Nuclear axial currents in chiral effective field theory

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One of the main advantage of the chiral effective field theory approach to the nuclear interaction is its ability to describe consistently nuclear forces *and* electroweak currents, since the latter are the Noether currents of chiral symmetry and, as such, strongly constrained by the same theoretical setting. Following the same scheme we adopted in the past for the case of the nuclear electromagnetic charge and current operators [1], we address in [2] the calculation of nuclear axial charge ρ_5^a and current $\mathbf{j}_5^a$ operators, with a the isospin index. An accurate theory of nuclear electroweak structure and dynamics is relevant in several areas of current interest. One such area is that of low-energy tests of physics beyond the Standard Model in β-decay experiments: phenomenologically, the weak interactions are known to couple only to left-handed neutrinos, and to violate parity maximally; deviations from these properties coming from new physics can in principle be detected, and must be interpreted having full control of the nuclear structure and weak interactions in nuclei. Furthermore, the low-energy inelastic neutrino scattering from nuclei is important in astrophysics and for neutrino detectors. Nuclear axial charge and current operators were originally derived up to one loop in heavy-baryon covariant chiral perturbation theory (HBChPT) by Park et al. [3]. Our framework is based instead on old-fashioned time-ordered perturbation theory, which allows for a clear identification of reducible and non-reducible contributions. According to Weinberg's prescription [4], the former have to be subtracted as they are generated by solving the Schroedinger or Lippman-Schwinger equation. Our formalism accounts for cancellations which occurr at a given order in the power counting between the contributions of irreducible diagrams and the contributions due to non-static corrections from energy denominators of reducible diagrams. Because of the different treatment of reducible diagrams we find differences compared to [3]. The method leads to nuclear operators which are not uniquely defined, due to an arbitrariness in the off-shell extension of the transition amplitude, but the resulting operators are nevertheless unitarily equivalent, and therefore the description of physical systems is not affected by this ambiguity. The weak axial charge and current operators at leading order consist of the single-nucleon contributions shown in Fig. 1, arising at order $O(Q^{-3}$ for the current and $O(Q^{-2})$ for the charge operator, with Q generically denoting the low-momentum scale. Two-

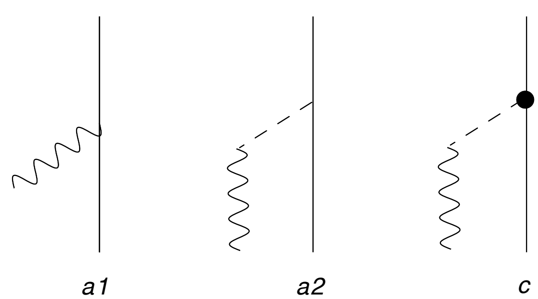

Figure 1. Diagrams a1 and a2 contribute to the one-body axial current operator at order $Q^{(-3)}$. Diagram c contributes to the one-body axial charge operator at order $Q^{(-2)}$. Nucleons, pions, and axial fields are denoted by solid, dashed, and wavy lines, respectively. Only a single time ordering is shown for diagrams a2 and c. The full dot in c stands for a $1/M_N$ correction to the πNN vertex.

body contributions from one- and two-pion exchange (OPE and TPE respectively) to the axial charge up to 1 loop are illustrated in Fig. 2. In addition there are four independent contact axial charge operators at order $O(Q)$, two of which absorb the divergencies of the loop diagrams, as evaluated in dimensional regularization.

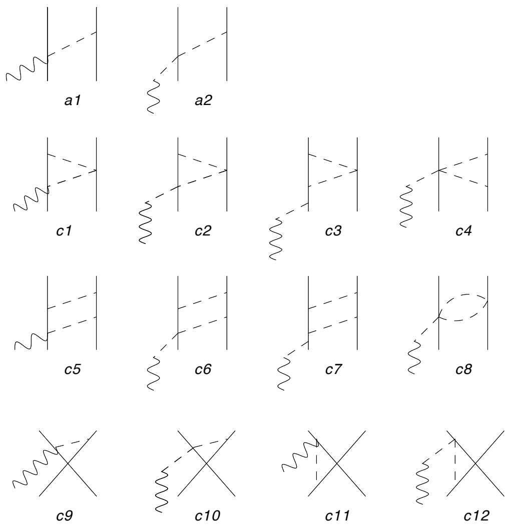

Figure 2. Diagrams contributing to the OPE axial charge at leading order Q^{-1} (panels a1 and a2), and to the TPE axial charge operator at order Q. Nucleons, pions, and axial fields are denoted by solid, dashed, and wavy lines, respectively. Only a single time ordering is shown for each topology.

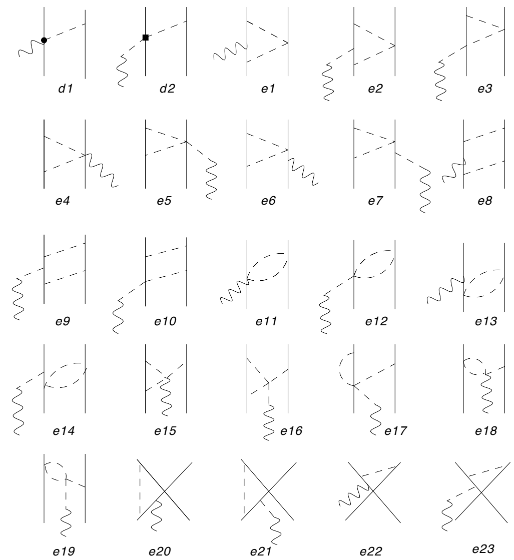

Figure 3. Diagrams contributing to the OPE axial current operator at order Q^0 and to the MPE axial current at order Q. Nucleons, pions, and axial fields are denoted by solid, dashed, and wavy lines, respectively. Only a single time ordering is shown for each topology.

For the axial current, we have a relativistic correction to the 1-body contributions arising at $O(Q^{-1})$, whose phenomenological relevance is well known e.g. for proton weak capture on ^{3}He, and two-body contributions at $O(Q^0)$ and $O(Q)$ displayed in Fig. 3. Only a single contact current operator arises at $O(Q^0)$, consistent with the fact that the loop contributions to the axial current are finite.

After renormalization of the pion and nucleon masses and fields, we find that all divergencies are absorbed by a redefinition of the contact low-energy constants (LECs). In particular the loop corrections to the OPE are renormalized by the subleading πN coupling constants, with the same values for the anomalous dimensions as obtained in the general heat-kernel formalism [5]. This provides a very non-trivial check of the consistency of our calculation.

We also checked that the current conservation,

$$\mathbf{q}\cdot\mathbf{j}_5^a = [H, \rho_5^a] \tag{1}$$

is fulfilled in the chiral limit ($m_\pi \to 0$) order by order in the low-energy expansion, where H represent the NN interaction calculated in the same chiral effective field theory setting, demonstrating once more the importance of using a consistent setting for interactions and currents. Studies for determining the involved LECs from phenomenology are in progress.

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