

Including the Z in an Effective Field Theory for dark matter at the LHC

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ABSTRACT: An Effective Field Theory for dark matter at a TeV-scale hadron collider should include contact interactions of dark matter with the partons, the Higgs and the Z . This note estimates the impact of including dark matter- Z interactions on the complementarity of spin dependent direct detection and LHC monojet searches for dark matter. The effect of the Z is small, because it interacts with quarks via small electroweak couplings, and the contact interaction self-consistency condition $C/\Lambda^2 < 4\pi/\hat{s}$ restricts the coupling to dark matter. In this note, the contact interactions between the Z and dark matter are parametrised by derivative operators; this is convenient at colliders because such interactions do not match onto low energy quark-dark matter contact interactions.

KEYWORDS: Cosmology of Theories beyond the SM, Beyond Standard Model

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Contents

1	Introduction	1
2	EFT, assumptions and operators	2
3	Estimated limits from the LHC	5
4	From the TeV to the MeV	7
5	Discussion	8

1 Introduction

Various experiments attempt to detect the particle making up the “dark matter” [1, 2] of our Universe. For instance, direct detection(DD) experiments [3–8], search for \sim MeV energy deposits due to scattering of dark matter particles from the galactic halo on detector nuclei. And the Large Hadron Collider (LHC) searches [9, 10] for dark matter pairs produced in multi-TeV pp collisions, which would materialise as an excess of events with missing energy and jets. The LHC and DD searches are at very different energy scales, so different Standard Model (SM) particles are present, and also the quantum interferences are different [11]. The expected rates can be compared in specific dark matter models [12], or, in recent years, several studies [9, 13–21] have compared the LHC and DD sensitivities using a contact interaction parametrisation of the dark matter interactions with the standard model particles.

The LHC bounds obtained in this way are restrictive, and probe smaller couplings than direct detection experiments searching for “spin dependent” interactions between partons and dark matter [6, 7]. These contact interaction studies are referred to as “Effective Field Theory” (EFT), and considered to be relatively model independent. However, the particle content is an input in EFT, and the restrictive LHC limits assume that the dark matter particle is the only new particle accessible at the LHC. Relaxing this assumption can significantly modify the experimental sensitivities [22–24]. This has motivated various simplified models for dark matter searches at the LHC [25–27]. Retaining this assumption, as will be done in this note, is only marginally consistent, because the contact interactions to which the LHC is sensitive would have to be mediated by strongly coupled particles. As recalled in the next section, this implies that colliders can exclude contact interactions of order their sensitivity, but not much larger.

Effective Field Theory (EFT) is a recipe to get the correct answer in a simple way [28, 29]. So this note attempts to compare LHC and DD constraints on dark matter, according to the prescriptions of [28, 29]. An EFT for dark matter at the LHC should parametrise all

possible SM-gauge invariant interactions of the dark matter with other on-shell particles. So first, contact interactions between the dark matter and the Higgs or Z should be included at the LHC. These can interfere with the contact interactions studied in previous analyses, but contribute differently at colliders from in direct detection, so the linear combination of operator coefficients constrained at high and low energy will be different. Secondly, an EFT contains in principle a tower of operators [30] organised in increasing powers of the inverse cutoff scale $1/\Lambda$, and higher orders can only be neglected if there is a sufficient hierarchy of scales: $\Lambda_{\text{NP}} \gg v$. This hierarchy is absent in dark matter production at the LHC. Addressing the importance of higher dimensional operators will be left to a subsequent publication.¹

This note focuses on including the Z in the EFT for dark matter at the LHC, and estimates analytically the consequences of including the lowest dimension operators allowing dark matter interactions with the Z .² Section 2 outlines a peculiar choice of operators for the Z vertex; they are proportional to the momentum-transfer-squared. This choice appears convenient, because the effects of the Z are therefore absent in direct detection. Section 3 estimates the impact of cancellations between Z exchange and dark matter contact interactions with quarks at the LHC, and section 4 recalls the direct detection bounds.

2 EFT, assumptions and operators

The low energy consequences of New Physics from above a scale Λ can be parametrised by contact interactions of coefficient C/Λ^n . Unitarity [20, 31] approximately implies that $C < 4\pi$, and the contact interaction approximation implies that the momentum exchange should be less than Λ . This means that an experiment can *exclude*

$$\frac{4\pi}{\hat{s}} > \frac{C}{\Lambda^2} > \text{sensitivity} \ , \quad (2.1)$$

where \hat{s} is the four-momentum-squared of the process. Low energy experiments, where $\hat{s} \rightarrow 0$, therefore can be taken to exclude everything above their sensitivity. However, the upper limit of eq. (2.1) is relevant for collider searches, where \hat{s} is the invariant mass of the invisibles. This upper limit is rarely taken into account in the literature.

The first step in the EFT recipe to parametrise New Physics from beyond the scale Λ it to add to the Lagrangian (at the scale Λ), all the non-renormalisable operators which can be constructed out of the fields present, consistently with the symmetries of the theory [28, 29]. The coefficients $C_O^{(n)}$ of these operators are unknown ‘‘coupling constants’’ which evolve with scale via Renormalisation Group Equations. This infinite set of operators would be unmanageable, so EFT is useful when there is a hierarchy between the experimental and NP scales. Then only the lowest dimension operators need be considered.

In this note, the dark matter is assumed to be the only new New Physics particle lighter than a TeV, and is taken to be a SM gauge singlet dirac fermion χ with a conserved

¹Higher dimensional operators can contain more fields and be suppressed by phase space, or contain Higgs fields and be suppressed by $\langle H \rangle^2/\Lambda^2$, or contain derivatives and be dangerous.

²Contact interactions between dark matter and the Z have been proposed in [27] as a benchmark model, assuming other contact interactions to be absent.

parity, and of mass $m_\chi \geq m_Z/2$ (maybe $\geq m_h/2$), to avoid bounds on the coupling to the Z from the invisible width³ of the Z (and Higgs). So the particle content of the EFT for χ at the LHC should be χ , plus all relevant particles of the SM, which I take to be the partons, the Higgs, and the Z .

The operators should be SM gauge invariant, to profit from our knowledge of the SM gauge sector. They are of dimension > 4 , and should attach a $\chi\bar{\chi}$ pair to partons, to the Higgs, or to the Z . The quark operators are taken generation diagonal; flavour-changing operators were considered in [33]. The quarks are chiral because the operators are SM gauge invariant, and also because opposite chiralities do not interfere at the LHC. The dark matter currents are taken in a vector, axial vector, etc basis because these do not interfere in direct detection, nor at the LHC in the limit where the χ mass is neglected, as done here.

I focus on operators of lowest dimension, that is six and seven. This is an arbitrary simplification, because $\Lambda \sim \text{TeV}$, which is the energy scale probed at the LHC. The contact interactions considered here therefore do not provide a “model-independent” parametrisation of the interactions of χ with the SM. This problem is left for a later publication. Concretely, Λ will be taken as 1- 2 TeV, for reasons discussed above eq. (3.2). Experimental limits on contact interactions will therefore be presented as limits on the dimensionless coefficient $C_O^{(n)}$.

At dimension six, there are vector and axial vector χ currents coupled to quarks:

$$\frac{C_{QX,V}}{\Lambda^2} \bar{\chi}\gamma_\mu\chi\bar{Q}_i\gamma^\mu P_X Q_i \quad , \quad -\frac{C_{QX,A}}{\Lambda^2} \bar{\chi}\gamma_\mu\gamma_5\chi\bar{Q}_i\gamma^\mu P_X Q_i \quad (2.2)$$

where the quarks Q_i are first generation SM multiplets $\{q_L, u_R, d_R\}$, and P_X is the appropriate chiral projector.

The contact interactions between the dark matter and the Z boson are taken as

$$\begin{aligned} -\frac{C_{Z,V}}{\Lambda^2} D^\mu B_{\mu\nu} \bar{\chi}\gamma_\mu\chi &\rightarrow s_w p_Z^2 \frac{C_{Z,V}}{\Lambda^2} Z^\mu \bar{\chi}\gamma_\mu\chi \\ \frac{C_{Z,A}}{\Lambda^2} D^\mu B_{\mu\nu} \bar{\chi}\gamma_\mu\gamma_5\chi &\rightarrow -s_w p_Z^2 \frac{C_{Z,A}}{\Lambda^2} Z^\mu \bar{\chi}\gamma_\mu\gamma_5\chi \end{aligned} \quad (2.3)$$

where to the right of the arrow is the resulting vertex, B^μ is the hypercharge gauge boson with coupling $g' = e \tan \theta_W \equiv e s_w / c_w$, $B^{\mu\nu} = \partial^\mu B^\nu - \partial^\nu B^\mu$, and a term $\propto p_Z \cdot Z$ was dropped after the arrow in the axial current operator, assuming the Z was produced by light quarks. There is in addition a “dipole moment” operator $B^{\mu\nu} \bar{\chi}\sigma_{\mu\nu}\chi$, which is neglected here because it also induces dark matter interactions with the photon [34, 35] which are more interesting.

Then at dimension seven, there are four-fermion operators:

$$\begin{aligned} \frac{C_{d,S}^{(7)}}{\Lambda^3} \bar{\chi}\chi \frac{1}{2} \left(\bar{q}_L H d + [\bar{q}_L H d]^\dagger \right) \quad , \quad \frac{C_{d,P}^{(7)}}{\Lambda^3} \bar{\chi}\gamma_5\chi \frac{1}{2} \left(\bar{q}_L H d + [\bar{q}_L H d]^\dagger \right) \\ \frac{C_{d,T}^{(7)}}{\Lambda^3} \bar{\chi}\sigma^{\mu\nu}\chi \frac{1}{2} \left(\bar{q}_L H \sigma_{\mu\nu} d + [\bar{q}_L H \sigma_{\mu\nu} d]^\dagger \right) \end{aligned}$$

³For $m_\chi < m_Z/2$, the invisible width of the Z (at “2 σ ”, so [32] $\Gamma(Z \rightarrow \chi\bar{\chi}) \leq 3 \text{ MeV}$) imposes that $|C_{Z,B}| < 8.9(\Lambda/\text{TeV})^2$, for $B = V, A$.

(and similarly for u quarks, but with a charge conjugate Higgs field), interactions with the gluons:

$$\frac{C_{gg,S}^{(7)}}{\Lambda^3} \bar{\chi} \chi G_{\mu\nu}^A G^{\mu\nu,A} \quad , \quad \frac{C_{g\tilde{g},P}^{(7)}}{\Lambda^3} \bar{\chi} \gamma_5 \chi G_{\mu\nu}^A \tilde{G}^{\mu\nu,A} \quad ,$$

and double-derivative interactions between dark matter and the Higgs:

$$\begin{aligned} H^\dagger D^\mu D_\mu H \bar{\chi} \chi &\rightarrow -m_W^2 W_\mu^+ W^{-\mu} \bar{\chi} \chi - m_Z^2 Z_\mu Z^\mu \bar{\chi} \chi + \frac{v p_h^2}{\sqrt{2}} h \bar{\chi} \chi \\ H^\dagger D^\mu D_\mu H \bar{\chi} \gamma_5 \chi &\rightarrow -m_W^2 W_\mu^+ W^{-\mu} \bar{\chi} \gamma_5 \chi - m_Z^2 Z_\mu Z^\mu \bar{\chi} \gamma_5 \chi + \frac{v p_h^2}{\sqrt{2}} h \bar{\chi} \gamma_5 \chi \end{aligned} \quad (2.4)$$

where $\langle H \rangle = v = 174$ GeV, p_h is the four-momentum of the physical Higgs particle h , and after the arrow are the interactions induced by the operator.

The Z and Higgs operators are chosen $\propto p^2$ so that they are relevant at the LHC where the Z and Higgs are external legs in the EFT, but do not contribute in the low-energy scattering of DD. This choice should be acceptable, because the operator basis can always be reduced by using the equations of motion [36]. Focussing for simplicity on the hypercharge boson B , and neglecting gauge-fixing terms, the equations of motion are [37]

$$D_\mu B^{\mu\nu} = g' y_H (H^\dagger D^\nu H - (D^\nu H)^\dagger H) + g' \sum_\psi y_\psi \bar{\psi} \gamma^\nu \psi \quad (2.5)$$

where ψ is a SM fermion of hypercharge y_ψ . Usually [37], operators containing the double derivative on the left of eq. (2.5) are dropped, and the operators containing the Higgs v.e.v. squared $\langle H^\dagger \overleftrightarrow{D}^\nu H \rangle$ are retained. In this usual basis, $\chi - Z$ interactions could be parametrised by $(\bar{\chi} \gamma^\mu \chi) H^\dagger \overleftrightarrow{D}^\nu H$, in which case the matrix element for Z exchange at the LHC is $\propto m_Z^2 / (p_Z^2 - m_Z^2)$, so negligible for $p_Z^2 \gg m_Z^2$. But Z exchange should be included in the quark- χ contact interaction used in direct detection, so the coefficient of the operators of eq. (2.2) would not be the same in direct detection as at the LHC. To avoid this discrepancy, I retain the derivative operators of eq. (2.3), and use eq. (2.5) to remove the operator $\propto \langle H^\dagger \overleftrightarrow{D}^\nu H \rangle$. This means that the Z couples significantly to χ at the LHC, but negligibly in DD, and the operator coefficients do not change when the Z is matched out of the theory.

In the case of the Higgs, the equation of motion is

$$D_\mu D^\mu H = \mu^2 H - \lambda H^\dagger H H - \bar{e} Y_e^\dagger P_L \ell - \bar{d} Y_d^\dagger q_L + \varepsilon \bar{q}_L Y_u u$$

where Y_f are Yukawa matrices. This has been used to exchange the more usual $(H^\dagger H)^2 \bar{\chi} \chi$, and $(H^\dagger H) \bar{\chi} \chi$ operators for the double-derivative interactions between dark matter and the Higgs given in eq. (2.4). Notice that it is possible to use the equations of motion to replace two operators $(\bar{\chi} \chi H^\dagger H$ and $\bar{\chi} \chi (H^\dagger H)^2)$ with one (involving the DM and $H^\dagger D^2 H$), because I am only interested in the h - χ - $\bar{\chi}$ interaction induced by these operators. The linear combination of operators $[\mu^2 H^\dagger H - \lambda (H^\dagger H)^2] \bar{\chi} \chi$, which is orthogonal to the combination in the Equations of Motion, gives a vanishing h - χ - $\bar{\chi}$ interaction, due to the minimisation condition of the Higgs potential. As in the case of the Z , the derivative operators of

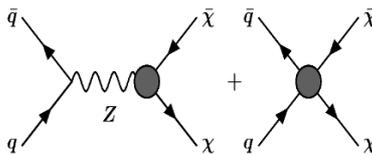


Figure 1. Effective interactions contributing to $q\bar{q} \rightarrow \chi\bar{\chi}$ at the LHC. The coefficient of the four fermion operator is $C_{q,AX}/\Lambda^2$, and the effective axial vector coupling of the Z to dark matter is $s_w p_Z^2 C_{Z,A}/\Lambda^2$.

eq. (2.4) are interesting, because they give a higgs coupling to dark matter $\propto p_h^2$, which has the desirable feature of being relevant at the LHC where the Higgs is in the effective theory, but not contributing at low energy.

This note focusses on the interactions of χ with the Z (eq. 2.3), and with the quark currents of eq. (2.2) which can interfere with Z exchange. So the dimension seven operators will be neglected in the following sections. However, it is interesting to first review the sensitivity to the coefficients of the operators of eq. (2.4). The dark matter interactions to W and Z pairs were studied in [38], who used $U(1)_{\text{em}} \times SU(3)$ invariant operators such that these contact interactions have dimension five with coupling $1/\Lambda_{\text{CHLR}}$. They find that the 8 TeV LHC with luminosity 25 fb^{-1} could probe $\Lambda_{\text{CHLR}} \lesssim \text{TeV}$. This constrains the coefficients of the operators of eq. (2.4) to be $\lesssim 1/(\text{TeV} m_W^2)$, which is not restrictive. For $m_\chi < m_h/2$, a more significant limit of 10 TeV^{-3} arises from requiring $\Gamma(h \rightarrow \chi\bar{\chi}) \lesssim \Gamma(h \rightarrow b\bar{b})$. This restriction should be reasonable [39, 40] because the Higgs is observed to decay to $b\bar{b}$.

3 Estimated limits from the LHC

Dark matter particles are invisible to the LHC detectors, so pair production of χ s can be searched for in events with missing transverse energy (\cancel{E}_T), which can be identified by jet(s) radiated from the incident partons. The principle Standard Model background for such “monojet” searches is Z + jet production, followed by $Z \rightarrow \bar{\nu}\nu$. The 8 TeV LHC is sensitive to dark matter contact interactions with $C/\Lambda^2 \sim \text{TeV}^{-2}$.

Given the operators of eqs. (2.2) and (2.3) at the LHC, the axial vector dark matter current can interact with quarks Q via the diagrams of figure 1, which can be written as a four-fermion interaction of coefficient

$$\begin{aligned}
 c_{QX,A} &= C_{QX,A} + g_X^Q \frac{g_{s_w}}{2c_w} \frac{p_Z^2}{p_Z^2 - m_Z^2} C_{Z,A} \\
 &\xrightarrow{p_Z^2 \gg m_Z^2} C_{QX,A} + g_X^Q \frac{g_{s_w}}{2c_w} C_{Z,A} \quad , \tag{3.1}
 \end{aligned}$$

where $g_X^Q = \{1 - \frac{4}{3}s_w^2, -\frac{4}{3}s_w^2, -1 + \frac{2}{3}s_w^2, \frac{2}{3}s_w^2\}$ for $\{u_L, u_R, d_L, d_R\}$ [32]. A similar expression can be obtained for the vector χ current. The Z exchange looks like a contact interactions for large $p_Z^2 = M_{inv}^2 \gg m_Z^2$, where M_{inv}^2 is the invariant mass-squared of the dark matter pair. This is a useful approximation, because the arguments below suggests that most $\chi\bar{\chi}$ events arise at larger M_{inv}^2 .

The aim here is to analytically estimate the invisible four-momentum-squared M_{inv}^2 , by comparing the partonic cross-sections for $\nu\bar{\nu}$ and $\chi\bar{\chi}$ production. I assume that the QCD part of the amplitude is identical in both cases, so it does not need to be calculated. This allows for an arbitrary number of jets, which is more difficult to simulate [41] (the data frequently contains more than one jet [9]). In the matrix element for jets $+\nu\bar{\nu}$ will appear

$$g_X^Q \frac{g^2}{4c_W^2} \frac{1}{p^2 - m_Z^2 + im_Z \Gamma_Z} (\bar{Q}\gamma^\alpha P_X Q)(\bar{\nu}\gamma_\alpha P_L \nu)$$

whereas, for DM production via the $\bar{\chi}\gamma^\mu\gamma_5\chi$ current, this is replaced by:

$$\frac{c_{QX,A}}{\Lambda^2} (\bar{Q}\gamma^\alpha P_X Q)(\bar{\chi}\gamma_\alpha\gamma_5\chi) .$$

Then the full matrix element must be squared, and integrated over the phase space of N jets and two invisible particles. The invisibles can be treated as a single particle of variable mass $p^2 = M_{inv}^2$, using the identity

$$d\Phi_{N+2} = \delta^4\left(P_{in} - \sum q_i - p\right) \prod_{i:1\dots N} \left(\frac{d^3q_i}{2E_i(2\pi)^3}\right) \\ \times (2\pi)^3 dp^2 \delta^4(p - p_\chi - p_{\bar{\chi}}) \frac{d^3p_\chi}{2E_\chi(2\pi)^3} \frac{d^3p_{\bar{\chi}}}{2E_{\bar{\chi}}(2\pi)^3} .$$

Neglecting spin correlations and the dark matter mass, the invisible phase space integral over the gamma-matrix trace for the invisible fermions gives $M_{inv}^2/(8\pi)$ for χ s, and $3M_{inv}^2/(16\pi)$ for neutrinos. For neutrinos in the final state, $M_{inv}^2 = m_Z^2$ due to the delta-function-like behaviour of the Z propagator-squared. However, for dark matter, the dM_{inv}^2 phase space integral will privilege larger values of M_{inv}^2 . Treating the N jets of the event as a particle of negligible mass, the upper bound on M_{inv}^2 is $\gtrsim 4\cancel{E}_T^2$, where \cancel{E}_T is the invisible transverse energy. The CMS study [9] uses the range $400 \text{ GeV} \leq \cancel{E}_T \lesssim \text{TeV}$. However, the assumption that the jet emission part of the cross-section is the same as for ν pairs will fail, if M_{inv}^2 is a significant fraction of the energy of the event. With the M_{inv} cutoff ranging from 800 GeV to 2 TeV, requiring that the dark matter contribute $\lesssim 1/6$ [9] of the SM background, gives an estimated bound $\Lambda \gtrsim 880 \rightarrow 2200 \text{ GeV}$, for $c_{uL,A} = c_{uR,A} = c_{dL,A} = c_{dR,A} = 1$. This compares favourably to the CMS bound of $\Lambda > 950 \text{ GeV}$, for $C_{uL,A} = C_{uR,A} = C_{dL,A} = C_{dR,A} = 1$. Since the analytical estimate is reasonable, most of the dark matter signal probably comes from $M_{inv}^2 \gg m_Z^2$, and the approximation (3.1) is consistent. However, the analytic bound is a bit restrictive (perhaps in part because it includes any number of jets), so in the remainder of the paper, the CMS limit of 950 GeV will be used.

There is also an upper limit on the C s which a collider can exclude, eq. (2.1), from requiring that the contact interaction approximation be self-consistent: $C/\Lambda^2 < 4\pi/M_{inv}^2$. Since the previous analytic estimate reproduces the CMS bound for $M_{inv}^2 \sim \text{TeV}^2$, the consistency condition is taken as $C < 4\pi$. For the axial χ current with $\Lambda = \text{TeV}$, the CMS

limit and eq. (2.1) give 3 independent bounds on $\{c_{qL,A}, c_{uR,A}, c_{dR,A}\}$:

$$\begin{aligned}
 4\pi &\lesssim \sqrt{\frac{2}{3}\left|C_{qL,A} + \frac{2}{15}C_{Z,A}\right|^2 + \frac{1}{3}\left|C_{qL,A} - \frac{1}{6}C_{Z,A}\right|^2} \lesssim \sqrt{2} \\
 4\pi &\lesssim \left|C_{uR,A} - \frac{1}{15}C_{Z,A}\right| \lesssim \sqrt{3} \\
 4\pi &\lesssim \left|C_{dR,A} + \frac{1}{30}C_{Z,A}\right| \lesssim \sqrt{6}
 \end{aligned} \tag{3.2}$$

where the first line is the summed contributions of u_L and d_L , the fractions are approximations $gg_{\chi}^Q s_w / 2c_w$, and the d to u pdf ratio is taken 1/2. Similar limits apply for the vector operator of eq. (2.2).

It can be seen already from eq. (3.2), that including the interactions with the Z will make little differences to the LHC limits on the $C_{QX,A}$: for the doublet quarks, the Z contribution cannot cancel simultaneously against the u_L and d_L contributions, and the Z contribution is irrelevant for the singlet quarks, because also $C_{Z,A}$ must be $\lesssim 4\pi$. The parameters ruled out by the first and second eqs. of (3.2) are represented as the central regions in figure 2.

4 From the TeV to the MeV

In direct detection, the dark matter scatters non-relativistically off nuclei. Therefore, to translate the EFT from the TeV to the MeV, the Z must be removed, the effects of QCD loops in running the operator coefficients should be included, and the quarks must be embedded in the nucleons.

To remove the Z , the Greens function for two quarks and two χ s in the effective theory with a Z , should be matching to the same Greens function in the theory without a Z . Since the matching is performed at zero momentum for the fermion legs, the contact interactions of eq. (2.3) do not contribute, and the coefficients of the four-fermion operators of eq. (2.2) remain the same after the Z is “matched out”. The Z vertices were taken $\propto p_Z^2$ to obtain this.

The light quark currents $\bar{q}\gamma^\mu P_X q$ are conserved in QCD, so do not run. Also, since χ is a SM gauge singlet and the only dark sector particle below the TeV, I suppose that the operators with vector and axial vector χ currents do not mix below the TeV. See e.g. [42] about loop effects mixing various operators involving dark matter and the SM.

Finally, the quark currents can be embedded in nucleons $N = \{p, n\}$ using identities [43] such as

$$\langle N | \bar{Q}_i \gamma^\mu Q_i | N \rangle = c_{V,i}^N \langle N | \bar{\psi}_N \gamma^\mu \psi_N | N \rangle$$

where $c_{V,u}^p = c_{V,d}^n = 2$, and $c_{V,d}^p = c_{V,u}^n = 1$, because this current counts valence quarks in the nucleon. The axial quark current is proportional to the nucleon spin:

$$\langle N | \bar{Q}_i \gamma^\mu \gamma_5 Q_i | N \rangle = 2s^\mu \Delta Q_i^N = \Delta Q_i^N \langle N | \bar{\psi}_N \gamma^\mu \gamma_5 \psi_N | N \rangle$$

where the proportionality constants are measured [44] as $\Delta u^p = \Delta d^n = 0.84$, $\Delta d^p = \Delta u^n = -0.43$. In the zero-momentum-transfer limit of non-relativistic scattering, the dark matter can have spin-dependent interactions via the axial current, or spin-independent interactions via the first component of the vector current.

The spin-independent scattering amplitude for χ on a nucleon, is a coherent sum of vector and scalar interactions, for quarks of both chiralities and all flavours. The experimental limit on the cross-section per nucleon is $\sigma_{SI} \lesssim 10^{-44} \text{ cm}^2$ for $m_\chi \sim 100 \text{ GeV}$ [4, 5]. For the proton ($C_{uR} \leftrightarrow C_{dR}$ for the neutron), with $C_{qR,V} = \frac{1}{3}(C_{dR,V} + 2C_{uR,V})$, this gives [43]

$$\sigma_{SI} \simeq \frac{1}{\pi} \left[\frac{3m_N}{2\Lambda^2} (C_{qL,V} + C_{qR,V} + \dots) \right]^2 \lesssim 3 \times 10^{-17} \text{ GeV}^{-2}$$

where the $+\dots$ contains scalar contact interactions neglected in this note. For $\Lambda = \text{TeV}$, this gives

$$[C_{qL,V} + \frac{1}{3}(C_{dR,V} + 2C_{uR,V}) + \dots] \lesssim 10^{-2} \quad (SI). \quad (4.1)$$

The spin dependent cross-section per proton is [43]

$$\sigma_{SD} \simeq m_p^2 \left[\frac{.42(C_{qL,A} + C_{uR,A} - 2C_{dR,A})}{2\Lambda^2} \right]^2 \lesssim \frac{10^{-10}}{4} \text{ GeV}^{-2}$$

where the experimental bound is for $m_\chi \sim 100 \text{ GeV}$. For $\Lambda = \text{TeV}$, this gives

$$|(C_{qL,A} + C_{uR,A} - 2C_{dR,A})| \lesssim 20 \quad (SD). \quad (4.2)$$

Comparing to eq. (3.2) shows that the contact interactions explored by SD direct detection experiments are mediated by physics which is not a contact interaction at the LHC, so are not excluded by the limits given in eq. (3.2). The limit (4.2) is represented in figure 2 as the vertical exclusions.

5 Discussion

From a bottom-up EFT point of view, it is important to include all operators which can interfere, when computing experimental constraints. This is to allow for cancellations. Including several operators which do not interfere improves the bound, but is not otherwise motivated. In this note, operators with vector and axial vector currents for the dark matter fermion χ were presented as an example, which illustrates two points.

First, the EFT at the LHC contains more particles than the light partons and dark matter that are relevant in direct detection. At the LHC, the Higgs and Z should also be included. Matching the high and low energy EFTs, as done in this note, suggests that the LHC constrains several combinations of operator coefficients that are different from direct detection, as can be seen by comparing eqs. (3.2) and (4.2). However, the contribution of the Z is relatively unimportant, because its couplings to singlet quarks are small, and it interferes with opposite sign with u_L and d_L . The LHC limits on the dark matter couplings to quarks and the Z are represented as the central exclusion areas of figure 2: the coupling

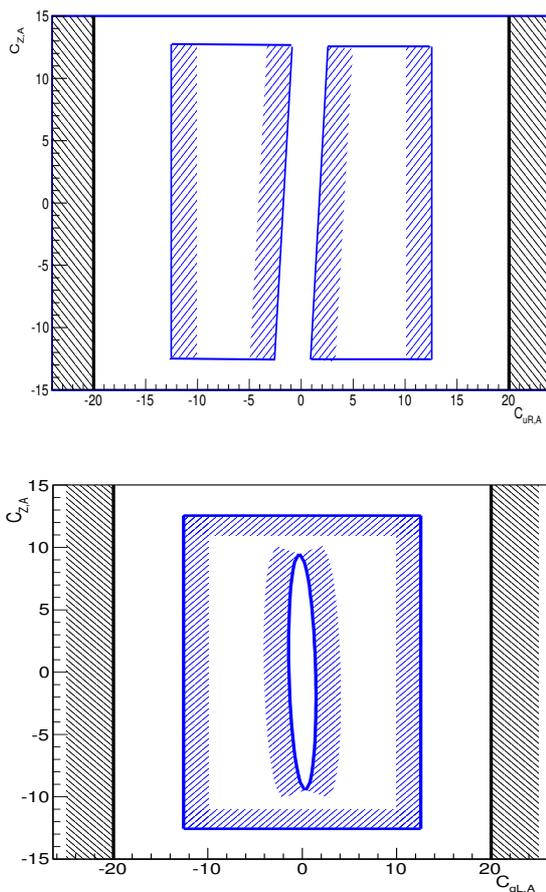


Figure 2. Parameter space excluded by spin dependent direct detection experiments (at either side), and the LHC (central region), for a dark matter fermion of mass ~ 100 GeV, with contact interactions with the Z parametrised by $C_{Z,A}$ (see eq. (2.3)), and with u_R quarks in the left plot, and the doublet q_L in the right plot (see eq. (2.2)). $\Lambda = \text{TeV}$, and all other coefficients are zero. The upper limit of the LHC exclusions is estimated from eq. (2.1).

to quarks is more constrained than the coupling to the Z , and arbitrary axial current dark matter interactions to quarks cannot be allowed by tuning the dark matter coupling to the Z . This is because there is a self-consistency upper bound on contact interaction coefficients at colliders $C/\Lambda^2 < 4\pi/\hat{s}$ (see eq. (2.1)). It is important to notice that this upper bound also implies that the LHC limits do not exclude the parameter space probed by spin dependent direct detect experiments.

Second, an interesting difference between direct detection and collider experiments, is that quarks of different chirality and flavour interfere in direct detection, whereas the LHC can constrain the interactions of dark matter with each flavour and chirality of quark individually. This is related to the relative unimportance of the Z : it cannot cancel separately against the contributions of u_L, d_L, u_R and d_R .

In summary, the rules of bottom-up Effective Field Theory say that one should include all operators up to some specified dimension. So to parametrise at dimension six the axial

vector interactions of dark matter with quarks, one should include contact interactions of dark matter with the quarks and with the Z . Including interactions with the Z that are $\propto p_Z^2$, as done here, suggests that these are not crucial.

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