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Electroweak Current Operators in Chiral Effective Field Theory

Hermann Krebs[*] [†]
Ruhr-Universität Bochum
Institut für Theoretische Physik II
E-mail: hermann.krebs@rub.de

In this proceeding I briefly review current status of the construction of nuclear electro-weak currents within chiral effective field theory. I show that gauge and chiral symmetry requirements lead to the well-known continuity equations for the current and charge operators which, however, get modified at higher orders. Regularization of the current will be also discussed. I demonstrate that implementation of a cutoff regulator in a naive way leads to violation of chiral symmetry. To respect the underlying symmetries I propose to use higher derivative regularization in the nuclear forces and currents.

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1. Nuclear Chiral Effective Field Theory

Chiral effective field theory (EFT) is an effective field theory of quantum chromodynamics (QCD) which works in the energy sector where momenta of pions/nucleons are much lower than chiral symmetry breaking scale $\Lambda_\chi \sim 1\,\text{GeV}$. Relevant symmetries of QCD like e.g. chiral symmetry is by construction implemented in chiral EFT in a most general way. In the chiral EFT all the processes are described by point-like pions and nucleons which gain their structure in a perturbative way from loop corrections. Due to confinement this is an efficient way to proceed since in the low energy sector these are the observed degrees of freedom.

Chiral EFT has been successfully applied to meson and a nucleon sector in entirely perturbative way. In the two- and more-nucleon case, however, perturbative approach is not appropriate to describe nuclei. These are bound states of nucleons which can be interpreted as poles in the S-matrix and close to the poles any perturbation theory does obviously not converge. Almost three decades ago, Weinberg in his seminal papers suggested to use chiral perturbation theory to calculate an effective interaction between nucleons (known as nuclear forces). Bound state energies and scattering off nuclei can be approached numerically by solving Schröding equation in a non-perturbative way [1], see also [2] for a review on this topic. This path has been followed in the last three decade by several groups such that chiral nuclear forces have been worked out up to next-to-next-to-next-to-next-to-leading-order (N^4LO) in chiral expansion. Two-nucleon observables calculated with N^4LO forces are described with an excellent precision [3]. At the same time the number of fitted parameters in N^4LO forces is significantly reduced compared to phenomenological potentials [3] which clearly underlines the importance of two-pion-exchange contributions coming as prediction in the chiral EFT framework.

Within the same formalism one can calculate nuclear electroweak current operators consistent with the nuclear forces. The field was pioneered by Park et al. [4, 5] and was matured by two groups who calculated leading one-loop corrections to electroweak current operators up to N^3LO using two different methods to account for off-shell nuclear effects: the unitary transformation technique (UT) used by the Bochum-Bonn group [6],[7] and the framework of time-ordered perturbation theory (TOPT) used by the Pisa-JLab group [8, 9, 10, 11].

The proceeding is structured in three parts. In the first part I will briefly review our activities on the construction of the electroweak current operator calculated within UT formalism. In the second part I will compare our results with the results discussed by Pisa-Jlab group. In the third part the emphasis will be on the symmetry preserving regulator of the current. I will demonstrate that a naive multiplication of the current operators by a cutoff regulator and its convolution with chiral EFT wave functions of the deuteron leads to violation of chiral symmetry. This calls for consistent regularization of forces and currents which preserve underlying symmetries. Higher derivative regularization introduced by Slavnov in the early seventies [13] seems to be a promising solution.

2. Electroweak Current Operators within UT

Unitary transformation technique is a powerful tool to decouple pion-nucleon and purely nucleonic states in the Fock space reducing in this way a quantum field theoretic problem to a quantum mechanical one. In order to formulate the problem we denote by λ and η projection operators which project the states to the states with at least one pion and no pions, respectively. The Schrödinger equation in the presence of external sources can be rewritten into the form

$$\begin{pmatrix} \eta H \eta & \eta H \lambda \\ \lambda H \eta & \lambda H \lambda \end{pmatrix} \begin{pmatrix} \eta\,|\Psi\rangle \\ \lambda\,|\Psi\rangle \end{pmatrix} = i\frac{\partial}{\partial t}\begin{pmatrix} \eta\,|\Psi\rangle \\ \lambda\,|\Psi\rangle \end{pmatrix}. \tag{2.1}$$

The idea is to apply a unitary transformation on the Hamilton operator H in order to blockdiagonalize the matrix on the lhs of Eq. (2.1). The transformed Schrödinger equation gets the form

$$\left[U^\dagger H U + \left(i\frac{\partial}{\partial t}U^\dagger\right)U\right]U^\dagger|\Psi\rangle = i\frac{\partial}{\partial t}U^\dagger|\Psi\rangle. \tag{2.2}$$

We require

$$\eta\,\left(U^\dagger H U\right)_s\lambda = \lambda\,\left(U^\dagger H U\right)_s\eta = 0, \tag{2.3}$$

where

$$O_s = O|_{a=0,v=0,s=m_q,p=0}, \tag{2.4}$$

and O stays for any operator. Here m_q is a light quark mass and a,v,s,p denote external axial, vector, scalar, pseudoscalar sources, respectively. We denote a strong interacting part of the Hamiltonian by

$$W = \eta\,\left(U^\dagger H U\right)_s\eta. \tag{2.5}$$

Note that, although possible, we do not require the full operator in the rectangular bracket of Eq. (2.2) to be block-diagonal. It is enough that the strong interacting part of the Hamiltonian is block diagonal (see Eq. (2.3)). The reason is that we are not interested in the Hamilton operator in the presence of a cloud of external axial, vector or pseudoscalar sources. We are rather interested in a Hamiltonian in the presence of just one (or, not in this proceeding, maybe two) external sources. This drastically simplifies a quantum field theoretical problem even without full block-diagonalization.

Nuclear current operators can be extracted from first functional derivative of the rotated Hamiltonian. In momentum space e.g. vector, axial and pseudoscalar vector operators are defined by

$$\tilde{V}^j_\mu(k) = \left.\frac{\delta H_{\rm eff}}{\delta\tilde{v}^\mu_j(k)}\right|_s, \quad \tilde{A}^j_\mu(k) = \left.\frac{\delta H_{\rm eff}}{\delta\tilde{a}^\mu_j(k)}\right|_s, \quad \tilde{P}^j(k) = \left.\frac{\delta H_{\rm eff}}{\delta\tilde{p}_j(k)}\right|_s, \tag{2.6}$$

where effective Hamiltonian is

$$H_{\rm eff} = U^\dagger H U + \left(i\frac{\partial}{\partial t}U^\dagger\right)U, \tag{2.7}$$

and Fourier transformed sources are defined by [6]

$$X(x) = \int d^4q\, e^{-iq\cdot x}\tilde{X}(q), \quad X \in \left\{v^j_\mu, a^j_\mu, p^j\right\}. \tag{2.8}$$

Since $H_{\rm eff}$ is not block-diagonalized the current operators are also not block-diagonalized which means that even if in the initial state we have a purely nucleonic state in the final state we can have a state with zero, one, or even more pions. However, since in the practical calculations the currents will be convoluted with nuclear wave functions we only need to consider purely nucleonic initial and final states. Other states will be important e.g. if we are interested in Compton scattering where we deal with two current operators. In this case other matrix elements like effective pion-electroproduction matrix-element of the vector current $\langle NN|V_\mu^j(k)|\pi NN\rangle$ need to be worked out.

In the derivation of the current operator we use unitary transformations which explicitly depend on external sources and for this reason are time-dependent such that in general a time derivative of the unitary transformation is non-zero. This leads to explicit energy-transfer dependence of the currents and for this reason to a modification of continuity equations for axial and vector current operators:

$$\Big[W,\tilde{\mathbf{V}}_0(\vec{k},0)-\frac{\partial}{\partial k_0}\vec{k}\cdot\vec{\tilde{\mathbf{V}}}(\vec{k},k_0)+\frac{\partial}{\partial k_0}\big[W,\tilde{\mathbf{V}}_0(\vec{k},k_0)\big]\Big] = \vec{k}\cdot\vec{\tilde{\mathbf{V}}}(\vec{k},0), \tag{2.9}$$

$$\begin{aligned}\Big[W,\tilde{\mathbf{A}}_0(\vec{k},0)-\frac{\partial}{\partial k_0}\vec{k}\cdot\vec{\tilde{\mathbf{A}}}(\vec{k},k_0)+\frac{\partial}{\partial k_0}\big[W,\tilde{\mathbf{A}}_0(\vec{k},k_0)\big]+m_q\,i\frac{\partial}{\partial k_0}\tilde{\mathbf{P}}(\vec{k},k_0)\Big]&\\ = \vec{k}\cdot\vec{\tilde{\mathbf{A}}}(\vec{k},0)-m_q\,i\tilde{\mathbf{P}}(\vec{k},0),&\end{aligned} \tag{2.10}$$

see [6] for derivation of these expressions[1]. Here we, as usual, denote by bold letters matrix elements in isospin space

$$\mathbf{X} = \vec{X}\cdot\vec{\tau}, \tag{2.11}$$

where τ_i with $i=1,2,3$ are Pauli matrices in isospin space. Note the direct consequence form Eq. (2.9) is that the knowledge of the current in the Breit frame, where $k_0=0$ is valid, is not enough to check the continuity equations. One needs also an information about energy-transfer derivatives of the current operators.

As suggested by Weinberg [1] we can use chiral perturbation theory in order to calculate $H_{\rm eff}$. This was used to calculate nuclear forces $W-\eta H_0\eta$, where H_0 denotes a free Hamiltonian, and nuclear current operators V_μ^j, A_μ^j, P^j. A power counting, that tells which Feynman diagram belongs to which order in the chiral expansion, can be derived from a naive dimensional analysis. Denoting by

$$Q \sim \{p/\Lambda_b, M_\pi/\Lambda_b\}, \tag{2.12}$$

where p stays for small momenta, M_π for pion mass and $\Lambda_b\sim 600$ MeV is a breakdown scale of the chiral expansion, we can extract the chiral dimension ν of the corresponding diagram which counts as Q^ν from a naive dimensional analysis. For nuclear forces the chiral dimension of a connected diagram is given by

$$\nu = -2+\sum_i V_i\kappa_i, \tag{2.13}$$

[1] We assume in these expressions a linear dependence on the energy transfer. For more complicated energy-transfer dependence of the currents the continuity equations look more complicated.

while the chiral dimension of the nuclear charge and current operators is given by

$$\nu = -3 + \sum_i V_i \kappa_i \,. \tag{2.14}$$

Here κ_i denotes inverse mass dimension of the coupling constant at the vertex "i" and V_i denotes how many times the vertex "i" appears in a considered diagram. The inverse mass dimension can be expressed in terms of the chiral dimension of the vertex d, the number of nucleon fields a, the number of pion fields b and the number of external sources c

$$\kappa = d + \frac{3}{2}a + b + c - 4. \tag{2.15}$$

The leading order for nuclear forces starts with one-pion-exchange and contact interactions with $\nu = 0$ and are by now calculated up to $\nu = 5$ which is N^{4}LO. Note that there are no contributions to the nuclear forces at $\nu = 1$ and next-to-leading-order (NLO) contributions starts with $\nu = 2$. For the vector and axial vector currents the leading order starts from $\nu = -3$[2]. These are the charge operator of single nucleon vector current and a current operator of the single nucleon axial vector current. The calculations for vector and axial vector current have been performed up to $\nu = 1$ which are next-to-next-to-next-to-leading-order (N^{3}LO) calculations. Note that similar to the nuclear forces there are no contributions at the order $\nu = -2$ and for this reason NLO contribution shows up first at $\nu = -1$. In tables 1 and 2, 3 and 4 all possible contributions up to N^{3}LO are summarized for vector and axial vector operators. Note that the nucleon mass m is counted as

$$m \sim \Lambda_b^2 / p. \tag{2.16}$$

3. Electroweak Currents within UT vs TOPT

As already mentioned in the introduction, in parallel to our activities within UT techniques electroweak currents have been calculated within TOPT technique by Pisa-JLab group [8, 9, 10, 11]. The current operators in both calculations should agree with each other modulo unitary transformation. For vector currents it has been shown that there exists a unitary transformation which transforms UT currents into TOPT currents [11, 12]. The situation is more complicated for the axial vector currents. In this case the UT and TOPT results disagree even at the point of vanishing momentum transfer. An extensive discussion on this issue can be found in [16]. It remains to be seen in the future if the currents are unitary equivalent. If this transformation exists it should depend explicitly on the axial vector sources. The reason is that the TOPT current satisfies (at least in the chiral limit) an ordinary continuity equation [8]

$$\begin{aligned} \left[W, \tilde{\mathbf{V}}_0(\vec{k}, 0)\right] &= \vec{k} \cdot \vec{\tilde{\mathbf{V}}}(\tilde{\mathbf{k}}, \mathbf{0}), \\ \left[W, \tilde{\mathbf{A}}_0(\vec{k}, 0)\right] &= \vec{k} \cdot \vec{\tilde{\mathbf{A}}}(\tilde{\mathbf{k}}, \mathbf{0}) - \mathbf{m_q}\, \mathbf{i}\, \tilde{\mathbf{P}}(\tilde{\mathbf{k}}, \mathbf{0}), \end{aligned} \tag{3.1}$$

which means that TOPT currents do not depend on energy transfer. Since our currents do depend on the energy transfer they satisfy continuity equations in the form of Eq. (3.1) and can only be transformed to TOPT currents (if possible) with source dependent unitary transformations.

[2]For single nucleon contributions we need to subtract three chiral dimension due to the delta function of the spectator nucleon.

Table 1: Chiral expansion of the nuclear electromagnetic current operator up to N^{3}LO. LO, NLO, N^{2}LO and N^{3}LO refer to chiral orders Q^{-3}, Q^{-1}, Q^0 and Q, respectively. The single-nucleon contributions are given in Eqs. (2.7) and (2.16) of [7].

order	single-nucleon	two-nucleon	three-nucleon
LO	—	—	—
NLO	$\vec{\mathbf{V}}_{1\mathrm{N:static}}$ $+\vec{\mathbf{V}}_{1\mathrm{N:}1/\mathrm{m}}$	$\vec{\mathbf{V}}_{2\mathrm{N:}1\pi}$, Eq. (4.16) of [14]	—
N^{2}LO	$\vec{\mathbf{V}}_{1\mathrm{N:static}}$	—	—
N^{3}LO	$\vec{\mathbf{V}}_{1\mathrm{N:static}}$ $+\vec{\mathbf{V}}_{1\mathrm{N:}1/\mathrm{m}}$ $+\vec{\mathbf{V}}_{1\mathrm{N:off-shell}}$	$\vec{\mathbf{V}}_{2\mathrm{N:}1\pi}$, Eq. (4.28) of [14] $+\vec{\mathbf{V}}_{2\mathrm{N:}2\pi}$, Eq. (2.18) of [15] $+\vec{\mathbf{V}}_{2\mathrm{N:cont}}$, Eq. (5.3) of [14]	—

Table 2: Chiral expansion of the nuclear electromagnetic charge operator up to N^{3}LO. LO, NLO, N^{2}LO and N^{3}LO refer to chiral orders Q^{-3}, Q^{-1}, Q^0 and Q, respectively. The single-nucleon contributions are given in Eq. (2.6) of [7].

order	single-nucleon	two-nucleon	three-nucleon
LO	$\mathbf{V}^0_{1\mathrm{N:static}}$	—	—
NLO	$\mathbf{V}^0_{1\mathrm{N:static}}$	—	—
N^{2}LO	$\mathbf{V}^0_{1\mathrm{N:static}}$	—	—
N^{3}LO	$\mathbf{V}^0_{1\mathrm{N:static}}$ $+\mathbf{V}^0_{1\mathrm{N:}1/\mathrm{m}}$ $+\mathbf{V}^0_{1\mathrm{N:}1/\mathrm{m}^2}$	$\mathbf{V}^0_{2\mathrm{N:}1\pi}$, Eq. (4.30) of [14] $+\mathbf{V}^0_{2\mathrm{N:}2\pi}$, Eq. (2.19) of [15] $+\mathbf{V}^0_{2\mathrm{N:cont}}$, Eq. (5.6) of [14] $+\mathbf{V}^0_{2\mathrm{N:}1\pi,1/\mathrm{m}}$, Eq. (4.30) of [14]	$\mathbf{V}^0_{3\mathrm{N:}\pi}$, Eq. (4.1) of [7] $+\mathbf{V}^0_{3\mathrm{N:}\pi}$, Eq. (4.2) of [7] $+\mathbf{V}^0_{3\mathrm{N:cont}}$, Eq. (4.3) of [7]

4. Towards Consistent Regularization of the Currents

Sofar all reported calculations of current operators have been performed by using dimensional regularization. Naively one could take these operators and start to look at their expectation values in order to study observables. This is indeed what has been done by various calculation with TOPT currents, see e.g. [17] for a review. All these calculations should be considered as a hybrid approach where no claim on consistency between nuclear forces and currents is made. Even if both nuclear forces and currents are calculated from the same framework of chiral EFT the use of different

Table 3: Chiral expansion of the nuclear axial current operator up to N^{3}LO. LO, NLO, N^{2}LO and N^{3}LO refer to chiral orders Q^{-3}, Q^{-1}, Q^0 and Q, respectively. All equation references are understood to be from [6].

order	single-nucleon	two-nucleon	three-nucleon
LO	$\vec{\mathbf{A}}_{\rm 1N:\, static}$, Eq. (4.2)	—	—
NLO	$\vec{\mathbf{A}}_{\rm 1N:\, static}$, Eq. (4.7)	—	—
N^{2}LO	—	$\vec{\mathbf{A}}_{\rm 2N:\, 1\pi}$, Eq. (5.7) $+\,\vec{\mathbf{A}}_{\rm 2N:\, cont}$, Eq. (5.8)	—
N^{3}LO	$\vec{\mathbf{A}}_{\rm 1N:\, static}$, Eq. (4.46) $+\,\vec{\mathbf{A}}_{{\rm 1N:\, 1}/m{\rm ,UT'}}$, Eq. (4.13) $+\,\vec{\mathbf{A}}_{\rm 1N:\, 1/m^2}$, Eq. (4.18)	$\vec{\mathbf{A}}_{\rm 2N:\, 1\pi}$, Eq. (5.13) $+\,\vec{\mathbf{A}}_{\rm 2N:\, 1\pi, UT'}$, Eq. (5.23) $+\,\vec{\mathbf{A}}_{{\rm 2N:\, 1}\pi,\, 1/m}$, Eq. (5.19) $+\,\vec{\mathbf{A}}_{\rm 2N:\, 2\pi}$, Eq. (5.29) $+\,\vec{\mathbf{A}}_{\rm 2N:\, cont,\, UT'}$, Eq. (5.43) $+\,\vec{\mathbf{A}}_{\rm 2N:\, cont,\, 1/m}$, Eq. (5.41)	$\vec{\mathbf{A}}_{\rm 3N:\, \pi}$, Eq. (6.2) $+\,\vec{\mathbf{A}}_{\rm 3N:\, cont}$, Eq. (6.6)

Table 4: Chiral expansion of the nuclear axial charge operator up to N^{3}LO. LO, NLO, N^{2}LO and N^{3}LO refer to chiral orders Q^{-3}, Q^{-1}, Q^0 and Q, respectively. All equation references are understood to be from [6].

order	single-nucleon	two-nucleon	three-nucleon
LO	—	—	—
NLO	$\mathbf{A}^0_{\rm 1N:\, UT'}$, Eq. (4.4) $+\,\mathbf{A}^0_{{\rm 1N:\, 1}/m}$, Eq. (4.10)	$\mathbf{A}^0_{\rm 2N:\, 1\pi}$, Eq. (5.3)	—
N^{2}LO	—	—	—
N^{3}LO	$\mathbf{A}^0_{\rm 1N:\, static,\, UT'}$, Eq. (4.14) $+\,\mathbf{A}^0_{{\rm 1N:\, 1}/m}$, Eq. (4.22)	$\mathbf{A}^0_{\rm 2N:\, 1\pi}$, Eq. (5.14) $+\,\mathbf{A}^0_{\rm 2N:\, 2\pi}$, Eq. (5.30) $+\,\mathbf{A}^0_{\rm 2N:\, cont}$, Eq. (5.34)	—

regularizations (cut off vs dimensional regularization) leads to a chiral symmetry violation in the very first iteration of the current with nuclear forces. Here is the explanation:

In order to solve the Schrödinger equation nuclear forces have to be regularized. The usual way is to use the cutoff regularization. Let us for example choose a semi-local regulator discussed in [3]. The regularized form of the long-range part of the leading order nuclear force, which is one pion exchange diagram, is given by

$$V_{1\pi,\Lambda} = -\frac{g_A^2}{4F_\pi^2}\tau_1\cdot\tau_2\frac{\vec\sigma_1\cdot\vec q\,\vec\sigma_2\cdot\vec q}{q^2+M_\pi^2}e^{-\frac{q^2+M_\pi^2}{\Lambda^2}}, \tag{4.1}$$

where $\vec q$ denotes momentum transfer between two nucleons. The nice property of this regulator is that it does not affect long range part of the nuclear force at any power of $1/\Lambda$. On the other hand a pion-pole contribution proportional to g_A of the relativistic correction of the axial vector two-nucleon current is given by

$$\begin{aligned}\vec{\mathbf{A}}^{(Q,g_A)}_{2\mathrm{N}:1\pi,1/m} =& \frac{g_A}{8F_\pi^2 m}i\,\tau_1\times\tau_2\frac{\vec q_1\cdot\vec\sigma_1}{q_1^2+M_\pi^2}\frac{\vec k}{k^2+M_\pi^2}\Big[i\vec k\cdot\vec q_1\times\vec\sigma_2-\vec k_1\cdot\vec q_1+\vec k_2\cdot(\vec q_1+\vec k)\Big]\\ &+ 1\leftrightarrow 2,\end{aligned} \tag{4.2}$$

where $\vec k$ is the momentum transfer of the axial vector current, and other momenta are defined by

$$\vec q_i = \vec p_i\,'-\vec p_i,\quad \vec k_i = \frac{\vec p_i\,'+\vec p_i}{2},\quad i=1,2, \tag{4.3}$$

and momenta $\vec p_i\,'$ and $\vec p_i$ correspond to the final and initial momenta of the i-th nucleon, respectively. Note that this is not the only contribution to the relativistic corrections of the current, but only that which is proportional to g_A. Complete expression (including terms proportional to g_A^3) for the relativistic corrections can be found in [6]. After we regularized the nuclear force and the axial vector current we can perform the first iteration and take $\Lambda\to\infty$ limit:

$$\begin{aligned}&\vec{\mathbf{A}}^{(Q,g_A)}_{2\mathrm{N}:1\pi,1/m}\frac{1}{E-H_0+i\varepsilon}V_{1\pi,\Lambda}+V_{1\pi,\Lambda}\frac{1}{E-H_0+i\varepsilon}\vec{\mathbf{A}}^{(Q,g_A)}_{2\mathrm{N}:1\pi,1/m} =\\ &\Lambda\frac{g_A^3}{32\sqrt{2}\pi^{3/2}F_\pi^4}(\tau_1-\tau_2)\frac{\vec k}{k^2+M_\pi^2}\vec q_1\cdot\vec\sigma_1+1\leftrightarrow 2+\mathcal{O}(\Lambda^0).\end{aligned} \tag{4.4}$$

Since the one loop amplitude should be renormalizable there should exist a counter term which absorbs the linear singularity in Λ. From Eq. (4.4) we see that this should be a contact two-nucleon interaction with one pion coupling to it. However, there is no counter term like this in chiral EFT. Such counter term requires derivative-less coupling of the pion which is forbidden by the chiral symmetry: There exists only a counter term proportional to $\vec k\cdot\vec\sigma_1$, but there is none which is proportional to $\vec q_1\cdot\vec\sigma_1$. Here $\vec k$ is the momentum of the pion coupling to the two-nucleon interaction.[3] If there is no counter term which absorbs the linear cutoff singularity there should be some cancelation in the amplitude with other terms. Indeed the same singularity but with opposite

[3] At higher orders one can construct derivative-less pion-four-nucleon interactions by multiplying low energy constants with M_π^2. They are coming from explicit chiral symmetry breaking by finite quark mass. However, at the order Q we can not construct a counter term like this.

sign we would get for the static limit of the axial vector current of the order Q if we would calculate the current by using cutoff regularization. Axial vector current at the order Q, however, is calculated by using dimensional regularization and is finite. It also remains finite if we just multiply the current with any cutoff regulator we want. So at the level of the amplitude the mismatch between the cutoff and the dimensional regularization used in the construction of operators leads to a violation of chiral symmetry at one-loop level which, however, is the order of accuracy of our calculations. So we see that it is dangerous to multiply the current operators calculated within dimensional regularization by some cutoff regulator and calculate expectation values of this. With the similar arguments one can show that dimensionally regularized three-nucleon forces at the level of N^{3}LO, which were published in [20, 21], can not be used in combination with the cutoff regularized two-nucleon forces at the same order. The mismatch between dimensional and cutoff regularization will lead also in this case to a violation of the chiral symmetry at the one loop level.

In order not to violate the chiral symmetry we need to calculate both nuclear forces and currents with the same regulator. On top of it the regulator which we choose should be symmetry preserving. One possibility to construct a regulator, which manifestly respects the chiral symmetry, was proposed more than four decades ago by Slavnov [13], where he introduced a so called higher derivative regularization in a study of non-linear sigma model. Recently, first applications of this technique to the chiral EFT have been discussed in the literature [18, 19]. Construction of consistent nuclear forces and currents within a similar approach is work in progress.

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