New Constraints on Dark Matter Effective Theories from Standard Model Loops

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We consider an effective field theory for a gauge singlet Dirac dark matter (DM) particle interacting with the Standard Model (SM) fields via effective operators suppressed by the scale $\Lambda \gtrsim 1 \text{ TeV}$. We perform a systematic analysis of the leading loop contributions to spin-independent (SI) DM–nucleon scattering using renormalization group evolution between $\Lambda$ and the low-energy scale probed by direct detection experiments. We find that electroweak interactions induce operator mixings such that operators that are naively velocity-suppressed and spin-dependent can actually contribute to SI scattering. This allows us to put novel constraints on Wilson coefficients that were so far poorly bounded by direct detection. Constraints from current searches are already significantly stronger than LHC bounds, and will improve in the near future. Interestingly, the loop contribution we find is isospin violating even if the underlying theory is isospin conserving.

Introduction. A Weakly Interacting Massive Particle (WIMP) is an appealing dark matter (DM) candidate [1–4]. The lack of evidence for New Physics at the Fermi scale motivates us to remain unbiased about the nature of DM and pursue model-independent approaches. Assuming that the DM is the only non-SM particle experimentally accessible is not always justified at colliders [5–8], and simplified models have been recently proposed to overcome this limitation [9–11]. Nevertheless, besides very specific cases (e.g. inelastic DM [12]), it is an excellent approximation for direct searches given the small energy exchanged with the target nuclei. Within this approach, DM interactions with SM fields can be parameterized by higher dimensional operators suppressed by the cutoff scale $\Lambda$, with the main strength of providing model-independent relations among distinct null DM searches [13–20]. However, different search strategies probe different energy scales, and such a separation of scales may have striking consequences when a connection between different experiments or ultraviolet (UV) complete models with experiments is attempted. Indeed, in some cases, loop corrections are known to dramatically alter direct detection (DD) rates [21–29].

In this article, we consider the case of a SM gauge singlet Dirac DM ($\chi$), with $m_\chi < \Lambda$, and we calculate the complete set of one-loop effects induced by SM fields for operators up to dimension 6 that contribute to spin-independent (SI) DM–nucleon scattering. The separation between $\Lambda$ and the DD scale is systematically taken into account via a proper renormalization group (RG) analysis. This procedure requires as a first step the computation of both electroweak (EW) and QCD running of both electroweak (EW) and QCD running contributions to spin-independent (SI) DM–nucleon scattering. This allows us to put novel constraints on Wilson coefficients that were so far poorly bounded by direct detection. Constraints from current searches are already significantly stronger than LHC bounds, and will improve in the near future. Interestingly, the loop contribution we find is isospin violating even if the underlying theory is isospin conserving.

SM$_\chi$ EFT. Our conceptual starting point is a renormalizable theory of New Physics, where interactions between the DM particle $\chi$ and SM fields are mediated by heavy messenger particles with masses of order $\Lambda \gtrsim 1 \text{ TeV}$. Upon integrating out the heavy mediators, the UV complete model is matched at the scale $\Lambda$ onto what we call “SM$_\chi$ EFT”, an effective field theory (EFT) whose dynamical degrees of freedom are $\chi$ and the SM fields. The resulting effective Lagrangian has the schematic form

$$\mathcal{L}_{\text{SM}_\chi} = \sum_{d>4} \mathcal{L}^{(d)}_{\text{SM}_\chi}, \quad \mathcal{L}^{(d)}_{\text{SM}_\chi} = \sum_\alpha C^{(d)}_\alpha O^{(d)}_\alpha. \quad (1)$$
Here, $\alpha$ runs over all possible operators of dimension $d$ allowed by the SM gauge symmetries, and $d = 1/\Lambda^{d-4}$. The Wilson coefficients $C^{(d)}_{\alpha}$ are dimensionless, in general scale dependent, and encode unresolved dynamics at higher scales. DM stability forbids operators with just one DM field, and we do not need more than two $\chi$ fields for our study. By applying Fierz identities, each operator can be expressed as the product of a DM bilinear and a SM-singlet operator built only with SM fields. A basis of operators for DD is obtained following the same procedure described in [32, 34] for pure SM fields. In what follows, we focus on operators up to dimension 6 generated at the matching scale $\Lambda$, and consistently and systematically derive their effects for DD.

At the matching scale where $SU(2)_L$ is unbroken, four effective operators contribute to DM–nucleon scattering at $d = 5$, i.e. the magnetic and electric dipole operators

\[ O^M_5 = \frac{1}{\Lambda^2} \chi \sigma^{\mu\nu} \chi B_{\mu\nu}, \quad O^E_5 = \frac{i}{\Lambda} \chi \sigma^{\mu\nu} \gamma^5 \chi B_{\mu\nu} \]  

(2)

and the Higgs operators

\[ O^S_{HH} = \frac{1}{\Lambda^2} \chi H^\dagger H, \quad O^\mu_{HH} = \frac{i}{\Lambda} \chi \gamma^5 \gamma^\mu H. \]  

(3)

$B_{\mu\nu}$ and $H$ are the $U(1)_Y$ field strength tensor and the SM Higgs doublet, respectively. At $d = 6$, tree-level exchange of messengers can generate interactions between DM currents and either quark or Higgs currents.

\[ O^{IJ}_{SS} = \frac{1}{\Lambda^2} \chi \Gamma^\mu_{IJ} q \gamma^\mu q, \quad O^{\mu\nu}_{HH} = \frac{i}{\Lambda^2} \chi \Gamma^\mu_{IJ} [H^\dagger D^\mu H], \]  

(4)

where $q$ runs over the quark flavors, while $I$ and $J$ stand for either $V$ or $A$, with $\Gamma^\mu_{V} = \gamma^\mu$ and $\Gamma^\mu_{A} = \gamma^\mu \gamma^5$. We define $H^\dagger D^\mu H = H^\dagger (D^\mu H) - (D^\mu H) H$. These are all operators at the scale $\Lambda$ up to dimension 6 that can contribute to the SI cross section. We now investigate their effects on the DD rates.

In the effective Lagrangian for elastic WIMP–nucleon scattering, the heavier SM fields (Higgs, $W$ and $Z$ bosons and $t, b, c$ quarks) have to be integrated out and the SM vacuum expectation value gives rise to quark masses. Therefore, among the operators above only $O^{VV}_{u,d,d}$ enter directly the SI cross section while threshold corrections from the dimension-5 $O^{SS}_{HH}$ generate dimension-7 scalar contributions. The DM–nucleon SI cross section accordingly reads (cf. [15, 22, 33, 36])

\[ \sigma_{SI}^N = \frac{m_N^2 m_N^2}{(m_\chi + m_N)^2 2\pi \Lambda^2} \left| \sum_{q = u,d} C^{VV}_{qq} f_q^N \right|^2 \]

\[ + \frac{m_N}{\Lambda} \left| \sum_{q = u,d,s} C^{SS}_{QQ} f_q^N - 12\pi C^{SS}_{qq} f_Q^N \right|^2, \]  

(5)

with $m_N$ denoting the nucleon mass, and scalar (vector) couplings $f_q^N$ ($f^N_\gamma$). For heavy quarks, the parameter $f^N_\gamma$ is induced by the gluon operator as discussed in [32], see also [36]. Here,

\[ O_{qq}^S = \frac{\alpha_s}{\Lambda^2} \chi \chi G_{\mu\nu} G^{\mu\nu}, \quad O_{qq}^{SS} = m_q \frac{\alpha_s}{\Lambda^3} \chi \chi (\bar{q} q), \]  

(6)

with $G_{\mu\nu}$ denoting the gluon field strength tensor. In the next section we will discuss how the Wilson coefficients of the operators in Eq. (2), Eq. (3) and Eq. (4) at the high scale $\Lambda$ are evolved down to the scale of DD and how they are connected to the Wilson coefficients of the low-scale operators in Eq. (4).

**Threshold corrections and mixing.** At dimension 5, $O^S_{AM}$, $O^S_{AT}$, and $O^S_{St,St}$ do not mix into other operators since they are the lowest dimensional ones, and therefore only threshold corrections have to be computed. The $Z$ boson in $B_{\mu\nu}$, once integrated out, generates $O^{V,Y}_{qq}$ at dimension 6. The photon field is also encoded in $B_{\mu\nu}$ but it is a degree of freedom of the low-energy theory, and the resulting long-range interaction between $\chi$ and nucleons severely constrains the Wilson coefficient of the dipole operator. The Higgs operator $O^S_{HH}$ gives rise to $O^{SS}_{QQ}$ after EW symmetry breaking, and upon integrating out the heavy quarks also the dimension-7 interaction with the gluon field strength $O_{gg}^S$ is generated. This leads to the following threshold corrections

\[ C_{qq}^S = \frac{1}{12\pi m_{\rho}^2} C_{HH}^S, \quad C_{qq}^{SS} = -\frac{\Lambda^2}{m_{\rho}^2} C_{HH}^S, \]  

(7)

whose form shows that the $O_{qq}^{SS}$ contribution induced by tree-level Higgs exchange is enhanced since it scales like $1/(\Lambda m_{\rho}^2)$ instead of $1/\Lambda^3$. Typical scattering cross sections involving DM effective couplings to the SM Higgs (like $C_{HH}^S$) are in the ballpark of current experimental limits [43, 44]. They may also contribute to mono-Higgs production at colliders [16, 17], and for light enough DM ($m_\chi < m_{h/2}$) to the invisible Higgs decay width [45, 51].

The evolution matrix for the operators defined in Eq. (6) only contains one non-vanishing off-diagonal entry, namely $O_{qq}^S$ mixes with $O_{qq}^{SS}$. Using Eq. (7), we find

\[ C_{qq}^S (\mu_0) = \left[ \frac{1}{12\pi} U_{m_\chi,m_\rho}^{(5)} + 2 U_{m_\chi,m_\rho}^{(4)} \right] - \frac{\Lambda^2}{m_{\rho}^2} C_{HH}^S, \]  

(8)

with

\[ U_{m_\chi,m_\rho}^{(n)} = \frac{3C_F}{\pi\beta_0} \ln \frac{\alpha_s (\Lambda)}{\alpha_s (\mu)} \]  

Here, $n_f$ is the number of active flavors, $\beta_0 = 11 - \frac{2}{3} n_f$ and $C_F = 4/3$. $\mu_0 < m_{h/2}$ is the low-energy scale relevant for DD. The mixing between $O_{qq}^{SS}$ and $O_{qq}^S$ has already been calculated in [23, 26, 30, 52]. We find that this has a numerically negligible impact on $\sigma_{SI}^N$. The reason is that it yields a contribution to $C_{qq}^S$ proportional to $C_{qq}^S$ but the effect of $C_{qq}^S$ in the cross section is enhanced by a factor of $12\pi$ compared to the scalar contribution.
Let us now turn to the dimension-6 operators (see Eq. [11]). Since we focus on SI interactions, only vector DM bilinears are relevant. Concerning quark currents, no QCD renormalization effect has to be taken into account: singlet quark vector currents are conserved under strong interactions and there is no one-loop RG contribution from the axial anomaly. However, EW corrections give rise to an interesting effect which has not been considered so far, namely the mixing of $O_{q q}^{V A}$ into $O_{H H D}^V$, which affects DD rates. $^6$ There are six diagrams contributing to this mixing, two of which are shown in Fig. 1. The result is proportional to the mass of the quark in the loop, i.e. to the Yukawa couplings $Y_q$, and it is therefore dominated by the top quark and to a less extent by the bottom quark. Solving the RG equation, we obtain

$$C_{H H D}^V (\mu) = C_{H H D}^V (\Lambda) - \frac{\alpha t N_c}{\pi} C_{t t}^{V A} (\Lambda) \ln \frac{\mu}{\Lambda} - (t \to b)$$

with $\alpha_t = Y_t^2/(4\pi)$. The relative sign between the last two terms is due to the fact that left-handed up- and down-type quarks have opposite eigenvalues of the third weak-isospin component. Here we keep only the top and bottom contributions to the loop. In applying this result, the running scale $\mu$ should be identified with the EW symmetry breaking scale, where the top and the Z are integrated out and the corresponding logarithm is frozen. A non-vanishing value of $C_{H H D}^V$ generates a finite threshold correction to $O_{q q}^{V V}$ and $O_{q q}^{V A}$ at the EW symmetry breaking scale by attaching a quark pair and integrating out the Z boson:

$$C_{u u}^{V V} \to C_{u u}^{V V} + \left(\frac{1}{2} - \frac{4}{3} s_w^2\right) C_{H H D}^V ,$$

$$C_{d d}^{V V} \to C_{d d}^{V V} + \left(-\frac{1}{2} + \frac{2}{3} s_w^2\right) C_{H H D}^V$$

where $s_w$ is the sine of the weak mixing angle. Combining Eq. (10) and Eq. (9), we find

$$C_{u u}^{V V} (\mu) = C_{u u}^{V V} (\Lambda) + \left(\frac{1}{2} - \frac{4}{3} s_w^2\right) C_{H H D}^V (\Lambda) + \left(-\frac{1}{2} + \frac{2}{3} s_w^2\right) C_{H H D}^V (\Lambda) + \left(\frac{\alpha t N_c}{\pi} C_{t t}^{V A} (\Lambda) \ln \frac{\mu}{\Lambda} - (t \to b)\right) ,$$

$$C_{d d}^{V V} (\mu) = C_{d d}^{V V} (\Lambda) + \left(-\frac{1}{2} + \frac{2}{3} s_w^2\right) C_{H H D}^V (\Lambda) + \left(\frac{\alpha t N_c}{\pi} C_{t t}^{V A} (\Lambda) \ln \frac{\mu}{\Lambda} - (t \to b)\right) ,$$

which means that a quark vector current is generated at the low scale, even if at the high scale there is only an axial-vector current. As an application of our results, in the next section we will present limits on $C_{q q}^{V A}$ previously bounded only by collider searches (see e.g. 53).

**Numerical Analysis.** We use our results to put constraints on Wilson coefficients that have not yet been bounded from direct searches. We first consider the scenario where $C_{q q}^{V A}$ is the only non-vanishing coefficient at the scale $\Lambda$, and assume flavor-universal DM–quark couplings. The regions in parameter space allowed by various experiments are shown in Fig. 2, where the matching scale $\Lambda$ is plotted as a function of the DM mass for $C_{q q}^{V A} = 1$. If loop effects are neglected, this operator generates a scattering amplitude which is both SD and velocity suppressed. For this reason, the best bound before our analysis came from collider searches (see e.g. 53), corresponding to the dashed orange line in Fig. 2. The RG induced contribution of $C_{q q}^{V V}$ to $C_{q q}^{V A}$ allows us to equally well constrain this operator from SI measurements. In order to use the experimental bounds on the WIMP–nucleon cross section given in 57, 58, we have to take care of the fact that these

![Fig. 1: Diagrams responsible for the mixing of $O_{q q}^{V A}$ into $O_{H H D}^V$. Graphs originated by crossing or reversing the fermion flow are not displayed.](image-url)
limits were obtained under the assumption of negligible isospin violation. However, as we see from Eq. 11, our loop contribution to \( C_{qq}^{VV} \) is isospin violating, i.e. \( \Delta C_{dd}^{VV} \approx -2 \Delta C_{uu}^{VV} \). Therefore, unlike the isospin-symmetric case, our WIMP–nucleus cross section does not scale just like \( A^2 \) (where \( A \) is the mass number of the target nucleus). The regions allowed by DD measurements are delimited by the green (XENON100) and red (LUX) lines. Remarkably, these bounds are one order of magnitude stronger than the ones from LHC searches (dashed orange). We also study the impact of future SI measurements, and show the projections for the allowed regions from SCDMS [59] (purple) and XENON1T [60] (blue). We also superimpose the line obtained by requiring an \( O_{qq}^{V,A} \)-dominated thermal freeze-out and observe that current experiments completely rule out the thermal window (for \( C_{qq}^{V,A} = 1 \)).

In the SMχ EFT, it is possible to assume that \( C_{qq}^{V,A} \neq 0 \) and \( C_{HHD}^{V} = 0 \) only at one fixed scale (in Fig. 2 this scale is \( \Lambda \)). We extend our analysis to the case where also \( O_{HHD}^{V} \) (and \( O_{HDD}^{VV} \)) is switched on, and we use the matching corrections in Eq. 10 to discuss the effect in terms of an effective \( C_{qq}^{VV} \) at the matching scale \( \Lambda \). In Fig. 3 we show the parameter space regions allowed by LUX in the \( (C_{qq}^{VV}, C_{qq}^{V,A}) \) plane for different values of \( \Lambda \). Any UV complete model generating only (axial-)vector operators must respect these bounds.

**Discussion and Outlook.** In this article we highlighted the importance of a systematic analysis of one-loop effects induced by SM fields to connect effective operators at the New Physics scale with DD rates. We computed all relevant one-loop effects for SI interactions up to dimension 6 (at the scale \( \Lambda \)) for a gauge singlet Dirac WIMP.

Previously known QCD corrections are numerically not very relevant in this case, although they can have drastic effects for electroweak charged candidates (e.g. wino, higgsino [24]). Instead, the new EW corrections that we computed allowed us to use DD data to significantly improve bounds on Wilson coefficients. More specifically, we put constraints on the SD and velocity-suppressed operator \( O_{qq}^{V,A} \). Our bounds are much stronger than LHC measurements and will significantly improve when new data will become available. For non-universal DM couplings, the mixing we computed between heavy and light quark currents induced by photon exchange allows us to constrain \( C_{QQ}^{VV} \) for heavy quarks \( Q = s, c, b, t \).

Although an analysis of UV complete models is beyond the scope of this article, we point out that our EW mixing effect can be relevant for \( Z' \)-portal models [31, 32], if the quarks couple to the \( Z' \) only through the axial current (as in some Eq. 4 GUT models [61, 62]). Kinetic and/or mass mixing between the \( Z \) and \( Z' \) will generate a contribution to \( C_{qq}^{VV} \) which is likely to be small compared to \( C_{qq}^{V,A} \), and has to obey the constraints in Fig. 3.

Our analysis systematically accounts for contributions from operators up to dimension 6 at the scale \( \Lambda \). At dimension 7, an important EW mixing effect is already known: the tensor operator \( O_{qq}^{T} = \frac{1}{\Lambda^2} \chi \sigma^{\mu\nu} \chi \overline{q} H \sigma_{\mu\nu} q \) mixes into the dimension-5 dipole operators \( O_{M}^{T}, O_{E}^{T} \) [28] and the predictions for SI DD rates get sizably affected. This motivates a systematic analysis of all one-loop effects at dimension 7 including EW corrections, building upon the work presented in this article.

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DM interactions in which lepton couplings are dominant are not considered here.

The Wilson coefficients of left-handed up and down quarks must be identical at the scale $\Lambda$ to respect $SU(2)_L$ gauge invariance.

The dimension-6 interactions between DM currents and $\partial^\mu B_{\mu\nu}$ can be removed by a field redefinition [63].

In the DM axial-current sector there is an analogous mixing between $O^{AA}_{qq}$ and $O^A_{HHD}$. Since this only affects SD scattering, it is not relevant for our study. However, for Majorana DM coupling mostly to heavy quarks it can be an important effect.

In the standard notation of [14], our operator $O^{V,A}_{q\bar{q}}$ corresponds to $D7$ with $M_* = \Lambda/\sqrt{C^{V,A}_{q\bar{q}}}$. 

We point out the remarkable fact that in our case isospin violation is entirely due to SM loops and is present even if the UV complete theory does not violate isospin.