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Induced-Tensor and Second-Order-Forbidden Terms*

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A model-independent result for the β-γ correlation in allowed β decay is given, including all second-order-forbidden corrections. A suggested method for measurement of the first-class contribution to the induced-tensor form factor is discussed.

I. INTRODUCTION

Recently, we have reported predictions for contributions to the induced-tensor term d_{I} in allowed nuclear β decay which arise from the conventional (first-class) axial-vector current.[1] Of course, if second-class axial currents[2] contribute to the semileptonic weak decays, as suggested by one interpretation of experiments on mirror nuclei,[3] they would make a substantial contribution to the induced tensor. However, in Ref. 1 we showed how, by using mirror transitions, the second-class component d_{II} can be separated out so that a measurement of the first-class contribution alone results. We found, correct to first order in recoil and assuming the nuclear impulse approximation,

$$d_I = A\langle\beta\| i\sum_i \sigma_i \times L_i \|\alpha\rangle ,$$

where A is the mass number, so that experimental measurement of d_I would provide an interesting test of the impulse approximation in recoil order.

However, Ref. 1 was incomplete in that second-order-forbidden terms of the order q/m, where q is the momentum transfer and m is the nucleon mass, were included but second-order-forbidden terms of the order q^2r^2, where r is the nuclear radius, were neglected. This may not be a good approximation, since all but one of the suggested decays is K-forbidden, thus tending to enhance matrix elements of operators with $\Delta K > 1$ relative to those of operators with $\Delta K \leq 1$. Here we report the calculation of the suggested β-γ and β-α correlations including all second-order-forbidden corrections in a model-independent form. Electromagnetic effects are again ignored except for final-state Coulomb interactions included in the standard Fermi function. In Sec. II we define the relevant form factors. Section III gives our results, and Sec. IV discusses the significance of these results for the transitions suggested in Ref. 1.

II. DEFINITIONS

We shall assume the validity of the usual current-current interaction. Then the β-decay amplitude is given (for electron decay – modifications suitable for positron decay will be included in the final formulas) by

$$T = \frac{G_V}{\sqrt{2}} \cos\theta_C \langle\beta_{p_2}| V_\mu + A_\mu |\alpha_{p_1}\rangle l^\mu , \qquad (1)$$

where G_V is the weak-coupling constant ($G_V m_p^{\,2} \simeq 10^{-5}$), θ_C is the Cabibbo angle, and l^μ is the matrix element of the lepton current,

$$l^\mu = \bar{u}(p)\gamma^\mu(1+\gamma_5)v(k) .$$

Let p_1, p_2, p, and k denote the four-momenta of parent nucleus, daughter nucleus, electron, and neutrino; and M_1 and M_2 represent the parent and daughter masses. We also define

$$P = p_1 + p_2 , \qquad q = p_1 - p_2 = p + k ,$$
$$M = \tfrac{1}{2}(M_1 + M_2) , \qquad \Delta = M_1 - M_2 .$$

We write for the amplitude of a general allowed transition including all form factors up to second order in the momentum transfer[4]

$$\begin{aligned}\langle\beta_{p_2}| V_\mu |\alpha_{p_1}\rangle l^\mu &= \left(\frac{a}{2M} P\cdot l + \frac{e}{2M} q\cdot l\right)\delta_{JJ'}\delta_{MM'} - i\frac{b}{2M} C^{M'k;M}_{J'1;J}(l\times q)_k \\ &\quad + C^{M'k;M}_{J'2;J}\left(\frac{f}{2M}\sum_{n,n'} C^{nn';k}_{11;2} l_n q_{n'} + \frac{g}{4M^2}\frac{P\cdot l}{2M}\sqrt{\tfrac{4}{5}\pi}\,Y^k_2(q)\right), \\ \langle\beta_{p_2}| A_\mu |\alpha_{p_1}\rangle l^\mu &= C^{M'k;M}_{J'1;J}\frac{\epsilon^{ijk}}{4M}\left(\epsilon_{ij\lambda\eta} l^\lambda (cP^\eta \mp dq^\eta) + \frac{h}{4M^2}\epsilon_{ij\lambda\eta} q^\lambda P^\eta q\cdot l\right) \\ &\quad + \frac{j_2}{4M^2} C^{M'k;M}_{J'2;J}\sum_{n,n'} C^{nn';k}_{12;2} l_n \sqrt{\tfrac{4}{5}\pi}\,Y^{n'}_2(q) + \frac{j_3}{4M^2} C^{M'k;M}_{J'3;J}\sum_{n,n'} C^{nn';k}_{12;3} l_n \sqrt{\tfrac{4}{5}\pi}\,Y^{n'}_2(q) ,\end{aligned} \qquad (2)$$

where J, J' are the spins of the parent and daughter nucleus, respectively, and M, M' represent the initial and final components of nuclear spin along some axis of quantization. Here repeated Latin indices are summed from 1 to 3, while repeated Greek symbols imply a four-vector contraction with the metric $g_{00} = -g_{ii} = 1$. The form factors $a, b, \ldots, j_3$ are, in general, functions of the momentum transfer q^2.

Our results, given in the next section, will be in terms of these 10 form factors and will be independent of specific nuclear models. However, in order to express our form factors in more conventional terms, we give here the impulse-approximation predictions for the nuclear form factors. We find for the vector current

$$\begin{aligned} a &= g_V M_F + \tfrac{1}{6} q^2 g_V M_{r^2} , \\ b &= g_V M_L + (g_V + g_M) M_{GT} , \\ e &= -\tfrac{1}{3} M\Delta g_V M_{r^2} , \\ f &= \sqrt{\tfrac{2}{3}} M\Delta g_V M_Q , \\ g &= -\tfrac{4}{3} M^2 g_V M_Q , \end{aligned} \qquad (3)$$

where $g_V \simeq 1$ and $g_M \simeq 4.7$ are the vector and weak-magnetism form factors for neutron decay, and

$$M_{\rm F} = \langle \beta \| \sum_i \tau_i^+ \| \alpha \rangle , \quad M_L = \langle \beta \| \sum_i \tau_i^+ L_i \| \alpha \rangle ,$$

$$M_{\rm GT} = \langle \beta \| \sum_i \tau_i^+ \sigma_i \| \alpha \rangle , \quad M_{r2} = \langle \beta \| \sum_i \tau_i^+ r_i^2 \| \alpha \rangle ,$$

$$M_Q = \sqrt{\tfrac{4}{5}\pi} \langle \beta \| \sum_i \tau_i^+ r_i^2 Y_2(\hat{r}_i) \| \alpha \rangle .$$

Thus, a is just the usual Fermi form factor, while b represents the weak-magnetism contribution. Form factors e and f make contributions to various spectral functions which are formally of order E/M, where E is the electron energy and M is the nuclear mass. However, they are both forbidden (to lowest order in recoil) on the conserved-vector-current hypothesis (CVC).[5] g is an isovector quadrupole term related, on the CVC hypothesis, to the difference of quadrupole moments of parent and daughter states if α, β are isotopic analogs.

For the axial-current terms, we have

$$c = g_A M_{\rm GT} + A \frac{\Delta}{2M} g_A M_{\sigma L} + \frac{1}{6\sqrt{10}} (2\Delta^2 + q^2) g_A M_{1y} ,$$

$$d = A g_A M_{\sigma L} + \frac{1}{\sqrt{10}} M \Delta g_A M_{1y} + M \Delta g_A M_{\sigma r} + A g_{\rm II} M_{\rm GT} ,$$

$$h = -\frac{2}{\sqrt{10}} M g_A M_{1y} + \frac{4M^2 g_A M_{\rm GT}}{m\pi^2 - q^2} , \qquad (4)$$

$$j_2 = -\tfrac{2}{3} M^2 g_A M_{2y} ,$$

$$j_3 = -\tfrac{2}{3} M^2 g_A M_{3y} ,$$

where $g_A \simeq 1.23$ and $g_{\rm II} = ?$ are the axial and induced-tensor form factors for neutron β decay, and

$$M_{GT} = \langle \beta \| \sum_i \tau_i^+ \sigma_i \| \alpha \rangle ,$$

$$M_{\sigma L} = \langle \beta \| i \sum_i \tau_i^+ \sigma_i \times L_i \| \alpha \rangle ,$$

$$M_{\sigma r} = \langle \beta \| \sum_i \tau_i^+ \sigma_i \cdot r_i r_i \| \alpha \rangle ,$$

$$M_{Ky} = \left\langle \beta \left\| \sum_i \tau_i^+ \frac{8\pi}{\sqrt{15}} C_{12;K} Y_1(\sigma_i) Y_2(r_i) \right\| \alpha \right\rangle .$$

Thus c is the usual Gamow-Teller form factor, while d represents the induced-tensor contribution. We note that the first-class contributions to d vanish between states which are exact isotopic analogs, as required by symmetry considerations.[6] h is the induced pseudoscalar and includes a pion-pole term.[7] j_2 and j_3 are additional axial form factors, which may be inportant to our results, as they have $\Delta K = 2, 3$, respectively.

III. RESULTS

Suppose a parent nucleus of spin J undergoes an allowed β decay to a daughter nucleus of spin J', which subsequently decays, emitting either a photon or an α particle to a final nucleus of spin J''. We assume the parent nucleus to be unpolarized, and we integrate over the neutrino direction, considering the spectrum in its dependence on the electron variables and on the photon (or α) direction. The latter is characterized by a unit vector $\hat{K}$ along the photon (or α) momentum as measured in the lab frame (rest frame of the parent). The spectrum is then found to be[8]

$$d\omega = F_{\mp}(Z,E) \frac{G_V^2 \cos^2\theta_C}{(2\pi)^5} (E_0 - E)^2 pE dE d\Omega_e d\Omega_K \left\{ f(E) + g(E) \frac{\hat{K} \cdot \vec{p}}{E} + h(E) \left[\left(\frac{\hat{K} \cdot \vec{p}}{E} \right)^2 - \frac{1}{3} \frac{p^2}{E^2} \right] \right\}, \qquad (5)$$

where $E(p)$ is the electron energy (momentum),

$$E_0 = \Delta \frac{1 + m_e^2 / 2M\Delta}{1 + \Delta/2M}$$

is the maximum electron energy, and $F_{\mp}(Z,E)$ is the usual Fermi function and accounts for the dominant Coulomb effects. The spectral functions f, g, h are found to be

$$f(E) = |a|^2 + |c|^2 - \frac{2}{3} \frac{E_0}{M} [|c|^2 \pm {\rm Re}\, c^*(b+d)] + \frac{2}{3} \frac{E}{M} (3|a|^2 + 5|c|^2 \pm 2{\rm Re}\, c^* b)$$
$$- \frac{1}{3} \frac{m_e^2}{ME} \left[2|c|^2 \pm {\rm Re}\, c^*(2b+d) - 3{\rm Re}\, a^* e - {\rm Re}\, c^* h \frac{E_0 - E}{2M} \right],$$

$$g(E) = \frac{2E_0}{3Mv^*} \{ -|a|^2 + \tfrac{1}{3} |c|^2 [1 - \tfrac{1}{10} \eta_{JJ'} \tau_{J',J''}(L)] \} - \frac{4E}{3Mv^*} \{ |a|^2 + \tfrac{5}{3} |c|^2 [1 - \tfrac{1}{100} \eta_{J,J'} \tau_{J',J''}(L)] \}, \qquad (6)$$

$$h(E) = \tau_{J',J''}(L) \frac{E}{M} \left\{ \eta_{J,J'} \tfrac{1}{20} [|c|^2 \pm {\rm Re}\, c^*(b-d)] - \Omega_{J,J'} \tfrac{1}{10} \left(2{\rm Re}\, a^* g \frac{E_0 + \frac{1}{2}E}{4M} + {\rm Re}\, a^* f \sqrt{\tfrac{3}{2}} \pm {\rm Re}\, a^* j_2 \sqrt{\tfrac{3}{2}} \frac{E_0 - E}{2M} \right) \right.$$
$$\left. - \Lambda_{J,J'} \tfrac{1}{10} \left(\pm {\rm Re}\, c^* g \sqrt{\tfrac{3}{2}} \frac{E_0 - E}{2M} \pm \tfrac{3}{2} {\rm Re}\, c^* f + 3{\rm Re}\, c^* j_2 \frac{E_0 - 2E}{4M} \right) - \Gamma_{J,J'} \tfrac{1}{10} 3{\rm Re}\, c^* j_3 \frac{E_0 + \frac{3}{2}E}{2M} \right\},$$

where the upper (lower) signs refer to electron (positron) decay and v^* is the photon (α) velocity in the center-of-mass frame of the daughter nucleus. The coefficients $\eta_{J,J'}$, $\Omega_{J,J'}, \ldots, \tau_{J',J''}(L)$ are given in Appendix A.

IV. DISCUSSION

In Ref. 1 we suggested a means to measure the first-class component of the induced tensor via measurement of β-γ (or β-α) correlation coefficients for certain mirror transitions. For example, if one measures the β-α correlation for the decay of Li^8 to the $J^P = 2^+$, 2.90-MeV excited state of Be^8 and that for the decay of B^8 to the same level, then (assuming no second-class vector currents)

$$\alpha_+ \equiv \left(\frac{h(E)}{f(E)}\right)^{B^8} + \left(\frac{h(E)}{f(E)}\right)^{Li^8}$$

$$= \frac{E}{M}\frac{1}{|c|^2}\left(|c|^2 - \mathrm{Re}\,c^* d_{\mathrm{I}} - \mathrm{Re}\,c^* j_2 \sqrt{\tfrac{1}{14}}\,\frac{3(E_0-2E)}{2M} - \frac{6}{7(35)^{1/2}}\mathrm{Re}\,c^* j_3 \frac{2E_0+5E}{2M}\right).$$

If the terms of second order in momentum transfer are omitted, as in Ref. 1, an experimental measure of $M_{\sigma L}$ results (assuming $|c|$ to be known from ft values). However, when $O(q^2r^2)$ terms are included there are two new features. First, we pick up additional contributions proportional to $\mathrm{Re}\,c^* j_2$, $\mathrm{Re}\,c^* j_3$ and, secondly, d_{I} is no longer just the interesting term $Ag_A M_{\sigma L}$ but receives first-class contributions also from M_{1y} and $M_{\sigma r}$.

In order to see whether experimental measurement of $M_{\sigma L}$ is still feasible, we have, for mass 8, calculated the relevant reduced matrix elements in the Nilsson model assuming the daughter state to belong to a pure rotational band with $K=0$ and the parent nucleus (Li^8 or B^8) to be in a linear combination of $K=1$ and $K=2$ states, with amplitudes γ and δ, respectively. Thus, for Li^8, we assume the ground state has an extra neutron in the Nilsson orbit having the form $|Njl\Omega\rangle = |1\,\tfrac{3}{2}\,1\,\tfrac{3}{2}\rangle$ in the limit of zero deformation and a vacancy to exist in the $|1\,\tfrac{3}{2}\,1\,\tfrac{1}{2}\rangle$ (for $K=1$) or $|1\,\tfrac{3}{2}\,1\,-\tfrac{1}{2}\rangle$ (for $K=2$) proton orbits. Harmonic-oscillator wave functions are used with radius parameter $\alpha = 0.75$.

For a deformation parameter $\beta = 0.3$, as suggested by nearby quadrupole moments, we find

$$\frac{j_2}{A^2c} \simeq 7\times 10^1\frac{\delta}{\gamma}, \quad \frac{j_3}{A^2c} \simeq 1.7\times 10^2\frac{\delta}{\gamma}, \quad \frac{d_{\mathrm{I}}}{Ac} \simeq \left(2.5 + 7.5\frac{\Delta}{m}\right) \simeq \frac{M_{\sigma L}}{M_{\mathrm{GT}}}, \tag{7}$$

so that

$$\alpha_+ \cong A\frac{E}{M}\left(\frac{1}{A} - \frac{M_{\sigma L}}{M_{\mathrm{GT}}} - \frac{\delta}{\gamma}40\frac{E_0-2E}{m} - \frac{\delta}{\gamma}25\frac{E_0+\frac{5}{2}E}{m}\right). \tag{8}$$

The experiments of Nordberg, Barnes, and Morinigo were performed with $E = 11$ MeV, whence the terms in the unknown mixing ratio δ/γ tend to cancel somewhat, but still yield a result which is of the same order as $M_{\sigma L}/M_{\mathrm{GT}}$. For mass 8 then, α_+ can yield only a rough estimate of d_{I}, and in general the contribution of terms in j_2 and/or j_3 can obscure completely the measurement of $M_{\sigma L}$. The cases of masses 24 and 28 which were suggested in Ref. 1 appear especially dubious when $O(q^2r^2)$ terms are included, since in these cases, involving $\Delta K = 4, 3$, respectively, j_3 is expected to play a dominant role. We conclude that although the nuclear impulse approximation produces a first-class contribution to the induced tensor, it is a difficult proposition to verify this prediction experimentally unless the absence of second-class currents is assumed, in which case a measurement of d itself – without separation of first-class and second-class components – is sufficient.[9]

APPENDIX A

We here quote values for the parameters used in Eq. (6). We find, for β-α correlations

$$\tau_{J',J''}(L) = \left(\frac{L(L+1)(2L+1)}{(2L-1)(2L+3)}\right)^{1/2}\left(\frac{(2J'-1)(2J'+1)(2J'+3)}{J'(J'+1)}\right)^{1/2} W(2J'LJ''; J'L). \tag{A1}$$

Explicit forms of $\tau_{J',J''}(L)$ for $L = 1, 2$ are given in Ref. 1.

Also, we have

$$\eta_{J,J'} = \begin{cases} -J'/(2J'+3), & J = J'+1 \\ 1, & J = J' \\ -(J'+1)/(2J'-1), & J = J'-1; \end{cases} \tag{A2}$$

$$\Omega_{J,J'} = [J'(J'+1)/(2J'-1)(2J'+3)]^{1/2}\delta_{J,J'};$$

$$\Lambda_{J,J'}=\frac{1}{\sqrt{2}\,(2J'+3)(2J'-1)}\begin{cases}-(2J'-1)[J'(J'+2)]^{1/2}, & J=J'+1\\ [3(2J'+3)(2J'-1)]^{1/2}, & J=J'\\ (2J'+3)[(J'-1)(J'+1)]^{1/2}, & J=J'-1;\end{cases}$$

$$\Gamma_{J,J'}=\frac{\sqrt{10}}{35}\left[\frac{(J'+2)(J'-1)}{(2J'-1)(2J'+3)}\right]^{1/2}\begin{cases}[J'(2J'+5)/(J'-1)(2J'+3)]^{1/2}, & J=J'+1\\ (3/2)^{1/2}, & J=J'\\ [(J'+1)(2J'-3)/(J'+2)(2J'-1)]^{1/2}, & J=J'-1.\end{cases} \tag{A2}$$

For β-γ correlations $\tau_{J',J''}(L)$ is multiplied by a factor $1-3/L(L+1)$ for $E(L)$ or $M(L)$ multipole radiation.

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[1]B. Holstein, Phys. Rev. C 4, 740 (1971).

[2]S. Weinberg, Phys. Rev. 132, 738 (1963).

[3]D. H. Wilkinson, Phys. Letters 31B, 447 (1970).

[4]In order to conform to the notation introduced in Ref. 1, d is defined with a change in sign for the case of positron decay.

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[6]Terms proportional to Δ vanish, since Δ is electromagnetic in origin, while $M_{\sigma L}$ is found to vanish for more subtle reasons. We have, if $U=e^{i\pi I_2}$ is the charge symmetry operation $U|I,I_3\rangle=(-)^{I+I_3}|I,I_3\rangle$. Thus if $|\beta\rangle=|I,I_3+1\rangle$ and $|\alpha\rangle=|I,I_3\rangle$,

$$\begin{aligned}\langle\beta\|i\sum_i\tau_i^{\dagger}\sigma_i\times L_i\|\alpha\rangle&=\langle\beta\|U^{-1}Ui\sum_i\tau_i^{\dagger}\sigma_i\times L_i\,U^{-1}U\|\alpha\rangle\\&=-\langle I,-I_3\|i\sum_i\tau_i^{\dagger}\sigma_i\times L_i\|I,-I_3-1\rangle\,.\end{aligned}$$

But by the Wigner-Eckart theorem

$$\frac{\langle\beta\|i\sum_i\tau_i^{\dagger}\sigma_i\times L_i\|\alpha\rangle}{\langle I,-I_3\|i\sum_i\tau_i^{\dagger}\sigma_i\times L_i\|I,-I_3-1\rangle}=\frac{C_{I\ 1\ I}^{I_3\ 1\ I_3+1}}{C_{I\ \ 1\ \ I}^{-I_3-1\ 1\ -I_3}}=1\,.$$

Hence the matrix element must vanish.

[7]We have used the partially conserved axial-vector-current value for the pion-pole term. See H. Primakoff in *High Energy Physics and Nuclear Structure*, edited by G. Alexander (North-Holland, Amsterdam, 1967).

[8]We have not included a small term $\propto \mathrm{Re}\,c^{*}j_3P_4(\hat{K}\cdot\vec{p})$.

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