Quantum Monte Carlo calculations of weak transitions in $A = 6–10$ nuclei

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$Ab initio$ calculations of the Gamow-Teller (GT) matrix elements in the $\beta$ decays of $^6$He and $^{10}$C and electron captures in $^7$Be are carried out using both variational and Green’s function Monte Carlo wave functions obtained from the Argonne $v_{18}$ two-nucleon and Illinois-7 three-nucleon interactions, and axial many-body currents derived from either meson-exchange phenomenology or chiral effective field theory. The agreement with experimental data is excellent for the electron captures in $^7$Be, while theory overestimates the $^6$He and $^{10}$C data by $\sim 2\%$ and $\sim 10\%$, respectively. We show that for these systems correlations in the nuclear wave functions are crucial to explain the data, while many-body currents increase by $\sim 2–3\%$ the one-body GT contributions.

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A major objective of nuclear theory is to explain the structure and dynamics of nuclei in a fully microscopic approach. In such an approach the nucleons interact with each other in terms of many-body (primarily, two- and three-body) effective interactions, and with external electroweak probes via effective currents describing the coupling of these probes to individual nucleons and many-body clusters of them. We will refer below to this approach as the basic model of nuclear theory.

For light nuclei (s- and p-shell nuclei up to $^{12}$C), quantum Monte Carlo (QMC) and, in particular, Green’s Function Monte Carlo (GFMC) methods allow us to carry out first-principles, accurate calculations of a variety of nuclear properties within the basic model. These calculations retain the full complexity of the many-body correlations induced by the Hamiltonians and currents, which have an intricate spin-isospin operator structure. When coupled to the numerically accurate QMC methods, the deceptively simple picture put forward in the basic model provides a quantitative and accurate description of the structure and dynamics of light nuclei over a broad energy range, from the keV’s relevant in nuclear astrophysical contexts to the MeV’s of low-lying nuclear spectra and radiative decay processes, to the GeV’s probing the short-range structure of nuclei and the limits of the basic model itself.

In the present study we focus on low-energy weak transitions in nuclei with mass number $A = 6–10$. In this mass range there are few microscopic calculations of GT matrix elements. Calculations based on the one-body GT operator have been performed, e.g., in Refs. (discussed below) and Refs. that report on the $^4$He $\beta$-decay. However, most of the calculations of $\beta$-decays and electron-capture processes in this mass range have mainly relied on relatively simple shell-model or cluster descriptions of the nuclear states involved in the transitions.

The shell model—its own approximation of the basic model—has typically failed to reproduce the measured Gamow-Teller (GT) matrix elements governing these weak transitions, unless use was made of an effective one-body GT operator, in which the nucleon axial coupling constant $g_A$ is quenched relative to its free value (ranging from $g_A^{\text{eff}} \simeq 0.85 g_A$ in the light nuclei under consideration here to $g_A^{\text{eff}} \simeq 0.7 g_A$ in heavy nuclei). More phenomenological models have been based on $\alpha$-nucleon-nucleon (for $A=6$) or $\alpha$-$\alpha$ and $\alpha$-$\alpha$-nucleon-nucleon (for $A=10$) clusterization, and have used Faddeev techniques with a separable representation of the nucleon-nucleon and $\alpha$-nucleon interactions or the resonating-group method or rather crude potential wells. While these studies provide useful insights into the structure of these light systems, nevertheless their connection to the basic model is rather tenuous. In particular, they do not explain whether the required quenching of $g_A$ in shell-model calculations reflects deficiencies in the corresponding wave functions—possibly due to the lack of correlations and/or to limitations in model space—or in the model adopted for the nuclear axial current, in which many-body terms are typically neglected.

The first QMC calculation of the $A = 6–7$ weak transitions in the basic model was carried out with the Variational Monte Carlo (VMC) method in Ref. (discussed below) and Refs. that report on the $^4$He $\beta$-decay. However, most of the calculations of $\beta$-decays and electron-capture processes in this mass range have mainly relied on relatively simple shell-model or cluster descriptions of the nuclear states involved in the transitions.
nuclear axial currents including, apart from the (one-body) GT operator, two-body operators, which arise naturally in a meson-exchange picture (π- and ρ-exchange, and ρπ-transition mechanisms) and when excitations of nucleon resonances (notably the Δ isobar) are taken into account. These two-body operators, multiplied by hadronic form factors so as to regularize their short-range behavior in configuration space, were then constrained to reproduce the GT matrix element contributing to tritium behavior in configuration space, were then constrained to reproduce and two-body (2b) operators are also listed (including fitted values of the N-to-Δ axial coupling constant). The χEFT axial current [27] consists of 1b, 2b, and three-body (3b) operators. The 1b operators read

\[ J^{1b}_{5,\pm} = -g_A \sum_{i=1}^{A} \tau_{i,\pm} \left( \sigma_i - \nabla_i \cdot \sigma_i - \sigma_i \cdot \nabla_i \right) \left( \sigma_i - \nabla_i \cdot \sigma_i - \sigma_i \cdot \nabla_i \right) / 2m^2, \]

where \( \tau_{i,\pm} = (\tau_{i,x} \pm i \tau_{i,y})/2 \) is the standard isospin raising (+) or lowering (−) operator, and \( \sigma_i \) and \( -i \nabla_i \) are, respectively, the Pauli spin matrix and momentum operator of nucleon \( i \). The 2b and 3b operators are illustrated diagrammatically in Fig. 1 in the limit of vanishing momentum transfer considered here. Referring to Fig. 1, the 2b operators are from contact [CT, panel (a)], one-pion exchange (OPE) [panels (b) and (f)], and multi-pion exchange (MPE) [panels (c)-(e) and (g)],

\[ J^{2b}_{5,\pm} = \sum_{i<j=1}^{A} \left( J^{CT}_{5,\pm}(ij) + J^{OPE}_{5,\pm}(ij) + J^{MPE}_{5,\pm}(ij) \right), \]

and the 3b operators are from MPE [panels (h)-(i)],

\[ J^{3b}_{5,\pm} = \sum_{i<j<k=1}^{A} J^{MPE}_{5,\pm}(ijk). \]

Configuration-space expressions for these 2b and 3b operators are reported in Ref. [34].

The 1b operator in Eq. (1) includes the leading order (LO) GT term and the first non-vanishing corrections to it, which come in at next-to-next-to-leading order (N2LO) [34]. The latter are nominally suppressed by two powers of the expansion parameter \( Q/\Lambda_\chi \), where \( Q \) specifies generically the low-momentum scale and \( \Lambda_\chi = 1 \) GeV is the chiral-symmetry-breaking scale, but being inversely proportional to the nucleon mass \( m \), itself of order \( \Lambda_\chi \), are in fact further suppressed than this naive power counting would indicate. Long-range 2b corrections from OPE enter at N3LO, panel (b) in Fig. 1 involving the low-energy constants (LECs) \( c_3 \) and \( c_4 \) in the sub-leading \( L_5^{(2)} \) chiral Lagrangian [35], as well as at N4LO, panel (f). In terms of the expansion parameter \( Q/\Lambda_\chi \), they scale as \( (Q/\Lambda_\chi)^3 \) and \( (Q/\Lambda_\chi)^4 \), respectively, relative to the LO. Loop corrections from MPE, panels (c)-(e) and (g), come in at N4LO, as do 3b currents, panels (h)-(i). Finally, the contact 2b current at N3LO, panel (a), is proportional to a LEC, denoted as \( z_0 \).
The short-range behavior of the 2b and 3b operators is regularized by including a cutoff \( C_A(k) = \exp(-k^2/\Lambda^2) \) in momentum space, and the values \( \Lambda = 500 \) and 600 MeV are considered in the present work. In correspondence to each \( \Lambda \) and to each set of \( \{c_3, c_4\} \), either \( \{c_3, c_4\} = (-3.2, 5.4) \) GeV\(^{-1}\) as reported in Ref. [36] or \( \{c_3, c_4\} = (-5.61, 4.26) \) GeV\(^{-1}\) as determined in Ref. [37], the LEC \( z_0 \) is constrained to reproduce the measured GT matrix element of tritium in hyperspherical-harmonics calculations based on the AV18+UIX Hamiltonian [38]. While the three-nucleon potential employed in that work, the Urbana IX (UIX) model [38], is different from the IL7 adopted here, nevertheless both Hamiltonians, AV18+UIX and AV18+IL7, reproduce the empirical values for the trinucleon binding energies and charge radii. In particular, we find that the GFMC tritium GT matrix element obtained with the 1b axial current and AV18+IL7 Hamiltonian is within \( \sim 3\% \) of the experimental determination.

Reduced matrix elements (RMEs) for the \( \beta \) decays between the \(^6\text{He}(0^+; 1)\) and \(^6\text{Li}(1^+; 0)\) ground states, and between the \(^{10}\text{C}(0^+; 1)\) ground state and \(^{10}\text{B}(1^+; 0)\) first excited state, and \( \epsilon \) captures of the \(^7\text{Be}(3/2^+; 1/2)\) ground state to the \(^7\text{Li}(3/2^-; 1/2)\) ground state and \(^7\text{Li}(1/2^-; 1/2)\) first excited state are listed in Table I (in parentheses are the spin-parity, \( J^\pi \), and isospin, \( T \), assignments for each state). All processes are allowed or superallowed, and are therefore driven (almost) exclusively by the axial current (and, additionally, the vector charge—the Fermi operator—for the transition between the ground states of \(^7\text{Be}\) and \(^7\text{Li}\)). Retardation effects from the momentum transfer dependence of the operators, and corrections from suppressed transitions, such as, for example, those induced in the \( A = 6 \) and 10 decays by the magnetic dipole associated with the vector current, are negligible [13]. Therefore the RMEs listed in Table I follow simply from

\[
RME = \sqrt{\frac{2J_f + 1}{g_A}} \frac{\langle J_f M | j^z_{5, \pm} | j_i M \rangle}{\langle J_i M; 0 | J_f M \rangle},
\]

where \( j^z_{5, \pm} \) is the \( z \)-component of the axial current \( j_5 \) (at vanishing momentum transfer) given above and \( \langle J_i M, 0 | J_f M \rangle \) are Clebsch-Gordan coefficients. The VMC results are obtained by straightforward Monte Carlo integration of the nuclear matrix elements above between (approximate) VMC wave functions; the GFMC results are from mixed-estimate evaluations of these matrix elements using previously generated GFMC configurations for the states under consideration, as illustrated in Ref. [14].

| \( ^6\text{He} \) \( \beta \)-decay | \( ^7\text{Be} \) \( \epsilon \)-capture (gs) | \( ^7\text{Be} \) \( \epsilon \)-capture (ex) | \( ^{10}\text{C} \) \( \beta \)-decay |
|---|---|---|---|
| LO | 2.168(2.174) | 2.294(2.334) | 2.083(2.150) |
| N4LO | 3.73(3.63) \( \times 10^{-2} \) | 6.07(4.98) \( \times 10^{-2} \) | 4.63(4.63) \( \times 10^{-2} \) |
| N4LO* | 3.62(3.43) \( \times 10^{-2} \) | 6.62(5.43) \( \times 10^{-2} \) | 5.31(5.38) \( \times 10^{-2} \) |
| MEC | 6.90(4.57) \( \times 10^{-2} \) | 10.5(10.3) \( \times 10^{-2} \) | 8.88(8.99) \( \times 10^{-2} \) |
| EXP | 2.1609(40) | 2.3556(47) | 2.1116(57) |

Table I: Gamow-Teller RMEs in \( A = 6, 7 \), and 10 nuclei obtained with chiral axial currents and GFMC (VMC) wave functions corresponding to the AV18+IL7 Hamiltonian model. Results corresponding to the one-body LO contribution (row labeled LO) and to the sum of all corrections beyond LO obtained with cutoff \( \Lambda = 500 \) MeV and 600 MeV (rows labeled respectively as N4LO and N4LO*), are listed. The sum of all two-body corrections obtained with conventional meson-exchange axial currents is listed in the row labeled MEC. Cumulative contributions, to be compared with the experimental data [16, 17, 39, 40] reported in the last row, are obtained by adding to the LO terms the contributions from either the chiral (N4LO or N4LO*) or the conventional (MEC) currents. Statistical errors associated with the Monte Carlo integrations are not shown, but are \( \sim 1\% \).

The sum of all contributions beyond LO, denoted as N4LO and N4LO* in Table I, leads approximately to a
2–3% increase in the LO prediction for the GT matrix elements of all processes under consideration. There is some cutoff dependence in these contributions, as indicated by the difference between the rows labeled N4LO and N4LO* in Table I which may be aggravated here by the lack of consistency between the \( \chi \) EFT currents and the phenomenological potentials used to generate the wave functions, i.e., by the mismatch in the short-range behavior of potentials and currents. The N4LO and N4LO* results in Table I correspond to the set \((c_3, c_4) = ( -3.2, 5.4)\) GeV\(^{-1}\) [39] in the OPE GT operator at N3LO. To illustrate the sensitivity of predictions to the set of \((c_3, c_4)\) values, we observe that use of the more recent determination \((c_3, c_4) = ( -5.61, 4.26)\) GeV\(^{-1}\) [37] would lead to an N4LO GFMC-calculated value of 6.71(2.89) \(\times 10^{-2}\) for the \(^7\)Be \(\epsilon\) capture to the \(^7\)Li ground (first excited) state for the choice of cutoff \(\Lambda = 500\) MeV, to be compared to the corresponding 6.07(4.63) \(\times 10^{-2}\) reported in Table I. Lastly, the N4LO contributions obtained with the more accurate GFMC wave functions are about 20% larger than those corresponding to VMC wave functions for the \(^6\)He and \(^7\)Be-to-\(^7\)Li ground-state transitions, although it should be emphasized that this is in relation to a small overall \(\sim 2\)% correction from 2b and 3b operators.

The contributions of the axial current order-by-order in the chiral expansion are given for the GT matrix element of the \(^7\)Be \(\epsilon\) capture in Table II. Those beyond LO, with the exception of the CT at N3LO, have opposite sign relative to the (dominant) LO. The loop corrections N4LO(2b) are more than a factor 5 larger (in magnitude) than the OPE. This is primarily due to the accidental cancellation between the terms proportional to \(c_3\) and \(c_4\) in the OPE operator at N3LO (which also occurs in the tritium GT matrix element [33]). It is also in line with the chiral filter hypothesis [41,45], according to which, if soft-pion processes are suppressed—as is the case for the axial current—then higher-order chiral corrections are not necessarily small. Indeed, the less than 3% overall correction due to terms beyond LO reported in Table II (row N4LO) comes about because of destructive interference between two relatively large (\(\sim 10\)% contributions from the CT and the remaining [primarily N4LO(2b)] terms considered here. Lastly, we note that going to N5LO (one order higher than in the present study) would introduce additional contact terms, thus reducing the predictive power of the theory, since the associated LECs would have to be determined by fitting, in addition to the \(^3\)H \(\beta\)-decay rate, some of the weak rates discussed here.

|          | \(g^s\)   | \(g^s\) |
|----------|-----------|---------|
| LO       | 2.334     | 2.150   |
| N2LO     | -3.18 \(\times 10^{-2}\) | 2.79 \(\times 10^{-2}\) |
| N3LO(CT) | 2.79 \(\times 10^{-4}\) | 2.36 \(\times 10^{-4}\) |
| OPE      | -2.99 \(\times 10^{-2}\) | -2.44 \(\times 10^{-2}\) |
| N4LO(2b) | -1.61 \(\times 10^{-3}\) | -1.34 \(\times 10^{-3}\) |
| N4LO(3b) | -6.59 \(\times 10^{-3}\) | -4.86 \(\times 10^{-3}\) |

TABLE II. Individual contributions to the \(^7\)Be \(\epsilon\)-capture Gamow-Teller RMEs obtained at various orders in the chiral expansion of the axial current (\(\Lambda = 500\) MeV) with VMC wave functions. The rows labeled LO and N2LO refer to, respectively, the first term and the terms proportional to \(1/m^2\) in Eq. (1); the rows labeled N3LO(CT) and OPE, and N4LO(2b) and N4LO(3b), refer to panel (a) and panels (b) and (f), and to panels (c)-(e), (g) and panel (h) in Fig. I, respectively.

Theoretical predictions corresponding to the \(\chi\) EFT axial current at LO (blue circles) and up to N4LO (magenta stars). Green squares indicate ‘unquenched’ shell model calculations from Ref. [17] based on the LO axial current.

The ratios of GFMC to experimental values for the GT RMEs in the \(^3\)H, \(^6\)He, \(^7\)Be, and \(^10\)C weak transitions. Theory predictions correspond to the \(\chi\) EFT axial current at LO and including corrections up to N4LO. The experimental values are those listed in Table II while that for \(^3\)H is 1.6474(24) [34]. These values have been obtained by using \(g_A = 1.2723(23)\) [40] and \(K/ [G_A^2 (1 + \Delta^g_R)] = 6144.5(1.4)\) sec [47], where \(K = 2\pi^3 \ln 2/m^*_e = 8120.2776(9) \times 10^{-10}\) GeV\(^{-4}\) sec and \(\Delta^g_R = 2.361(38)%\) is the transition-independent radiative correction [47]. In the case of the \(\beta\) decays, but not for the \(\epsilon\) captures, the transition-dependent \((\delta^R_R)\) radiative correction has also been accounted for. Lastly, in the \(\epsilon\) processes the rates have been obtained by ignoring the factors \(B_K\) and \(B_{L1}\) which include the effects of electron exchange and overlap in the capture from the \(K\) and \(L1\) atomic subshells. As noted by Chou et al. [17], following Bahcall [18,49], such an approximation is expected to be valid in light nuclei, since these factors only account for a redistribution of the total strength among the different subshells (however, it should be noted that \(B_K\) and \(B_{L1}\) were retained in Ref. [13], and led to the extraction of...
We find overall good agreement with data for the $^6$He $\beta$-decay and $\epsilon$ captures in $^7$Be, although the former is overpredicted by $\sim 2\%$, a contribution that comes almost entirely from 2b and 3b chiral currents. The experimental GT RMEs for the $^{10}$C $\beta$-decay is overpredicted by $\sim 10\%$, with two-body currents giving a contribution that is comparable to the statistical GFMC error. The presence of a second (1$^+; 0$) excited state at $\sim 2.15$ MeV can potentially contaminate the wave function of the $^{10}$B excited state at $\sim 0.72$ MeV, making this the hardest transition to calculate reliably. In fact, a small admixture of the second excited state ($\sim 6\%$ in probability) in the VMC wave function brings the VMC reduced matrix element in statistical agreement with the the measured value, a variation that does not spoil the overall good agreement we find for the reported branching ratios of 98.54(14)$\%$ ($< 0.08\%$) to the first (second) (1$^+; 0$) state of $^{10}$B$^\ast$. Because of the small energy difference of these two levels, it would require a computationally expensive GFMC calculation to see if this improvement remains or is removed; in lighter systems we have found that such changes of the trial VMC wave function are removed by GFMC.

We note that correlations in the wave functions significantly reduce the matrix elements, a fact that can be appreciated by comparing the LO GFMC (blue circles in Fig. 2) and the LO shell model calculations (green squares in the same figure) from Ref. [17]. Moreover, preliminary variational Monte Carlo studies, based on the Norfolk two- and three-nucleon chiral potentials [6, 50, 51] and the LO GT operator, bring the $^{10}$C prediction only $\sim 4\%$ above the experimental datum [52], indicating that the $\sim 10\%$ discrepancy we find here may indeed be attributable to deficiencies in the AV18+IL7 wave functions of $A = 10$ nuclei.

In the present study we have shown that weak transitions in $A = 6$–10 nuclei can be satisfactorily explained in the basic model, without having to “quench” $g_A$. Clearly, in order to resolve the mismatch in the short-range behavior between potentials and currents alluded to earlier, GFMC calculations based on the Norfolk chiral potentials of Refs. [6, 51] and consistent chiral currents are in order. Work along these lines is in progress.

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