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Quenching of nuclear matrix elements for $0\nu\beta\beta$ decay by chiral two-body currents

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We examine the leading effects of two-body weak currents from chiral effective field theory on the matrix elements governing neutrinoless double-β decay. In the closure approximation these effects are generated by the product of a one-body current with a two-body current, yielding both two- and three-body operators. When the three-body operators are considered without approximation, they quench matrix elements by about 10%, less than suggested by prior work, which neglected portions of the operators. The two-body operators, when treated in the standard way, can produce somewhat larger quenching. In a consistent effective field theory, however, these two-body effects become divergent and must be renormalized by a contact operator, the coefficient of which we cannot determine at present.

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Neutrinoless double-β ($0\nu\beta\beta$) decay is a still hypothetical process in which two neutrons decay to two protons and two electrons without emitting neutrinos [1]. Its discovery would show that neutrinos are their own antiparticles and could both pin down uncertain neutrino masses and discover entirely new particles. Experiments to observe the decay are thus growing in size and cost. Interpreting them, however, requires us to know the values of the nuclear matrix elements that figure in the decay rate via Fermi's golden rule. These cannot be measured, only calculated, and theorists have worked increasingly hard to compute them accurately; see Refs. [2,3] for reviews and, e.g., Refs. [4–11] for original work.

Because $0\nu\beta\beta$ decay has never been observed, one really ought to calculate its matrix elements from first principles with ingredients that allow an error estimate. The standard scheme for doing this is chiral effective field theory (EFT) [12]. Roughly speaking, one writes down all interactions among nucleons and pions that are consistent with spontaneously broken chiral symmetry. There are infinitely many of these but a power-counting scheme in nuclear momenta or the pion mass (both denoted by Q) divided by a QCD scale Λ near a GeV allows one to fit all the terms necessary to achieve any desired level of accuracy, at least in principle. The counting is not rigorous but usually works well.

The weak nuclear current can also be represented in this way. The leading piece involves the usual Gamow-Teller and Fermi operators associated with a single nucleon. Three orders down in the counting, two-body current operators appear [13]. Two-body axial weak currents are currently receiving a lot of attention because they appear [14] to mostly explain the longstanding tendency of nuclear theorists to overpredict single-β decay rates [15,16], which forces them to adopt an effective value for the axial-vector coupling constant g_A that is significantly smaller than the bare value. Recent suggestions [17] that g_A should exhibit similar quenching in $0\nu\beta\beta$ matrix elements, where it is squared and would thus have a larger impact, have led theorists to examine the effects of two-body current operators in $0\nu\beta\beta$ decay. Reference [18] was the first work on the issue. The authors and those of the later Ref. [19] normal ordered the two-body operators with respect to the noninteracting ground state of spin- and isospin-symmetric nuclear matter to obtain an effective density-dependent one-body current that quenched $0\nu\beta\beta$ matrix elements by roughly 30%, less than one might fear because the quenching was less effective when the virtual neutrino exchanged between nucleons in the process transferred a significant amount of momentum. The assumptions underlying the conclusions—that an effective one-body operator is sufficient and that normal ordering with respect to a simple nuclear-matter state is sufficient to obtain it—have never been examined, however.

Here we carry out a more comprehensive analysis. We construct the explicit product of the one-body and two-body current operators, the leading contribution from two-body currents to the $0\nu\beta\beta$ matrix element in the closure approximation (which in tests is accurate to 10% or so [20,21]), to obtain two- and three-body $0\nu\beta\beta$ operators. After an illustrative calculation in symmetric nuclear matter, we evaluate the matrix elements of these operators between reasonable approximations to full shell-model wave functions in ^{76}Ge and ^{76}Se, which have been used in many experiments; see, e.g., Ref. [22]. We find that the obvious sources of quenching, involving three nucleons (only two of which decay), have even smaller effects than the effective-operator approach suggests. Contributions from pairs of nucleons that both generate the two-body current and decay themselves turn out to be more problematic, however.

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In $0\nu\beta\beta$ decay the weak current acts twice. The nuclear matrix element that governs the decay is given by

$$M = \frac{4\pi R}{g_A^2} \int \frac{d\boldsymbol{x}_1 d\boldsymbol{x}_2 d\boldsymbol{q}}{(2\pi)^3} \frac{e^{i\boldsymbol{q}\cdot(\boldsymbol{x}_1-\boldsymbol{x}_2)}}{q(q+E_d)} \times \langle 0_F^+ | \hat{\mathcal{J}}^{\mu}(\boldsymbol{x}_1)\hat{\mathcal{J}}_{\mu}(\boldsymbol{x}_2) | 0_I^+ \rangle, \quad (1)$$

where $\hat{\mathcal{J}}(\boldsymbol{x})$ is the nuclear current, $R \equiv 1.2A^{1/3}$ fm is the nuclear radius, $g_A \approx 1.27$, $\boldsymbol{q}$ labels the momentum transfer, and $E_d \equiv \bar{E} - (E_I + E_F)/2$ is an average excitation energy to which the matrix element is not sensitive [23] ($\bar{E}$ is an absolute average energy). Up to third order in Q/Λ, the nuclear current $\hat{\mathcal{J}}^{\mu}$ can be written as $\hat{\mathcal{J}}^{\mu} = \hat{\mathcal{J}}^{\mu}_{1\text{b}} + \hat{\mathcal{J}}^{\mu}_{2\text{b}}$ where the two terms in the sum are the one- and two-body pieces of the current. The first of these is [13,24,25]

$$\hat{\mathcal{J}}^{\mu}_{1\text{b}}(\boldsymbol{x}) = \sum_{n=1}^{A} [\delta_{\mu 0} J_{n,0}(q^2) - \delta_{\mu j} \boldsymbol{J}_{n,j}(q^2)] \tau_n^- \delta(\boldsymbol{x} - \boldsymbol{r}_n). \quad (2)$$

Here $\boldsymbol{r}_n$ is the coordinate of the nth nucleon, $\boldsymbol{q} \equiv i\nabla$, and

$$J_{n,0}(q^2) = g_V + \cdots,$$
$$\boldsymbol{J}_n(q^2) = g_A \boldsymbol{\sigma}_n + i(g_M + g_V)\frac{\boldsymbol{\sigma}_n \times \boldsymbol{q}}{2m_N} - g_P(q^2)\frac{\boldsymbol{q}\boldsymbol{\sigma}_n \cdot \boldsymbol{q}}{2m_N} + \cdots, \quad (3)$$

where $g_V = 1$, $g_M \approx 3.706$, $g_P(q^2)$ is given, e.g., in Ref. [18], and m_N is the nucleon mass. In what follows, we will be looking at the axial current and so neglect contributions of $J_{n,0}(q^2)$. The terms indicated by ellipses can be shown [18] to contribute negligibly to the matrix element in Eq. (1).

In considering the two-body current, we neglect the term with coefficient c_6 [13] and terms with two-body pion poles [26] but otherwise keep the full momentum dependence of Ref. [13]). Fourier transforming Eqs. (A5) and (A6) of that paper with, following Ref. [27], an additional factor of $-1/4$ in the contact term gives the leading space piece of the axial two-body current operator in coordinate space,

$$\hat{\mathcal{J}}_{2\text{b}}(\boldsymbol{x}) = \sum_{k<l}^{A} \boldsymbol{J}_{kl}(\boldsymbol{x}), \quad (4)$$

$$\begin{aligned}\boldsymbol{J}_{kl}(\boldsymbol{x}) = \frac{2c_3 g_A}{m_N F_\pi^2} \left\{ m_\pi^2 \left[\left(\frac{\boldsymbol{\sigma}_l}{3} - \boldsymbol{\sigma}_l \cdot \hat{\boldsymbol{r}}\hat{\boldsymbol{r}} \right) Y_2(r) - \frac{\boldsymbol{\sigma}_l}{3} Y_0(r) \right] + \frac{\boldsymbol{\sigma}_l}{3}\delta(\boldsymbol{r}) \right\} \tau_l^- \delta(\boldsymbol{x} - \boldsymbol{r}_k) + (k \leftrightarrow l) \\ + \left(c_4 + \frac{1}{4} \right) \frac{g_A}{2m_N F_\pi^2} \left\{ m_\pi^2 \left[\left(\frac{\boldsymbol{\sigma}_\times}{3} - \boldsymbol{\sigma}_k \times \hat{\boldsymbol{r}}\boldsymbol{\sigma}_l \cdot \hat{\boldsymbol{r}} \right) Y_2(r) - \frac{\boldsymbol{\sigma}_\times}{3} Y_0(r) \right] + \frac{\boldsymbol{\sigma}_\times}{3}\delta(\boldsymbol{r}) \right\} \tau_\times^- \delta(\boldsymbol{x} - \boldsymbol{r}_k) + (k \leftrightarrow l) \\ - \frac{g_A}{4m_N F_\pi^2} [2\hat{d}_1(\boldsymbol{\sigma}_k \tau_k^- + \boldsymbol{\sigma}_l \tau_l^-) + \hat{d}_2 \boldsymbol{\sigma}_\times \tau_\times^-] \delta(\boldsymbol{r}) \delta(\boldsymbol{x} - \boldsymbol{r}_k), \end{aligned} \quad (5)$$

where $F_\pi = 92.4$ MeV is the pion decay constant, m_π is the pion mass, $\boldsymbol{r} = \boldsymbol{r}_k - \boldsymbol{r}_l$, and $\hat{\boldsymbol{r}} \equiv \frac{\boldsymbol{r}}{r}$. The Yukawa functions Y are $Y_0(r) = \frac{e^{-m_\pi r}}{4\pi r}$ and $Y_2(r) = \frac{1}{m_\pi^2} r \frac{\partial}{\partial r} \frac{1}{r} \frac{\partial}{\partial r} Y_0(r)$, and the compound spin and isospin operators are $\boldsymbol{\sigma}_\times = \boldsymbol{\sigma}_k \times \boldsymbol{\sigma}_l$ and $\tau_\times^- = (\tau_k \times \tau_l)^-$ [13]. The product of currents in Eq. (1) for the $0\nu\beta\beta$ matrix element can be broken up into contributions from one- and two-body currents. The leading piece, from two one-body currents acting as in diagram (a) of Fig. 1, is what has been considered almost exclusively in prior work. The first correction comes from diagrams such as (b) and (c), in which one of the one-body currents is replaced by a two-body current of either long range [diagram (b) with an internal pion] or short range [diagram (c)]. Reference [18] first considered these contributions but only with approximations that we avoid here.

To get an idea of what to expect in real nuclei, we begin with a more schematic discussion of nuclear matter, modeled after that in Ref. [18]. To simplify matters here (and only here), we neglect all but the d_1 and d_2 contact pieces of the two-body current [see Eq. (5)] and evaluate all the current operators at $q = 0$.

In nuclear matter, the one-body–two-body contributions just alluded to can be represented by the Goldstone diagrams in Fig. 2. The top row of diagrams, in which one nucleon in the two-body current is a spectator, was treated in Ref. [18]. The spectators contribute coherently, leading to a factor of the nuclear density in the matrix element and allowing one to replace the two-body current in the diagram by a density-dependent one-body effective current. Three-body operators need never be considered explicitly in such a procedure.

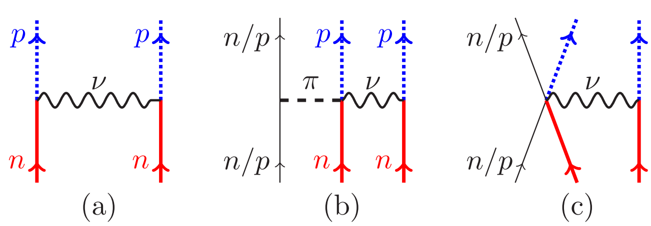

FIG. 1. $0\nu\beta\beta$ decay with electron lines omitted. Diagram (a) shows the leading contribution in which the one-body current acts twice, turning two neutrons (thick solid red lines) into two protons (dotted blue lines) via the exchange of a Majorana neutrino. Diagram (b) shows the action of the pion-exchange two-body current at one vertex; the thin solid black line on the left represents either a proton or a neutron. In diagram (c) the contact current replaces the pion-exchange current.

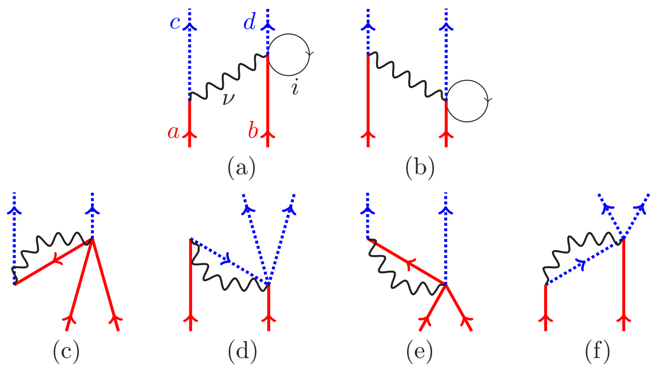

FIG. 2. Goldstone-diagram contributions to the $0\nu\beta\beta$ matrix element from the two-body current in symmetric nuclear matter. The thick solid red lines represent neutrons, the dotted blue lines represent protons, the thin black lines represent either neutrons or protons, and the wiggly black lines represent the exchanged neutrino. Upward-pointing arrows signify particles and downward-pointing arrows holes. The diagrams (a) and (b) in the top row represent the contributions considered in Ref. [18]. The diagrams (c)–(f) in the bottom row have not been considered before.

The bottom row has not been examined before. These diagrams contract creation and annihilation operators from different vertices and superficially are perhaps not as coherent. But the internal hole and particle lines are summed, and it is not obvious that their contributions will be much smaller. It *is* obvious, however, that diagrams (e) and (f) have the same sign as the top row of diagrams and that diagrams (c) and (d) have the opposite sign. A diagram's sign contains a factor of $S = (-1)^{n_h+n_l}$, where n_h (n_l) is the number of hole lines (nucleon loops). The diagrams in the top row have one hole line and one nucleon loop, and thus $S = 1$. Diagrams (e) and (f) have no hole lines or nucleon loops ($S = 1$), and diagrams (c) and (d) have one hole line and no nucleon loops ($S = -1$). The net effect once all terms are taken into account remains to be seen.

We evaluate the diagrams in the closure approximation, that is, by neglecting the variation in the energies of the intermediate particles and holes in the bottom row of diagrams. To simplify matters, we set E_d in Eq. (1) to zero so that the energy denominators contain just $1/q^2$ associated with the neutrino. We take the external momenta $\boldsymbol{k}_a$, $\boldsymbol{k}_b$, $\boldsymbol{k}_c$, and $\boldsymbol{k}_d$, which are to represent those of valence nucleons, to lie on the Fermi surface ($k = k_F$), although in evaluating the angle average of $1/|k_a - k_c|^2$ in the top row of diagrams we let the magnitude of one of the two momenta be distributed with equal probability in a symmetric interval of width k_F around the Fermi surface (to avoid a divergent result). With these assumptions, the amplitude represented by each of the diagrams has the form

$$X\delta(\boldsymbol{k}_a + \boldsymbol{k}_b - \boldsymbol{k}_c - \boldsymbol{k}_d)\,\langle f|\,\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2\tau_1^-\tau_2^-\,|i\rangle \tag{6}$$

for some constant X where the matrix element refers just to the spin-isospin part of the initial (i) and final (f) wave functions. We separately sum diagrams (a) and (b), (c) and (d), and (e) and (f) (the members of each pair are equal). The results are

$$\begin{aligned} X_{ab} &\equiv X_{(a)} + X_{(b)} \approx -\frac{2C(2+2\ln 2)k_F}{3\pi^2},\\ X_{cd} &\equiv X_{(c)} + X_{(d)} = \frac{3Ck_F}{4\pi^2} \approx -\frac{1}{2}X_{ab},\\ X_{ef} &\equiv X_{(e)} + X_{(f)} \approx -\frac{6C(\Lambda - k_F)}{4\pi^2} \approx 2X_{ab}, \end{aligned} \tag{7}$$

where C is a constant containing d_1, d_2, R, F_π, g_A, and m_N and where we take Λ, the momentum at which we cut off the integral over particle states, to be $3k_F$. The relative signs of the contributions reflect the discussion above. Avoiding the closure approximation—i.e., modifying the energy denominators to include the energies of the intermediate particle lines—would reduce the contributions of diagrams (e) and (f) by about 20%, but the integral would still grow with Λ. (The closure approximation is more drastic here than anywhere else because the intermediate-particle energies are bounded from above by Λ not by k_F. The states that carry these high intermediate-particle energies do not contribute strongly when only one-body currents are at play because the second current cannot act on the same nucleon as the first one.)

We can break the X's in Eq. (7) into contributions of three-body operators with $n \neq k, l$ in the products of the currents in Eqs. (1), (2), and (4), and two-body operators with $n = k$ or l. In addition to the producing the quenching contributions X_{ab} discussed in Ref. [18], three-body operators also contribute exactly twice X_{cd} so that the net quenching produced by the three-body operators nearly vanishes. Two-body operators produce $X_{ef} - X_{cd}$, which is about $5/2\ X_{ab}$ (a number, that, again, would be a bit smaller without closure) so the final overall quenching is greater than obtained in prior work. As we see next, conclusions much like these still hold when we use realistic nuclear wave functions, nucleon form factors, and the full two-body current.

One might argue that in computing X_{ef} we should not use a cutoff to regulate the integral. In a more consistent chiral effective field theory such as that in Refs. [28,29], in which all two-body processes are evaluated in isolation and the results subsequently embedded in a many-body calculation [so that Eq. (1) is not the starting point], that is standard practice; dimensional regularization restricts the momenta in loops to be low. But that procedure introduces counterterms with unknown coefficients at chiral orders below those considered here. We are simply trying to assess the quenching induced by two-body currents alone, and a cutoff simulates the effects of nucleon form factors in the sum over intermediate states in a realistic calculation. Of course, the use of a form factor to eliminate divergences, in conjunction with chiral currents, is not consistent; if we really want to do EFT we will require explicit counterterms. We return to this issue later.

First, however, we present realistic shell-model-like calculations ^{76}Ge and ^{76}Se in the usual $f_{5/2}pg_{9/2}$ oscillator valence space. Here, without the ability to include a complete set of intermediate-nucleus states, we need to work to evaluate the matrix elements of three-body operators. We do so by combining the three-body matrix elements of the operator $\hat{O}_{3b}$ [representing the three-body part of $\hat{\mathcal{J}}^\mu(\boldsymbol{x}_1)_{1b}\hat{\mathcal{J}}_\mu(\boldsymbol{x}_2)_{2b} +$

$\hat{\mathcal{J}}^{\mu}(\boldsymbol{x}_1)_{2b}\hat{\mathcal{J}}_{\mu}(\boldsymbol{x}_2)_{1b}]$ with three-body transition densities to obtain

$$M^{3b} = -\sum_{abcdef} \langle abc|\, \hat{O}_{3b}\, |def\rangle\, \rho^{3b}_{abc,def}, \tag{8}$$

where

$$\rho^{3b}_{abc,def} = \langle 0_F^+|\, a_a^\dagger a_b^\dagger a_c^\dagger a_d a_e a_f\, |0_I^+\rangle. \tag{9}$$

Here the subscripts $a, b, \dots$ represent full single-particle labels, e.g., a stands for the set $\{\tau_a, n_a, l_a, j_a, m_a\}$, i.e., the isospin, harmonic-oscillator radial quantum number, orbital angular momentum, total angular momentum, and z projection associated with the level in question. M^{3b} is thus the three-body piece of the matrix element M in Eq. (1). Other than the terms already mentioned, only those pieces of the product containing tensors in both the one-body and two-body currents are neglected.

We obtain the three-body matrix elements of $\hat{O}_{3b}$ in much the same way as Refs. [30–32] obtained those of three-body interactions, i.e. by first computing them in a large three-body Jacobi basis and then transforming to a coupled product basis. To get ρ^{3b} we use the generator coordinate method (GCM) to approximate shell-model wave functions [33]. As in Ref. [34], we use the Hamiltonian GCN2850 [4,35], take E_d to be 7.72 MeV, and include both axial deformation and an isoscalar pairing amplitude [36] as generator coordinates. We assume that the valence space sits atop an inert core of 56 filled oscillator orbitals. If all three nucleons acted on by $\hat{O}_{3b}$ are in the valence space, the densities ρ^{3b} are the matrix elements between the initial and the final GCM states of three creation and three annihilation operators. If one of the three nucleons comes from the shell-model core, on the other hand, then the ρ^{3b} reduce to simpler two-body valence-space transition densities. The corresponding contributions to M^{3b} are what one would obtain by normal ordering the product of currents with respect to the inert shell-model core, a more realistic version of the symmetric nuclear-matter state considered in Ref. [18]. The contractions generated by the normal ordering can be either between creation and annihilation operators within the two-body current as in the top row of Fig. 2 and in Ref. [18] or between operators from different currents as in the bottom row of Fig. 2.

Figure 3 shows the ratio M^{3b}/M^0, where M^0 is the leading part of the matrix element that comes from one-body currents at both vertices [Fig. 1(a)] for the decay of ^{76}Ge with the GCM wave functions described in the previous paragraph. These wave functions are not quite as complex as those in Ref. [34]; they are linear combinations of states with a single value for the isoscalar pairing amplitude and seven values for the axial deformation parameter β. The resulting matrix element—3.47—is reasonably close to the exact result of 2.81 [4]. The different panels in the figure correspond to different values for the couplings c_3 and c_4, and we present them as functions of $c_D \equiv d_1 + 2d_2$. The values of $c_3 = -3.2,\ c_4 = 5.4$ are from Ref. [37], the values of $c_3 = -4.78,\ c_4 = 3.96$ are from Ref. [38], and the values of $c_3 = -3.4,\ c_4 = 3.4$ are from Ref. [39]. To get the results on the left side of the figure (labeled same), we include only the contributions of contractions of creation and annihilation operators from within the same (two-body) current, such as those of Ref. [18] or diagrams (a) and (b) in Fig. 2. {Note, however, that Ref. [18] omitted a factor of $-1/4$ in the last line of Eq. (5).} We include all possible contractions to obtain the results on the right. The dashed and dotted lines show approximate results in which the we have discarded three-body operators that survive normal ordering with respect to the inert core (the discarded terms are those in which all three nucleons are in the valence shell) and with respect to an ensemble containing the full GCM ground states of ^{76}Ge and ^{76}Se, weighted equally. The ideas on which this ensemble normal ordering is based are presented in Ref. [40].

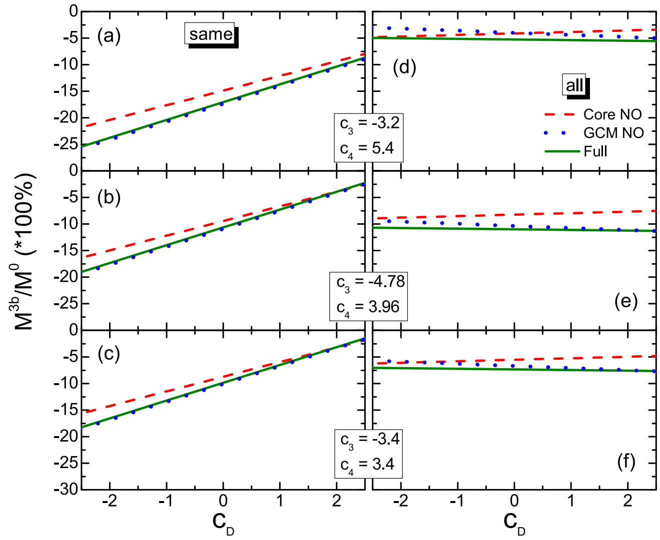

FIG. 3. Relative effects on the $0\nu\beta\beta$ matrix element from the three-body-operator parts of diagrams involving chiral two-body currents [as shown in Figs. 1(b) and 1(c), and Eq. (8) with several sets of coefficients c_3, c_4, and as a function of c_D for ^{76}Ge. The solid line represents the full results, the dashed line represents the approximate results when three-body operators are discarded after normal ordering with respect to the intercore, and the dotted line represents the results when the normal ordering is with respect to an ensemble containing the GCM ^{76}Ge and ^{76}Se ground states. The results in the panels on the left include only contributions from the contraction of creation an annihilation operators at the same vertex in Fig. 1. See the text for details.

The figure shows that with only the contractions from within the two-body current, the three-body operators quench the matrix element by 5%–25% for $|c_D| \leqslant 2$. This level of quenching is what one would obtain with the density-dependent effective-operator treatment of Ref. [18] at a somewhat lower nuclear density than that used there. A similar level of quenching holds in single-β decay as discussed in Ref. [14]. When all the contractions are included, the quenching decreases, just as in our nuclear-matter results for the contact part of the current. In the bottom two panels it does not decrease very much, but in the top panel, it decreases significantly. The full results are also nearly independent of c_D, bearing out the almost complete cancellation between the different three-body contractions we found in the nuclear-

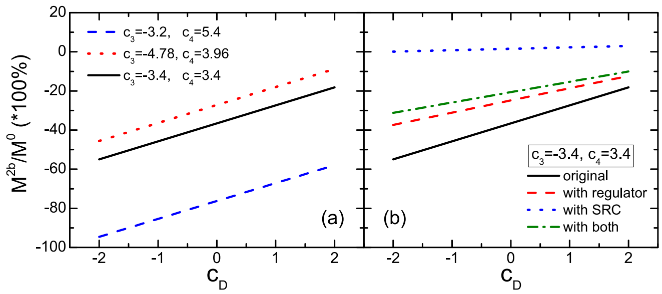

FIG. 4. Relative effects on the $0\nu\beta\beta$ matrix element from the two-body-operator parts of diagrams involving chiral two-body currents [as shown in Figs. 1(b) and 1(c)] as a function of c_D. The results in panel (a) are for three sets of values for c_3 and c_4. Panel (b) shows the effects of several modifications to the operator, both separately and when combined. See the text for details.

matter calculation. When all is said and done, the three-body operators end up quenching matrix element by 5 or 10%.

A final observation regarding Fig. 3: the normal ordering with respect to the inert core indeed provides most of the matrix element with the configurations in which all three nucleons are in the valence shell contributing relatively little. The realistic reference ensemble usually makes the normal ordering even better. That is good news for many-body calculations in which three-body operators are problematic.

We turn finally to the troublesome two-body operators in the product of one-body and two-body currents. As already noted, without nucleon form factors or other regulators the loops that produce these operators cause divergences. The operator that comes from the contact current, for example, is

$$\hat{O}^{2b}_{c_D} = \frac{2c_D R}{\pi m_N F_\pi^2} \sum_{k\neq l}^{A} \int dq \frac{\left[q g_A^2(q^2) - \frac{q^3 g_A(q^2) g_P(q^2)}{6m_N}\right]}{(q+E_d) g_A^2} \times \boldsymbol{\sigma}_k \cdot \boldsymbol{\sigma}_l \tau_k^- \tau_l^- \delta(\boldsymbol{r}). \quad (10)$$

The integral diverges if g_A has no q dependence. Here, for the purposes of estimation, we assign the dipole nucleon form factors given in Ref. [25] and used in nearly every prior calculation. Figure 4(a) shows the relative effects on the nuclear matrix elements from all the two-body operators and with the GCM wave functions described earlier (again for several chiral interactions). These operators can quench the matrix element substantially by up to 90% in the figure. The amount of quenching, however, is very sensitive to c_3, c_4, and c_D and can be much less. Furthermore, the quenching is due almost entirely to the zero-range parts of the two-body operators in Eq. (10) and in the analogous contact associated with pion exchange. (The longer-range pion-exchange Yukawa functions have very little effect.) As a result, modifications to physics at short distances, omitted from the matrix elements in panel (a) but typically included in calculations such as ours, have significant effects. Panel (b) shows what happens when we include them. A consistent EFT calculation requires regulation; the two-body regulator from Ref. [13] with $\Lambda =$ 500 MeV smears out the contact terms in the operator and decreases the quenching from the two-body matrix element by about a third. An explicit short-range correlation function is needed if the model space omits high-momentum states; the figure indicates that a Jastrow function of the "Argonne" type from Ref. [41]) drastically reduces the quenching. When the regulator and short-range correlation function are used together, the latter has a much smaller effect because of the smearing by the regulator. It is not, of course, consistent to use a short-range correlation function from a calculation with the Argonne potential in chiral EFT, but the regulator makes the precise form of the correlation function irrelevant.

What is the meaning of these various results? One might take the final dot-dashed curve in panel (b) to be a rough estimate of the quenching from two-body operators, but because those contributions diverge without form factors or a cutoff, a consistent calculation will contain additional short-range contributions from a counterterm. Unfortunately, the coefficient of that term is unknown with no obvious way to fix it from data. Only if it is small, i.e., if short-range repulsion fully and faithfully represents the effects induced by high-energy virtual neutrinos, will the dot-dashed line represent reality.

Interestingly, the counterterm is already a part of the $\beta\beta$ EFT of Ref. [28] where it occurs one order below that of the two-body currents. Within a cutofflike scheme, such as ours, that is the order required to cancel the divergent loops. With more typical dimensional regularization, however, the two-body operators in the product of currents, after removal of the divergence, would naturally contribute at most at the same order as the two-body currents themselves. We might even expect them to have less of an effect than the three-body operators in the product because the factor of Q from the low momenta in loops is a little smaller than k_F, which one obtains from the third particle in the three-body operators. But even with dimensional regularization one would need a counterterm at higher order than with cutoff regulation but with an equally unknown coefficient to remove the divergence.

How can we determine the unknown coefficient? The same question arises for a coefficient at leading order according to Ref. [29]. In either case, the coefficient can, in principle, be fixed through a calculation in lattice QCD or from data on pion double-charge exchange [42]. Either possibility is difficult to realize, however. Until a lattice calculation becomes feasible, we may have to resort to models to provide estimates that will be hard to verify.

The unknown coefficient does not weaken our conclusions about the three-body operators in the product of currents. They probably quench matrix elements by about 10%—we say probably because of the potential effects of pion poles, higher-order contributions in χEFT, etc., which eventually should be examined carefully. A 10% quenching is less than previous work suggests and nearly independent of c_D. Furthermore, in the future we can compute the effects of these operators to a good approximation by discarding all but the normal-ordered two-body pieces. All in all, many-body currents in χEFT are unlikely to produce really severe quenching. To say more, we will need a way to determine the coefficients of two-body counterterms.

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