIMP mass and the loop function $G$. As long as $m_\chi^2 \ll m^2_\Phi$, the nEDM constraint on $\text{Im}(l_1^1 l_2^2)/m^2_\Phi$ thus scales like $1/m_\chi$. Let us consider the cases $m_\chi \ll m_N$ and $m_\chi \gg m_N$ separately. In the first case $c_{11}$ is dominated by the contribution $\propto m_N/m_\chi$ given in the third line of Table III. The constraint on $|c_{11}|^2$ then scales like $1/(m_\chi^4)$ and $\mu^2_N \simeq m^2_\chi$, hence the constraint on the scattering cross section varies $\propto 1/m_\chi^2$. For $m_\chi \gg m_N$, $c_{11}$ is dominated by the contribution from the last line in Table III, i.e. the nEDM constraint on $|c_{11}|^2$ scales like $1/m^2_\chi$. Since now $\mu^2_N \simeq m_N$ the nEDM constraint on the cross section again scales $\propto 1/m_\chi^2$. Since the two contributions to $c_{11}$ have opposite sign, there is a cancellation for $m_\chi \sim m_N$, leading to a very strong upper bound on the cross section.

For $m_\chi \simeq m_\Phi$ the nEDM constraint becomes a bit weaker again, chiefly due to the $1/(m^2_\Phi - m^2_\chi)^2$ factor from the squared mediator propagator. Altogether we nevertheless see that the nEDM constraint is again several orders of magnitude below the neutrino floor, i.e. the contribution from $O_{11}$ to the WIMP-nucleon cross section can be neglected.

Although we have shown above that for Model II, $O_{10}$ and $O_{12}$ are expected to be sub-leading due to the presence of $O_1$ and $O_{11}$, in Fig. 4 we repeat the exercise of Fig. 3, but now for $O_{10}$ (top frames) and $O_{12}$ (bottom frames). The current XENON1T limit as well as the neutrino floor for $O_{10}$ are quite similar to those for $O_{10}$ in Model I, see Fig. 2.
FIG. 4: DM-nucleon cross section for $O_{10}$ (the top four subfigures) and $O_{12}$ (the bottom four subfigures) in Model II as a function of DM mass for four values of the mediator mass, $m_{\phi} = 0.1, 0.5, 1$ and 10 TeV. The shaded regions in green and pink denote the regions excluded by current limits on the nEDM. The orange curves denoting the neutrino floor for $O_{10}$ and $O_{12}$ have been taken from Ref. [48].
Table III shows that in Model II the coefficient of $O_{10}$ also receives two contributions, one of which scales like $m_N/m_\chi$; hence the dependence of the nEDM constraint on $m_\chi$ is similar to that in Fig. 3. However, since the contribution that is not proportional to $m_N/m_\chi$ has a much smaller coefficient, the cancellation now happens at significantly larger WIMP mass than in Fig. 3. On the other hand, $O_{12}$ only receives one contribution, which is independent of $m_\chi$ as long as $m_\chi^2 \ll m_N^2$. The shape of the nEDM constraint on the scattering cross section therefore resembles that for Model I shown in Fig. 2. Evidently in all cases the nEDM constraint once again implies that the contributions from these P- and T-odd operators to WIMP-nucleon scattering can safely be neglected.

V. CONCLUSIONS

We studied the detection prospects of P- and T-odd operators arising in the NREFT formalism for WIMP-nucleus elastic scattering. These operators occur only if one includes terms that are suppressed by powers of relative velocity or three-momentum transfer. They can therefore be expected to be significant only if the coefficient of the leading spin-independent operator $O_1$ is very small or vanishes. Moreover, T-odd NREFT operators can only arise in the low-energy limit of a relativistic theory if the latter violates CP invariance. This can give rise to very stringent constraints on the theory, in particular from electric dipole moments.

We analyzed these issues in the framework of two simplified models, taken from ref. [36]. They introduce one or more $s-$channel mediator(s) carrying electric and color charge, with spin 1/2 (Model I) or 0 (Model II); the dark matter particle then has spin 0 or 1/2, respectively. These models will cause FCNC at one-loop unless each mediator only couples to one quark. Moreover, we pointed out that generically the coefficient of $O_1$ is not suppressed in these models. Since by now strong constraints on spin-independent WIMP-nucleon scattering exist for several isotopes, the new operators will be significant only if the coefficient of $O_1$ is suppressed for both neutrons and protons. This generically requires two independent cancellations, which should hold to one part in $10^3$ or better. We also note that one has to introduce several mediators for these cancellations to be possible at all, i.e. in minimal versions of these models, with a single mediator, the contribution from $O_1$ will always dominate.

We therefore assumed that the new couplings and mediator masses are universal for all generations. Since in this case the new terms in the Lagrangian respect strong isospin invariance the two cancellation conditions become almost the same. However, we showed that these scenarios give rise to one-loop contributions to the electric dipole moment of the neutron (nEDM) which are proportional to exactly the same combination of new Yukawa couplings which appears in the coefficients of the new P- and T-odd NREFT operators. The resulting indirect constraint lies well below the irreducible neutrino background. This implies that the contributions from the new T-odd operators on the WIMP-nucleon scattering can safely be neglected in the analyses of both current and future experiments. This conclusion can only be avoided if one relaxes the universality assumption on the new couplings and mediator masses. However, one will then have to impose three independent conditions, in order to cancel the coefficients of $O_1$ for protons and neutrons as well as the neutron EDM. Since these cancellations are not enforced by any symmetry, we conclude that it is very unlikely indeed that the new T-odd operators can contribute significantly. We emphasize that, for a given size of the Wilson coefficients, these are the most important of the total eleven new operators. This indicates that at least in models with charged mediator coupling to light quarks, WIMP-nucleon scattering can safely be analyzed using the traditional operators $O_1$ and $O_4$ only.

Charged mediators coupling exclusively to heavy ($c, b, t$) quarks will contribute to WIMP-nucleon scattering only at one-loop level. On the other hand, the neutron EDM constrains the EDMs of heavy quarks only weakly; the nEDM may then receive dominant contributions at the two-loop level, e.g. via Weinberg’s three-gluon operator [58]. In this case the contribution to the nEDM would still only occur at one order higher in perturbation theory than WIMP-nucleon scattering, as is the case in the scenarios we analyzed here, so the conclusions are likely to be similar. Finally, the P- and T-odd operators can also be generated from simplified models not containing charged mediators [36]. We leave the analysis of these scenarios to future work.

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Appendix A: Xenon1T limits

The latest results from the XENON1T collaboration involved a run time of 278.8 days with a target mass of 1.3 tonnes. To compute the limits, we use here the data reported for ‘0.65t’ of fiducial region in Table I of [1], which has a substantially smaller background. Using the computed differential event rate in (2.3), we calculate the total number of expected signal events $N_s$ using the procedure outlined in [59, 60]:

$$N_s = \text{Exposure} \times \int_{S_{1\text{low}}}^{S_{1\text{upp}}} \sum_{n=1}^{\infty} \text{Gauss}(S1|n, \sigma\sqrt{n}) \times \int_0^\infty \frac{dR}{dE_R} \epsilon(E_R) \text{Pois}(n|n_{ph}(E_R)) \, dS1. \quad (A.1)$$

The differential recoil energy spectrum is converted to a differential rate in the number of photoelectrons $n$ by convoluting it with the product of a Poisson distribution with mean $n_{ph}(E_R)$, which is the expected number of photoelectrons given a recoil energy $E_R$, and the detection efficiency $\epsilon(E_R)$. We obtain $n_{ph}(E_R)$ from the left panel of Fig. 3 of [1] where the constant recoil energy contours intersect the middle of the median nuclear recoil band and the $2\sigma$ quantile curve. We digitize the detection efficiency as a function of recoil energy from the black curve of Fig. 1 of [1] and weigh it with a factor of 0.5, since we only consider the fiducial region, and a factor of 0.475, since we only consider events in the reference region in $(cS2_s,cS1)$ as in Fig. 3 of [1].

In order to finally obtain the event rate differential in the S1 signal, we apply to each photoelectron a convolution with a Gaussian distribution of width $\sigma = 0.4$ PE which parameterizes the response of the XENON1T photomultiplier tubes [60]. We then integrate over the S1 signal region as reported in [1] from $S_{1\text{low}} = 3$ PE to $S_{1\text{upp}} = 70$ PE and multiply by the exposure to get the total number of expected events $N_s$. The experiment observed $N_s = 2$ events with a best fit of number of expected background events $N_b = 0.83$ and therefore, we employ a fairly simple Poisson upper limit on the number of signal events at 90 % C.L. after using (A.1) to place exclusions on different NRFF operators. Using this procedure, we obtain satisfactory agreement (off by a factor of roughly 2) with the XENON1T exclusion limit [1] for the SI response, which is based on more sophisticated likelihood fits.

Appendix B: Fierz Identities

We use the following Fierz rearrangement identities as derived in [61, 62] to derive the relativistic effective operators for Model II.

\begin{align}
(q\chi)(\bar{q}g) &= -\frac{1}{4} \left[ \bar{q}q\chi\chi + \bar{q}\gamma^\mu q\gamma_\mu \chi + \bar{q}\gamma^5 q\chi\gamma^5 \chi - \bar{q}\gamma^\mu \gamma^5 q\chi\gamma_\mu \chi + \frac{1}{2} \bar{q} \sigma^{\mu\nu} q \chi \sigma_{\mu\nu} \chi \right], \quad (B.1) \\
(q\gamma^5 \chi)(\bar{q}\gamma^5 g) &= -\frac{1}{4} \left[ \bar{q}q\chi\chi + \bar{q}\gamma^\mu q\gamma_\mu \chi + \bar{q}\gamma^5 q\chi\gamma^5 \chi + \bar{q}\gamma^\mu \gamma^5 q\chi\gamma_\mu \chi + \frac{1}{2} \bar{q} \sigma^{\mu\nu} q \chi \sigma_{\mu\nu} \chi \right], \quad (B.2) \\
(\bar{q}q)(q\gamma^5 \chi) &= -\frac{1}{4} \left[ \bar{q}q\gamma^5 \chi + \bar{q}\gamma^\mu q\gamma_\mu \gamma^5 \chi + i\epsilon_{\mu\nu\alpha\beta} \bar{q} \sigma^{\mu\nu} q \chi \sigma^{\alpha\beta} \chi + \bar{q}\gamma_\mu \gamma^5 q\chi\gamma_\mu \chi + \bar{q}\gamma^5 q\chi\chi \right], \quad (B.3) \\
(\bar{q}\gamma^5 q)(\bar{q}q) &= -\frac{1}{4} \left[ \bar{q}q\gamma^5 \chi - \bar{q}\gamma^\mu q\gamma_\mu \gamma^5 \chi + i\epsilon_{\mu\nu\alpha\beta} \bar{q} \sigma^{\mu\nu} q \chi \sigma^{\alpha\beta} \chi - \bar{q}\gamma^\mu \gamma^5 q\chi\gamma_\mu \chi + \bar{q}\gamma^5 q\chi\chi \right]. \quad (B.4)
\end{align}

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