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Chiral power counting of one- and two-body currents in direct detection of dark matter

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ABSTRACT

We present a common chiral power-counting scheme for vector, axial-vector, scalar, and pseudoscalar WIMP–nucleon interactions, and derive all one- and two-body currents up to third order in the chiral expansion. Matching our amplitudes to non-relativistic effective field theory, we find that chiral symmetry predicts a hierarchy amongst the non-relativistic operators. Moreover, we identify interaction channels where two-body currents that previously have not been accounted for become relevant.

1. Introduction

Elucidating the nature of dark matter is one of the most pressing challenges in contemporary particle physics and astrophysics. Still, one of the dominant paradigms rests on a weakly-interacting massive particle (WIMP), such as the neutralino in supersymmetric extensions of the standard model (SM). A WIMP can be searched for at colliders, in annihilation signals, or in direct-detection experiments, where the recoil energy deposited when the WIMP scatters off nuclei is measured. Recent years have witnessed an impressive increase in sensitivity, e.g., from XENON100 [1], LUX [2], and SuperCMDS [3], which will further improve dramatically with the advent of ton-scale detectors, XENON1T [4] and LZ [5]. In the absence of a signal, direct-detection experiments provide more and more stringent constraints on the parameter space of WIMP candidates. To derive these constraints and to interpret a future signal, it is mandatory that the nucleon matrix elements and the nuclear structure factors, which are required when transitioning from the SM to the nucleon to the nucleus level, be calculated systematically and incorporate what we know about QCD.

Effects at the level of the nucleus can be described by an effective field theory (EFT) whose degrees of freedom are non-relativistic (NR) nucleon and WIMP fields [6,7]. This NREFT has been recently used in an analysis of direct-detection experiments [8]. In this approach, scales related to the spontaneous breaking of chiral symmetry of QCD are integrated out, with the corresponding effects subsumed into the coefficients of the EFT. In the context of nuclear forces, such an EFT is called pionless EFT. To derive limits on the WIMP parameter space, information from QCD has then to be included in the analysis in a second step.

Alternatively, one can start directly from chiral EFT (ChEFT) to incorporate the QCD constraints from chiral symmetry [9–16], which makes predictions for the hierarchy among one- and two-body currents. Based on ChEFT, scalar and axial-vector two-body currents were recently considered in [10] and [11,12], respectively. Moreover, lattice QCD can be used to constrain the couplings of two-body currents [17].

The goal of this Letter is to combine vector, axial-vector, scalar, and pseudoscalar interactions in a common chiral power counting, collect all relevant one- and two-body matrix elements, and match the result onto NREFT. This combines our knowledge of QCD at low energies: the one-body matrix elements correspond to the standard decomposition into form factors, while the two-body scalar [9,10], vector [18–20], and axial-vector [15,21] currents have been calculated as well, the vector current even at one-loop order. Here, we combine these results for their application in direct detection, extending the axial-vector two-body currents to finite momentum transfer and generalizing to the three-flavor case where appropriate. By matching to the NREFT, we find that the chiral symmetry of QCD predicts a hierarchy among the different operators and that two-body currents can be as important as one-body currents in some channels.

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2. Effective Lagrangian and kinematics

We start from the following dimension-6 and -7 effective Lagrangian for the interaction of the WIMP χ, assumed to be a SM singlet, with the SM fields [22]

$$\begin{aligned}\mathcal{L}_\chi = \frac{1}{\Lambda^3}\sum_q\Big[&C_q^{SS}\bar\chi\chi\, m_q\bar q q + C_q^{PS}\bar\chi i\gamma_5\chi\, m_q\bar q q\\ &+C_q^{SP}\bar\chi\chi\, m_q\bar q i\gamma_5 q + C_q^{PP}\bar\chi i\gamma_5\chi\, m_q\bar q i\gamma_5 q\Big]\\ +\frac{1}{\Lambda^2}\sum_q\Big[&C_q^{VV}\bar\chi\gamma^\mu\chi\,\bar q\gamma_\mu q + C_q^{AV}\bar\chi\gamma^\mu\gamma_5\chi\,\bar q\gamma_\mu q\\ &+C_q^{VA}\bar\chi\gamma^\mu\chi\,\bar q\gamma_\mu\gamma_5 q + C_q^{AA}\bar\chi\gamma^\mu\gamma_5\chi\,\bar q\gamma_\mu\gamma_5 q\Big]\\ +\frac{1}{\Lambda^2}\sum_q\Big[&C_q^{TT}\bar\chi\sigma^{\mu\nu}\chi\,\bar q\sigma_{\mu\nu} q + \tilde C_q^{TT}\bar\chi\sigma^{\mu\nu}i\gamma_5\chi\,\bar q\sigma_{\mu\nu} q\Big]\\ +\frac{1}{\Lambda^3}\Big[&C_g^{S}\bar\chi\chi\,\alpha_s G^a_{\mu\nu}G_a^{\mu\nu} + C_g^{P}\bar\chi i\gamma_5\chi\,\alpha_s G^a_{\mu\nu}G_a^{\mu\nu}\\ &+\tilde C_g^{S}\bar\chi\chi\,\alpha_s G^a_{\mu\nu}\tilde G_a^{\mu\nu} + \tilde C_g^{P}\bar\chi i\gamma_5\chi\,\alpha_s G^a_{\mu\nu}\tilde G_a^{\mu\nu}\Big],\end{aligned} \tag{1}$$

where the Wilson coefficients C_i parameterize the effect of new physics associated with the scale Λ (organizing the interactions in this way assumes Λ to be much larger than the typical QCD scale of 1 GeV). To render the scalar and pseudoscalar matrix elements renormalization-scale invariant we included explicitly the quark masses m_q in the definition of the respective operators. We further assumed χ to be a Dirac fermion (in the Majorana case, $C_q^{VV} = C_q^{VA} = C_q^{TT} = 0$), and defined the dual field strength tensor as

$$\tilde G_a^{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\lambda\sigma}G^a_{\lambda\sigma}, \tag{2}$$

with sign convention $\epsilon^{0123} = +1$. Compared to the operator basis used in [23] we do not include the dimension-8 operators related to the traceless part of the QCD energy–momentum tensor. As shown in [23], these operators become relevant for heavy WIMPs and contribute to spin-independent interactions, decreasing significantly the single-nucleon contribution. Finally, we will ignore the tensor operators in (1) and concentrate on the chiral predictions for the V, A, S, P channels.

The kinematics for the WIMP–nucleon scattering process are taken as

$$N(p) + \chi(k) \to N(p') + \chi(k'), \tag{3}$$

the momentum transfer is defined as

$$q = k' - k = p - p', \qquad q^2 = t, \tag{4}$$

and the pion, η, nucleon, nucleus, and WIMP masses will be denoted by M_π, M_η, m_N, m_A, and m_χ, respectively (Dirac spinors are normalized to 1). We will also need

$$P = p + p', \qquad K = k + k'. \tag{5}$$

The cross section differential with respect to momentum transfer for the elastic WIMP–nucleus scattering process in the laboratory frame can be expressed as

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\mathbf{q}^2} = \frac{1}{8\pi v^2(2J+1)}\sum_{\text{spins}}|\mathcal{M}_{\text{NR}}|^2 + \mathcal{O}(\mathbf{q}^0), \tag{6}$$

with nucleus spin J, WIMP velocity v, and NR amplitude $\mathcal{M}_{\text{NR}}$ defined as

$$\mathcal{M} = 2m_A 2m_\chi\mathcal{M}_{\text{NR}} + \mathcal{O}(\mathbf{q}^2), \tag{7}$$

where $\mathcal{M}$ is the relativistic scattering amplitude. In the Majorana case, (6) receives an additional factor of 4.

3. Chiral power counting

We use the standard chiral power counting [24,25]

$$\partial = \mathcal{O}(p), \qquad m_q = \mathcal{O}(p^2), \qquad a_\mu, v_\mu = \mathcal{O}(p), \tag{8}$$

with axial-vector and vector sources a_μ and v_μ. The velocity distribution in dark matter halo models indeed suggests to count the momentum transfer $q \lesssim M_\pi$ as $\mathcal{O}(p)$ [10]. In the baryon sector we depart from the standard counting in chiral perturbation theory (ChPT) and adopt the more conventional ChEFT assumption (see, e.g., [26–28]) for the scaling of relativistic corrections

$$\frac{\partial}{m_N} = \mathcal{O}(p^2). \tag{9}$$

This counting is appropriate for a break-down scale around 500 MeV. As far as the WIMP is concerned, a chiral counting is only required for the NR expansion of the spinors. We assume the same counting as in the nucleon case, but display the corresponding additional powers explicitly. If $m_\chi \gtrsim m_N$, the suppression will be more pronounced, for $M_\pi \lesssim m_\chi \lesssim m_N$ the counting should be adapted, and for even smaller m_χ the naive counting breaks down.

For most of the channels it suffices to consider the leading-order Lagrangian to determine at which chiral order a given contribution starts. For the one-body matrix elements higher orders are subsumed into the nucleon form factors, which are obtained by their chiral expansion or could be taken from phenomenology. In this work, we consider all contributions up to $\mathcal{O}(p^3)$. Since the leading two-body terms start at $\mathcal{O}(p^2)$, this leaves the possibility that the next-to-leading-order (NLO) pion–nucleon Lagrangian involving the low-energy constants c_i [29] could be required, and this is indeed the case for the spatial component of the axial-vector current [11,12] (indicated by "2b NLO" in Table 1). In the same channel, NN contact terms d_i [30] enter. We define both c_i and d_i in the conventions of [21] (with dimensionless c_6 and c_7).

As a preview of our results, the leading chiral orders of one- and two-body currents for time and space components of the axial-vector and vector currents, as well as for the scalar and pseudoscalar operators, are listed in Table 1. The suppression by two powers ("+2") originating from the WIMP spinors is displayed separately. In the following sections, we give results for all one- and two-body currents involved in Table 1.

4. Nuclear matrix elements

4.1. Scalar

At zero momentum transfer the scalar couplings of the heavy quarks $Q = c, b, t$ can be determined from the trace anomaly of the QCD energy–momentum tensor [31]

$$\begin{aligned}&\theta^\mu{}_\mu = \sum_q m_q\bar q q + \frac{\beta_{\text{QCD}}}{2g_s}G^a_{\mu\nu}G_a^{\mu\nu}, \qquad \langle N|\theta^\mu{}_\mu|N\rangle = m_N,\\ &\frac{\beta_{\text{QCD}}}{2g_s} = -\left(11 - \frac{2N_f}{3}\right)\frac{\alpha_s}{8\pi} + \mathcal{O}(\alpha_s^2).\end{aligned} \tag{10}$$

For $N_f = 3$ active flavors, one obtains

$$\langle N|m_Q\,\bar Q\,Q\,|N\rangle = -\frac{\alpha_s}{12\pi}\langle N\big|G^a_{\mu\nu}G_a^{\mu\nu}\big|N\rangle = m_N f_Q^N, \tag{11}$$