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Asymmetry of Beta-Ray Angular Distribution in Polarized Nuclei and *G*-Parity Nonconservation

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We have derived an equation for the β-ray angular distribution including Coulomb corrections, radiative corrections, induced effect, and higher-order nuclear matrices. With this equation and the experimental data on β-ray asymmetries in polarized ^{12}B and ^{12}N, we conclude that the strength of the second-class induced tensor is $f_T/f_A = -(0.96 \pm 0.35)\times 10^{-3}$ in the limit of the impulse approximation. A possible modification of this value due to mesonic corrections is discussed.

Since Weinberg proposed a measurement of the ft-value ratio in mirror β decays to test a possible existence of the second-class currents in weak interactions,[1] there have been a number of articles published on this subject. Among those, Wilkinson and his co-workers have made an extensive search for the asymmetries of the ft values experimentally.[2] The results were originally thought to be a direct indication for the induced tensor interaction. Later on, these were, however, recognized as the sum of the effects due to nuclear structure, the induced tensor term f_T [see Eq. (4)], and possible meson-exchange currents. In a model calculation,[3] Kubodera, Delorme, and Rho adopted the meson-exchange effect due to $\omega\rightarrow\pi e\nu$, and they gave a ratio

$$(ft)_+/(ft)_- - 1 = \delta_{\mathrm{exp}} = \delta_{\mathrm{scc}} + \delta_{\mathrm{nucl}}.$$

Here

$$\delta_{\mathrm{scc}} = -4(\lambda/f_A)J + (2/3f_A)(\lambda L - 2\zeta)(E_0^- + E_0^+)$$

and δ_{nucl} represents the nuclear-structure effect. The effect of the second-class current, δ_{scc}, is also dependent on the nuclear model through J and L, while λ is a combination of the strong-coupling constants including the ω-ρ mixing parameter, and ζ is nearly equal to f_T. Wilkinson made an analysis of δ_{exp} in the region of mass number $A = 8$–30 systematically, using available nuclear models. He obtained as limits for the parameters[2]

$$|\zeta| \leq 2.5\times 10^{-3}\ \mathrm{MeV}^{-1} \text{ and } |\lambda| \leq 10\times 10^{-3}.$$

Here the combination $\zeta = \lambda = 0$ is not necessarily excluded, while the theory does not give us a single value of λ, since it contains a logarithmic divergence.

Less ambiguous information for the second-class current can be obtained from the measurement of the β-ray asymmetries in polarized nuclei. A long-awaited experiment on ^{12}B and ^{12}N was recently performed successfully by Sugimoto, Tanihata, and Göring.[4] In order to derive a conclusion about G-parity nonconservation from this experiment, we have to be careful to include all induced effects, higher-order corrections, etc., in the equation for the angular distribution of β rays.[5-8] In particular, the radiative corrections are important, at least for the β-ray spectrum,[9] and the Coulomb corrections are nonnegligible for the electron waves, especially with $j \geq \frac{3}{2}$.[5] Since we found several features insufficient for our purpose, first we state the method for determining the equation correctly, and next, we find the strength of G-parity nonconservation. Finally, we make a comment on possible modification due to meson-exchange effects, since our equation is based on the impulse approximation.

The higher-order corrections including Coulomb corrections are given in extended form by one of the present authors whose equations[5] include no induced tensor coupling. The effect of this G-parity-nonconserving second-class current is, however, easily taken into account by replacing in the published equations[5] the factors [in terms of Eq. (4) below]

$$C_A \int \vec{\sigma} \text{ by } (f_A - E_0 f_T) \int \vec{\sigma} \text{ and}$$
$$C_A(i \int \gamma_5 \vec{r}) \text{ by } (f_A + 2M f_T)(i \int \gamma_5 \vec{r}), \qquad (1)$$

where $C_V = f_V(0)$ and $C_A = f_A(0)$. We also adopt the nonrelativistic approximation for nuclear matrix elements:

$$\int \vec{\alpha} \times \vec{r} = x[(1 + \mu_p - \mu_n)/M] \int \vec{\sigma},$$
$$i \int \gamma_5 \vec{r} = -y(1/2M) \int \vec{\sigma}. \qquad (2)$$

Here $1 + \mu_p - \mu_n = 4.7$, and M is the nucleon mass. The parameters x and y are nuclear-model dependent, and they are given by

$$x = 1 + (1/4.7)(\int \vec{r} \times \vec{p})/(\int \vec{\sigma}),$$
$$y = 1 + i2c[\int \vec{r}(\vec{\sigma} \cdot \vec{p})]/(\int \vec{\sigma}). \qquad (3)$$

Here c is 1 for f_A and 0 for f_T. The reason for no nucleon-recoil term for f_T is seen in Eq. (7) below. In the case of ^{12}B and ^{12}N, we have $x \approx 1$,[10] and $y \approx 1.5$ for $c = 1$.[11]

We adopt the interaction Hamiltonian density

$$H = \{\bar{\psi}_p[\gamma_\lambda(f_V - f_A\gamma_5) + \sigma_{\lambda\nu}k_\nu(f_W + f_T\gamma_5) + ik_\lambda(f_S + f_P\gamma_5)]\psi_n\}[\bar{\psi}_e\gamma_\lambda(1+\gamma_5)\psi_\nu]/\sqrt{2} + \text{H.c.}, \qquad (4)$$

with $k = k_p - k_n$ and $\sigma_{\lambda\nu} = [\gamma_\lambda, \gamma_\nu]/2i$. The coupling constants for β decay are

$$f_V(0) = (3.001 \pm 0.002) \times 10^{-12} \hbar^3/m^2 c,$$
$$f_A(0) = -(1.239 \pm 0.011) f_V(0), \qquad (5a)$$
$$f_W = -(\mu_p - \mu_n) f_V/2M = -3.7 f_V/2M, \quad f_S = 0,$$

from conservation of vector current,

$$f_P \approx f_A/25, \qquad (5b)$$

from partial conservation of axial-vector current. The induced tensor term, f_T, introduced by Weinberg[1] is now being revealed. Equation (1) can be derived by comparing the space and time components of f_A and f_T in Eq. (4),

$$(f_A - E_0 f_T)(\psi_p^\dagger \vec{\sigma} \psi_n) \cdot [\psi_e^\dagger \vec{\sigma}(1+\gamma_5)\psi_\nu] \text{ and } -[f_A(\psi_p^\dagger \gamma_5 \psi_n) + if_T(\psi_p^\dagger \vec{\gamma}\gamma_5 \psi_n) \cdot \vec{p}][\psi_e^\dagger(1+\gamma_5)\psi_\nu], \qquad (6)$$

where $\vec{p}$ is the differential operator, and adopting the nonrelativistic approximation with two-component spinor u,

$$(\psi_p^\dagger \gamma_5 \psi_n) = -(\tfrac{1}{2}M)[(u_p^\dagger \vec{\sigma} u_n) \cdot \vec{p} + 2(u_p^\dagger \vec{\sigma} \cdot \vec{p} u_n)], \qquad (7)$$

$$i(\psi_p^\dagger \vec{\gamma}\gamma_5 \psi_n) \cdot \vec{p} = -(u_p^\dagger \vec{\sigma} u_n) \cdot \vec{p} \approx 2M(\psi_p^\dagger \gamma_5 \psi_n). \qquad (8)$$

Equation (8) is useful for deriving rule (1), by neglecting the nucleon recoil. In practice, a correction in Eq. (3) should be considered.

The explicit form of angular distribution of β rays in polarized nuclei is very lengthy; one must take into account Coulomb corrections, higher-order nuclear matrices of momentum type as well as coordinate type, radiative corrections,[12] and f_T terms simultaneously. These can be obtained from Ref. 5 by the replacements (1) and (2). To save space, we only give an expression for the $1^+ \to 0^+$ transition by omitting the contributions of $\int \vec{\sigma} r^2$ and $\int (\vec{\sigma} \cdot \vec{r})\vec{r}$ (the finite de Broglie wavelength effect is, however, included in the main L_0 and Λ_1 terms[5]):

$$W(\theta)=|f_A|^2|\int\vec{\sigma}|^2\Big(\{[1\mp 2E_0(f_T/f_A)][1+F(E,E_0)][\tfrac{1}{2}(1+\gamma)\mp\tfrac{10}{3}VR^2E\mp(2/3E)VR^2-(p^2/3)R^2]$$
$$\mp\tfrac{4}{3}a[E_0-2E+(1/E)\mp(\tfrac{2}{3})b[E_0-(1/E)\pm 3V]\}$$
$$\mp P(p/E)\{[1\mp 2E_0(f_T/f_A)][1+G(E,E_0)][\tfrac{1}{2}(1+\gamma)\mp\tfrac{10}{3}VR^2E-(p^2/3)R^2]$$
$$\mp\tfrac{2}{3}a[2E_0-5E\mp 6V]\pm\tfrac{2}{3}b[E_0-E\pm 3V]\mp a'VR[p+(4/p)]\}\cos\theta$$
$$\pm A\{\tfrac{2}{3}(a-b)[E-(1/E)]\mp a'VRp\}\tfrac{1}{2}(3\cos^2\theta-1)\Big). \tag{9}$$

with

$$\gamma=(1-\alpha^2Z^2)^{1/2},\quad V=\alpha Z/2R,$$
$$P=a_1-a_{-1},$$
$$A=1-3a_0,$$
$$a=x\,\mathrm{Re}(f_Ag_W{}^*)/|f_A|^2,\quad g_W=f_W-(f_V/2M), \tag{10}$$
$$b=\mathrm{Re}(f_Ag_T{}^*)/|f_A|^2,\quad g_T=f_T\pm y(f_A/2M),$$
$$a'=\mathrm{Im}(f_Ag_W{}^*)/|f_A|^2.$$

Here F and G are radiative corrections,[12] and V is half of the Coulomb energy at the nuclear surface. p, E, and E_0 are the momentum, energy, and maximum energy of the electron, respectively. The upper sign refers to the electron, and the lower sign refers to the positron. The angle θ is the direction of the electron emission with respect to the nuclear polarization axis; p is the nuclear polarization and A is the nuclear alignment. The a_i's are the magnetic substate densities ($\sum_i a_i=1$). The units $\hbar=m_e=c=1$ are adopted. Assuming small real values for a and b, we can rewrite Eq. (9) as

$$W(\theta)=\mathrm{const}[1+F(E,E_0)][1\pm\tfrac{8}{3}aE]$$
$$\times\{1\mp P(p/E)[1\pm\tfrac{2}{3}(a-b)E]\cos\theta\pm A\tfrac{2}{3}(a-b)E\tfrac{1}{2}(3\cos^2\theta-1)\}. \tag{11}$$

Here the energy-dependent radiative corrections[12] $F(E,E_0)$ and the V and E_0 terms are factorized in the first line, while the minor terms, $1/E$, $\alpha Z/E$, and R^2, are simply omitted. It is noticed here that the factorization of the electromagnetic interactions in Eq. (11) is by no means self-evident without calculations. In fact, this is not the case, if $\int(\vec{\sigma}\cdot\vec{r})\vec{r}$ and $\int\vec{\sigma}r^2$ are explicitly taken into consideration.[5]

Experimentally, the asymmetry of the β-ray angular distribution is given in a form

$$\mathcal{A}=[W(0)-W(\pi)]/[W(0)+W(\pi)]=\mp P(p/E)[1+\alpha_\mp(1-A)E]$$

and

$$\alpha_-(^{12}\mathrm{B})-\alpha_+(^{12}\mathrm{N})=(0.52\pm0.09)\%/\mathrm{MeV}. \tag{12}$$

From this we have

$$f_T/f_A=-(=0.96\pm0.35)\times10^{-3} \tag{13}$$

with $x=1$. It is interesting to note that this is also given by $f_T=-(3.5\pm1.3)f_A/2M$, which is closely related to the strength of weak magnetism. The $f_A/2M$ term in g_A of Eq. (10) has a reasonable magnitude but with an opposite sign from the experimental data, $\alpha_-+\alpha_+$, at the present stage[4] [$\alpha_-(^{12}\mathrm{B})=(0.31\pm0.06)\%/\mathrm{MeV}$ and $\alpha_+(^{12}\mathrm{N})=-(0.21\pm0.07)\%/\mathrm{MeV}$]. A possible small effect of $\int(\vec{\sigma}\cdot\vec{r})\vec{r}$ and $\int\vec{\sigma}r^2$ is under investigation.

Up to this point, we have made the impulse approximation. We can take into account possible meson-exchange effects by regarding x, y, and c in Eqs. (3) as parameters. (y for f_T is no longer unity.) Equivalently, we regard a and b in Eqs. (13) as parameters, and the coupling constants in Eqs. (10) as those modified by the exchange currents. In fact, Eq. (13) coincides with the equation in the elementary-particle approach.[7] In this case, f_T in Eqs. (10) may be written as $f_T{}^{\mathrm{eff}}$ which is not equal to f_T in Eq. (4) anymore. Although the relation of f_T to $f_T{}^{\mathrm{eff}}$ is model dependent, the experimental data in Eq. (12) do not lose their importance by giving us a nonzero value for $f_T{}^{\mathrm{eff}}$, which is of the second class anyhow. This is because the magnitude of a including the mesonic corrections is known in the experiment.[10] Adopting the experimental value [$\frac{16}{3}a=(1.07\pm0.24)\%/\mathrm{MeV}$], we have $f_T{}^{\mathrm{eff}}/f_A{}^{\mathrm{eff}}=-(3.5\pm1.5)/2M$.

We have assumed the point nuclear charge in Eq. (9). The effect of the finite nuclear size can be evaluated in a way described in the references.[5] A contribution of the pseudoscalar term in Eq. (4) is very small as is given elsewhere.[13]

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Hyperfine Structure of $2s$ $^3He^+$ by an Ion-Storage Technique*

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An electrostatic confinement device has provided resonance linewidths $\simeq 1$ kHz for the hyperfine transition $F=1$, $m_F=0$ to $F=0$, in metastable $2s$ $^3He^+$. The state-selection and resonance-detection scheme is the same used in an earlier ion-beam experiment; however, ion storage has yielded a resonance linewidth narrower by a factor of 100. Our result for the $2s$ hyperfine structure is $\Delta\nu_2 = 1083.354\,969(30)$ MHz. Comparison with the $1s$ hyperfine structure yields a test of state-dependent terms in the theory.

It is well known that in the theory of the hyperfine structure of atomic hydrogen, uncertainty in the size of the nuclear-structure correction limits its comparison with experiment to the level of about 3 ppm. This far exceeds the experimental precision of $\simeq 1\times 10^{-6}$ ppm, and, for example, precludes a good test of the quantum-electrodynamic (QED) correction term proportional to $\alpha(Z\alpha)^2$ which is calculated[1] to be 2.27(62) ppm. It is possible, however, to reduce the importance of nuclear structure if one compares the hfs in the $2s$ and $1s$ states. In particular the quantity $D_{21} \equiv (8\Delta\nu_2 - \Delta\nu_1)$, where $\Delta\nu_2$ and $\Delta\nu_1$ are the $2s$ and $1s$ hfs, has a contribution in hydrogen due to the $\alpha(Z\alpha)^2$ term of about 2%, whereas the nuclear structure is not expected to contribute more than about 0.01%. The obvious drawback to this strategy is the requirement for two precision measurements.

In the case of $^3He^+$, in a unique and pioneering experiment, Novick and Commins[2] measured $\Delta\nu_2 = 1083.354\,99(20)$ MHz and, by a novel ion-storage technique, Schuessler, Fortson, and Dehmelt[3] measured $\Delta\nu_1 = 8665.649\,867(10)$ MHz. This yields $D_{21} = 1.1901(16)$ MHz. One sees that the uncertainty in $\Delta\nu_2$ is responsible for virtually all the uncertainty in D_{21}. It was the goal of the present work to determine $\Delta\nu_2$ more accurately. Our experiment uses the same method of state selection and resonance detection as the work of No-