Chiral two-body currents and neutrinoless double-β decay in the quasiparticle random-phase approximation

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(Received 30 March 2014; revised manuscript received 19 May 2014; published 13 June 2014)

We test the effects of an approximate treatment of two-body contributions to the axial-vector current on the quasiparticle random-phase approximation (QRPA) matrix elements for neutrinoless double-beta decay in a range of isotopes. The form and strength of the two-body terms come from chiral effective-field theory. The two-body currents typically reduce the matrix elements by about 20%, not as much as in shell-model calculations. One reason for the difference is that standard practice in the QRPA is to adjust the strength of the isoscalar pairing interaction to reproduce two-neutrino double-beta decay lifetimes. Another may be the larger QRPA single-particle space. Whatever the reasons, the effects on neutrinoless decay are significantly less than those on two-neutrino decay, both in the shell model and the QRPA.

DOI: 10.1103/PhysRevC.89.064308 PACS number(s): 21.60.Jz, 23.40.Hc

I. INTRODUCTION

The observation of neutrinoless double-beta ($0\nu\beta\beta$) decay would mean that neutrinos are Majorana particles. It would also tell us the overall neutrino mass scale if the nuclear matrix elements that help govern the decay could be calculated with sufficient accuracy. At present, the matrix elements from reasonable calculations differ from one another by up to factors of 3. The true uncertainty might be larger if there is physics that none of the calculations capture.

One familiar source of uncertainty is the way in which the axial-vector coupling constant is "renormalized in medium." The renormalization has several apparent sources, not all of which are directly connected to the weak interactions. Truncation of the many-nucleon Hilbert space, for instance, appears to reduce the matrix elements of spin operators, whether or not they stem from weak interactions [1]. But a separate source affects weak currents themselves: many-body operators that arise because nucleons are effective degrees of freedom. The many-body currents reduce matrix elements as well, though by amounts that are still in dispute [2].

Here, we focus on two-body currents in a very particular framework: chiral effective field theory (χEFT) in combination with the quasiparticle random-phase approximation (QRPA). We build on several published papers. References [3] and [4] extract coefficients of the χEFT interaction, and Refs. [5] and [6] present the form of the χEFT one- and two-body weak currents that follow. References [7] and [8] use an isospin-symmetric Fermi gas model to substitute approximate effective one-body current operators for the complicated two-body operators, and use the renormalized one-body current to correct shell-model calculations of double-beta decay (in a particular limit that we describe later) [7] and WIMP-nucleus scattering (in more generality) [8]. We will use the effective one-body operators from both references to calculate double-beta matrix elements in the QRPA.

There is some reason, before beginning, to believe that the effects of two-body currents will be smaller in the QRPA as usually applied than in the shell model. The correlations in the QRPA are simpler than in the shell-model and in recent years QRPA practitioners have compensated by fitting a parameter in the interaction—the strength of the isoscalar particle-particle interaction—to reproduce measured two-neutrino double-beta ($2\nu\beta\beta$) decay rates. Any change in the axial current operator is compensated to retain the correct $2\nu\beta\beta$ matrix element, and the compensation should carry over, in some measure at least, to $0\nu\beta\beta$ decay. We shall see below the degree to which that occurs.

The remainder of this paper is structured as follows: Section II describes the ingredients of our calculation, including the new two-body currents. Section III displays our results for $0\nu\beta\beta$ matrix elements in a wide range of isotopes. Section IV is a conclusion.

II. METHODS

A. One- and two-body currents

In χEFT, interactions and currents for nucleon and pion degrees of freedom are expanded in powers of momentum transfer p divided by a breakdown scale $\Lambda_\chi \approx 500$ MeV. Following Ref. [7], we equate $\mathcal{O}(p/m)$, where m is the nucleon mass, with $\mathcal{O}((p/\Lambda_\chi)^2)$. The interactions and currents should be derived consistently, either through fitting or by matching onto the predictions of QCD. We will not use a χEFT interaction, but can still do an approximately correct calculation by using χEFT currents. Of course those currents will comprise operators for three, four, ... nucleons and there is no guarantee that the higher-order terms in the chiral expansion that generates these operators will be small in a many-body

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system. Truncating the expansion for the interaction at low order yields reasonable results, however (see, e.g., Ref. [9]), and it is worth exploring a similar approximation in the currents.

In a nonrelativistic framework, one can write a general one-body current in the form

$$J^{0\dagger}(\mathbf{r}) = \sum_{i=1}^{A} J^0_{i,1b}\delta(\mathbf{r}-\mathbf{r}_i)\tau_i^+,$$
$$\mathbf{J}^{\dagger}(\mathbf{r}) = \sum_{i=1}^{A} \mathbf{J}_{i,1b}\delta(\mathbf{r}-\mathbf{r}_i)\tau_i^+. \tag{1}$$

where an operator with subscript i acts only on the ith nucleon and τ^+ changes a neutron into a proton. To third order in the counting, the one-body charge-changing weak current operators can be written as [7]

$$J^0_{i,1b} = g_V(p^2) - g_A(0)\frac{\mathbf{P}\cdot\boldsymbol{\sigma}_i}{2m} + g_P(p^2)\frac{E\boldsymbol{\sigma}_i\cdot\mathbf{p}}{2m},$$
$$\mathbf{J}_{i,1b} = -g_A(p^2)\boldsymbol{\sigma}_i + g_P(p^2)\frac{\mathbf{p}(\boldsymbol{\sigma}_i\cdot\mathbf{p})}{2m}$$
$$\quad - i(g_M(0)+g_V(0))\frac{\boldsymbol{\sigma}\times\mathbf{p}}{2m} + g_V(0)\frac{\mathbf{P}}{2m}, \tag{2}$$

where $\mathbf{p} = \mathbf{p}_i - \mathbf{p}'_i$, $\mathbf{P} = \mathbf{p}_i + \mathbf{p}'_i$, $E = E_i - E'_i$, and, to the same order in the chiral expansion,

$$g_V(p^2) = 1 - 2\frac{p^2}{(850\text{ MeV})^2},$$
$$g_A(p^2) = g_A(0)\left(1 - 2\frac{p^2}{(1040\text{ MeV})^2}\right), \quad g_A(0) = 1.27,$$
$$g_P(p^2) = 2\frac{g_{\pi pn}F_\pi}{p^2+m_\pi^2} - \frac{4mg_A(0)}{1040\text{ MeV}^2}, \quad g_M(0) = 3.70 \tag{3}$$

with $F_\pi = 92.4$ MeV, $m_\pi = 138.04$ MeV, and $g_{\pi pn} = 13.05$. In all these expressions $\mathbf{p}_i$ and $\mathbf{p}'_i$ stand for $-i\nabla_i$ acting on the left and right of the delta functions in Eq. (1).

As mentioned in the introduction, we will use two separate approximations schemes for two-body currents, one presented in Ref. [7] for $0\nu\beta\beta$ matrix elements in the shell model and an improved version presented by the same group in a paper on spin-dependent WIMP-nucleus scattering [8]. The first approximation scheme neglects the difference between the momentum transfers to the two nucleons. We use it nonetheless because it is the only scheme applied so far to double-beta decay and we wish to compare our results with those of Ref. [7]. The two schemes will turn out to yield only minor differences. Occasional disparities between the Gamow-Teller matrix elements in the two formulations are largely compensated by disparities in the tensor matrix elements.

Both schemes involve an effective correction to the one-body current through the assumption that one of the two nucleons in the two-body current lies in a spin-and-isospin symmetric core. The resulting approximation is crude but probably reasonable. Reference [7] neglects tensor-like terms in the current, leading to a renormalization of g_A but not g_P. References [8] does a more complete calculation that leads to a separate renormalization of g_P. Here we write explicitly only the effective current of [7]:

$$\langle\mathbf{p_i}|\,\mathbf{J}^{\text{eff}}_{i,2b}(\mathbf{r})\,|\mathbf{p_i'}\rangle$$
$$= -g_A(p^2)\boldsymbol{\sigma}_i\left\{\frac{\rho}{F_\pi^2}\left[\frac{c_D}{g_a\Lambda_\chi} + \frac{2}{3}c_3\frac{p^2}{4m\pi^2+p^2}\right.\right.$$
$$\left.\left. + I(\rho,P)\left(\frac{1}{3}(2c_4-c_3)+\frac{1}{6m}\right)\right]\right\}e^{-i\mathbf{p}\cdot\mathbf{r}}. \tag{4}$$

with

$$I(\rho,P) = 1 - \frac{3m_\pi^2}{2k_F^2} + \frac{3m_\pi^3}{2k_F^3}\,\text{acot}\left[\frac{m_\pi^2+\frac{P^2}{4}-k_F^2}{2m_\pi k_F}\right]$$
$$+ \frac{3m_\pi^2}{4k_F^3 P}\left(k_F^2+m_\pi^2-\frac{P^2}{4}\right)$$
$$\times \ln\left[\frac{m_\pi^2+\left(k_F-\frac{P}{2}\right)^2}{m_\pi^2+\left(k_F+\frac{P}{2}\right)^2}\right]. \tag{5}$$

In these equations k_F is the Fermi momentum and P is the center-of-mass momentum of the decaying nucleons, which can be set to zero without altering $I(\rho,P)$ significantly [7]. The constants c_3, c_4, and c_D are the χEFT parameters, fit to data in light nuclei. Their values depend on how the fit is carried out.

The above can be captured by defining new *effective* one-body current operators $\mathcal{J}^{\mu\dagger}$ as the operators $J^{\mu\dagger}$ from Eq. (1) but with the factor $g_A(p^2)$ multiplying $\boldsymbol{\sigma}_i$ in Eq. (3) replaced by an effective coupling $g_A^{\text{eff}}(p^2)$, given by

$$g_A^{\text{eff}}(p^2) = g_A(p^2)\left\{1 - \frac{\rho}{F_\pi^2}\left[\frac{c_D}{g_a\Lambda_\chi} + \frac{2}{3}c_3\frac{p^2}{4m_\pi^2+p^2}\right.\right.$$
$$\left.\left. + I(\rho,0)\left(\frac{1}{3}(2c_4-c_3)+\frac{1}{6m}\right)\right]\right\}, \tag{6}$$

where P has been set to zero.

The treatment of WIMP scattering in Ref. [8] is more complete and involves much longer expressions. We refer the reader there for details on the renormalization of both g_P and g_A.

B. Decay matrix elements

The two kinds of double-beta decay—$2\nu\beta\beta$ and $0\nu\beta\beta$—transfer very different amounts of (virtual) momentum among nucleons. Two-neutrino decay is simply two successive virtual beta decays, with very little momentum transfer. Its matrix element can be written with excellent accuracy as

$$M'^{2\nu} = \left(\frac{g_A^{\text{eff}}(0)}{g_A(0)}\right)^2 \sum_{N,i,j}\frac{\langle F|\,\boldsymbol{\sigma}_i\tau_i^+\,|N\rangle\cdot\langle N|\,\boldsymbol{\sigma}_j\tau_j^+\,|I\rangle}{E_N - \frac{E_I+E_F}{2}}, \tag{7}$$

where the $|N\rangle$ are states in the intermediate nucleus with energy E_N, and $|I\rangle$ and $|F\rangle$ are the initial and final nuclear ground states, with energies E_I and E_F. This matrix element and the neutrinoless version to follow differ from the unprimed $M^{2\nu}$ (and $M^{0\nu}$) used elsewhere in that g_A is always set to $g_A(0) = 1.27$ (and not to some effective value) in the phase space factor multiplying the matrix element, so that all effects of g_A modification are in the matrix element itself.

Neutrinoless decay, in contrast to its two-neutrino counterpart, creates a virtual neutrino that typically carries 100–200 MeV of momentum. The expression for its matrix element involves an integral over all neutrino momenta:

$$M'^{0\nu} = \frac{R}{2\pi^2 g_A(0)^2} \sum_N \int d^3x\, d^3y\, d^3p\, e^{i\mathbf{p}\cdot(\mathbf{x}-\mathbf{y})} \times \frac{\langle F|\, \mathcal{J}^{\mu\dagger}(\mathbf{x})\,|N\rangle\, \langle N|\, \mathcal{J}^{\dagger}_{\mu}(\mathbf{y})\,|I\rangle}{p\left(p+E_N-\frac{E_I+E_F}{2}\right)}, \quad (8)$$

where R is the nuclear radius, inserted to make the matrix element dimensionless. Details on the evaluation of this still rather abstract expression appear, e.g., in Refs. [10] and [11]. The important point is that the $0\nu\beta\beta$ matrix element depends on $g_A^{\rm eff}(p^2)$ (and in the formulation of Ref. [8] on $g_P^{\rm eff}(p^2)$ as well) because of the two current operators $\mathcal{J}^{\mu\dagger}$, defined just above Eq. (6), and the integral over momentum.

III. RESULTS

The values of the parameters c_3, c_4, and c_D come from fits to data in systems with very few nucleons. They depend on details of the fitting procedure; for this reason Ref. [7] gives several sets of possible values. It also evaluates the effective one-body current for a range of Fermi-gas densities ρ (the gas represents the nuclear core) because the nuclear density, though roughly constant in the nuclear interior, is not exactly so. As a result, it finds a range of final shell-model $0\nu\beta\beta$ matrix elements, with the correct one probably somewhere within the range. Here we will use c parameters and densities at the extremes of the reasonable range to set probable upper and lower limits on the effects of two-body currents.

Tables I and II present the results of our calculations. The headings a, b, c, and d in these tables refer to various prescriptions for fixing the χEFT parameters (see caption). The last column averages the quenching of the $0\nu\beta\beta$ matrix element over these entries, leading to a mean effect of about 20%, either with the parametrization of Ref. [7] or the more complete one in Ref. [8]. Figure 1 summarizes the same results, comparing $M'^{0\nu}$ with one-body and one-plus-two-body currents for all the nuclei we consider.

The minimum quenching from the set of choices in Ref. [7] occurs when $c_3 = -3.2$, $c_4 = 5.4$, and $c_D = 0$ [3], a combination we call EM below (following Ref. [7]) while the maximum quenching corresponds to $c_3 = -2.4, c_4 = 2.4$, and $c_D = 0$ [4], a combination we call EGM+δc_i. (The units of the c parameters are GeV^{-1}.) These choices result in values for $g_A^{\rm eff}(0)/g_A$ in a range 0.66–0.85 that brackets the empirical value of that ratio derived from the analysis of ordinary Gamow-Teller beta decay; see, e.g., Refs. [12] and [13]. Reference [14] points out that single-beta and two-neutrino double-beta decay observables can be described simultaneously in the QRPA with $g_A^{\rm eff}(0)/g_A$ in that range, implying that two-body currents can completely account for the renormalization of g_A. On the other hand, older meson-exchange models [1] suggest that the effects of many-body currents on allowed beta decay are small. The source of the disagreement between the strength of two body currents in χEFT and exchange models is not completely clear to us.

The degree of quenching here is noteworthy for two additional reasons. First, the $0\nu\beta\beta$ quenching is much less than its $2\nu\beta\beta$ counterpart, which with the same currents is closer to 40% or 50%. Second, it is noticeably less then the quenching of $0\nu\beta\beta$ decay in the shell model, which from the figures in Ref. [7], appears to be about 30%. There are several reasons for the first fact. As Ref. [7] shows, the degree of quenching decreases with increasing momentum transfer. An as we noted earlier, $2\nu\beta\beta$ decay involves almost no momentum transfer by the currents, while $0\nu\beta\beta$ decay involves momentum transfers that are typically 100–200 MeV

TABLE I. The $0\nu\beta\beta$ matrix element $M'^{0\nu}$ with one- and two-body nucleon current operators from the text and the Argonne V18 G-matrix-based QRPA. We use several sets of values for the χEFT parameters and two nuclear densities, and both the simplest and more complete versions of the effective one-body current, from Refs. [7] and [8] respectively. $\langle M'^{0\nu}\rangle$ is the matrix element averaged over these possibilities; its variance is in parentheses. The columns labeled "a" through "d" correspond to different EFT-parameter choices (defined in Ref. [7]) and nuclear-density choices. These choices are a: EGM+δc_i, $\rho = 0.10$ fm^{-3}; b: EGM+δc_i, $\rho = 0.12$ fm^{-3}; c: EM, $\rho = 0.10$ fm^{-3}; d: EM, $\rho = 0.12$ fm^{-3}. The last column contains the percent suppression ε of $\langle M'^{0\nu}\rangle$ with respect to the value without two-body currents (displayed in the first column).

Nucleus	$M'^{0\nu}$ 1bc	$M'^{0\nu}$ 2bc								$\langle M'^{0\nu}\rangle$ with quenching	ε (%)
		Parameters of Ref. [7]				Parameters of Ref. [8]					
		a	b	c	d	a	b	c	d		
^{48}Ca	0.684	0.641	0.629	0.580	0.558	0.637	0.637	0.596	0.592	0.61(0.03)	11
^{76}Ge	5.915	5.121	4.932	4.369	4.084	5.050	4.914	4.412	4.206	4.64(0.41)	22
^{82}Se	5.313	4.570	4.393	3.863	3.583	4.506	4.378	3.906	3.701	4.11(0.39)	23
^{96}Zr	3.224	2.999	2.913	2.636	2.506	2.946	2.894	2.651	2.573	2.76(0.19)	14
^{100}Mo	6.287	5.552	5.370	4.801	4.510	5.437	5.314	4.813	4.618	5.05(0.41)	20
^{110}Pd	6.575	5.795	5.607	5.037	4.758	5.673	5.540	5.030	4.833	5.28(0.41)	20
^{116}Cd	4.485	3.894	3.754	3.342	3.127	3.812	3.701	3.331	3.126	3.51(0.31)	22
^{124}Sn	3.974	3.599	3.511	3.231	3.118	3.521	3.464	3.211	3.143	3.35(0.19)	16
^{130}Te	4.610	4.031	3.890	3.445	3.216	3.949	3.855	3.465	3.313	3.65(0.32)	21
^{136}Xe	2.570	2.249	2.169	1.920	1.791	2.190	2.136	1.915	1.829	2.02(0.18)	21

TABLE II. The same as Table I, but for the CD-Bonn interaction instead of the Argonne V18 interaction.

Nucleus	$M'^{0\nu}$ 1bc	$M'^{0\nu}$ 2bc								$\langle M'^{0\nu}\rangle$ with quenching	ε (%)
		Parameters of Ref. [7]				Parameters of Ref. [8]					
		a	b	c	d	a	b	c	d		
^{48}Ca	0.649	0.615	0.605	0.561	0.542	0.606	0.606	0.570	0.569	0.58(0.03)	10
^{76}Ge	5.849	5.086	4.904	4.356	4.082	4.990	4.858	4.371	4.175	4.60(0.40)	21
^{82}Se	5.255	4.538	4.366	3.848	3.577	4.453	4.327	3.867	3.669	4.08(0.38)	22
^{96}Zr	3.144	2.953	2.872	2.608	2.485	2.883	2.835	2.603	2.532	2.72(0.18)	12
^{100}Mo	6.164	5.469	5.295	4.747	4.469	5.326	5.208	4.726	4.542	4.97(0.39)	19
^{110}Pd	6.532	5.772	5.589	5.029	4.758	5.629	5.497	4.998	4.806	5.26(0.40)	19
^{116}Cd	4.474	3.888	3.749	3.338	3.125	3.796	3.685	3.317	3.149	3.51(0.31)	22
^{124}Sn	4.024	3.646	3.556	3.273	3.158	3.553	3.494	3.239	3.170	3.29(0.20)	16
^{130}Te	4.642	4.063	3.921	3.473	3.242	3.958	3.861	3.468	3.313	3.66(0.32)	21
^{136}Xe	2.602	2.276	2.196	1.943	1.812	2.206	2.149	1.926	1.837	2.04(0.18)	21

and still contribute non-negligibly at several hundred MeV. In addition, the $0\nu\beta\beta$ matrix element contains a Fermi part, for which we have assumed no quenching. While this assumption may not be completely accurate, it is implied at low momentum transfer by conservation of the vector current (CVC). The overall quenching of the vector current is certain to be less than that of the axial-vector current. (In the results listed in Tables I and II the Fermi matrix elements are smaller than in some other calculations because the isovector particle-particle interaction was adjusted as explained in Ref. [15] to reflect isospin symmetry).

Why is the QRPA $0\nu\beta\beta$ quenching less than that in the shell model? Part of the reason, as we noted in the introduction, is that in the QRPA the strength of the isoscalar pairing interaction, which we call $g_{pp}^{T=0}$, is adjusted to reproduce the measured $2\nu\beta\beta$ rate. The suppression of $2\nu\beta\beta$ decay by two-body currents implies that the value of $g_{pp}^{T=0}$ is smaller than it would be without those currents. The smaller $g_{pp}^{T=0}$ in turn implies less quenching for the $0\nu\beta\beta$ matrix element.

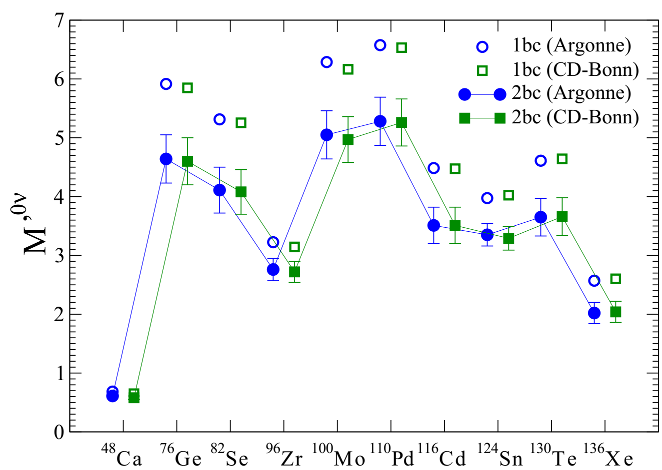

FIG. 1. (Color online) Nuclear matrix elements $M'^{0\nu}$ for all the nuclei considered here. The empty circles and squares represent the results with the one-body current only, and the solid circles and squares the average of the results with two-body currents included. The error bars represent the dispersion in those values (see text).

Figure 2 illustrates this idea. The upper panel shows the $2\nu\beta\beta$ matrix element, with (solid red) and without (dashed blue) two-body currents. The two vertical lines indicate the values of $g_{pp}^{T=0}$ needed to reproduce the "measured" matrix element [16], defined as that which gives the lifetime under the assumption that g_A is unquenched. The value of $g_{pp}^{T=0}$ that works with the two-body currents is smaller. The lower panel shows the consequences for $0\nu\beta\beta$ decay. The longer (purple) arrow represents the quenching that would obtain if $g_{pp}^{T=0}$ were not adjusted for the presence of the two-body currents (as is the case in the shell model, where the interaction is fixed ahead of time). The shorter arrow represents the same quenching after adjusting $g_{pp}^{T=0}$. The requirement that we reproduce $2\nu\beta\beta$ decay thus means that the $0\nu\beta\beta$ matrix element is quenched noticeably less than it would otherwise be.

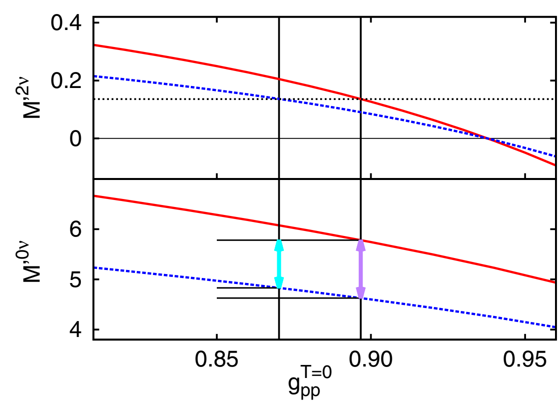

FIG. 2. (Color online) The quenching of $2\nu\beta\beta$ and $0\nu\beta\beta$ decay by two-body currents in χEFT. Top: $M'^{2\nu}$ vs $g_{pp}^{T=0}$, the strength of isoscalar pairing. The solid (red) line is the unquenched matrix element and the dashed (blue) line the matrix element with quenching caused by two-body currents, with the parametrization EGM+δc_i from Ref. [7]. The dotted black line is the measured matrix element [16] under the assumption that g_A is unquenched. The vertical lines are the values of $g_{pp}^{T=0}$ that reproduce the measurement with and without two-body currents. Bottom: The same, for $M'^{0\nu}$ (without a measured value). The long (purple) arrow represents the quenching when $g_{pp}^{T=0}$ is not readjusted to reproduce $2\nu\beta\beta$ decay. The short (cyan) arrow is the quenching when $g_{pp}^{T=0}$ is readjusted.

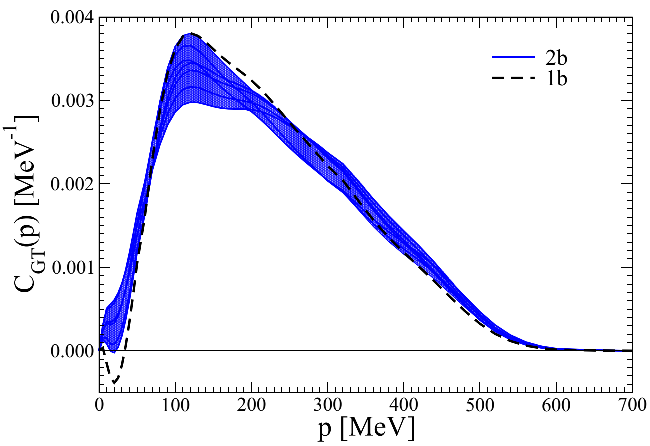

FIG. 3. (Color online) Normalized momentum transfer distribution $C_{GT}(p)$ for the Gamow-Teller part of $M'^{0\nu}$ in ^{136}Xe. The dashed line is the unquenched (one-body current) distribution and the filled area represents the range of the distributions produced by the variants of the calculation that include two-body currents. For these latter, the average momentum $\langle p \rangle$ is about 230 MeV and $\sqrt{\langle p^2 \rangle} \sim 255$ MeV, vs about 190 MeV and 225 MeV in Ref. [7].

Another difference between the QRPA and the shell model is that the QRPA works in a much larger single-particle space (at the price of working with only a particular kind of correlation). This larger space presumably means larger contributions at high momentum transfer. Since the quenching decreases with momentum transfer, the contributions of the high-angular-momentum multipoles are less affected by the two-body currents than their low-angular-momentum counterparts. The large QRPA model space therefore makes the quenching of $0\nu\beta\beta$ decay less than it would be in a shell model calculation. To demonstrate that fact, we show in Fig. 3 the distribution in momentum transfer (normalized to 1) of the Gamow-Teller part of the 0ν matrix element for ^{136}Xe; the inset in Fig. 2 of Ref. [7] shows the same distribution. The shapes of our curve and that in Ref. [7] are visibly different and, indeed, the averages $\langle p \rangle$ and $\sqrt{\langle p^2 \rangle}$ are 15 or 20 percent larger in QRPA than in the shell model, both for the unquenched and quenched variants.

IV. DISCUSSION

It is clear, in today's terminology, that some of the quenching of spin operators in nuclei is due to the use of restricted model spaces and some to many-body currents. Model-space truncation can exclude strength that may be pushed to high energies, and the omission of two-body currents leaves delta-hole excitations, among other things, unaccounted for. The question of which effect is more important is still open. If two-body currents are behind most of the quenching, as recent fits of the c parameters seem to suggest, then the spin operators in $2\nu\beta\beta$ decay (and ordinary beta decay as well) are very likely more quenched than those in $0\nu\beta\beta$ decay, and existing calculations of $0\nu\beta\beta$ decay that do not include quenching are at least roughly correct. We have seen that the quenching of $0\nu\beta\beta$ decay is mild in the QRPA, even a bit milder than in the shell model, and in sharp contrast to the severe quenching discussed, e.g., in Ref. [17].

It is of course possible that, as older meson-exchange models suggest [1], the effects of many-body currents are small at all momentum transfer. In that event the quenching of $0\nu\beta\beta$ decay would be unrelated to the two-body currents and could be similar in magnitude to the quenching of $2\nu\beta\beta$ decay, a state of affairs that would make $0\nu\beta\beta$ experiments less sensitive to a Majorana neutrino mass than we currently believe. A strong argument that this state of affairs is real, however, has yet to be presented. It seems likely to us that the quenching of $0\nu\beta\beta$ matrix elements is around the size indicated by the χEFT plus QRPA analysis carried out here or the χEFT plus shell-model analysis of Ref. [7].

ACKNOWLEDGMENTS

This work was supported by the US Department of Energy through Contract No. DE-FG02-97ER41019. F.Š. acknowledges the support by the VEGA Grant agency of the Slovak Republic under the Contract No. 1/0876/12 and by the Ministry of Education, Youth and Sports of the Czech Republic under Contract No. LM2011027.

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