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# Mesonic enhancement of the weak axial charge and its effect on the half-lives and spectral shapes of first-forbidden $J^+ \leftrightarrow J^-$ decays

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ABSTRACT

The effects of the enhancement of the axial-charge matrix element $\gamma_5$ were studied in medium heavy and heavy nuclei for first-forbidden $J^+ \leftrightarrow J^-$ decay transitions using the nuclear shell model. Noticeable dependence on the enhancement $\epsilon_{\rm MEC}$ of the axial-charge matrix element, as well as on the value of the axial-vector coupling constant $g_{\rm A}$ was found in the spectral shapes of $^{93}$Y, $^{95}$Sr, and $^{97}$Y. The importance of the spectrum of $^{138}$Cs in the determination of $g_{\rm A}$ is discussed. Half-life analyses in the $A \approx 95$ and $A \approx 135$ regions were done, and consistent results $g_{\rm A} \approx 0.90$, 0.75, and 0.65, corresponding to the three enhancement scenarios $\epsilon_{\rm MEC} = 1.4$, 1.7, and 2.0, were obtained. Connection to the reactor-antineutrino anomaly is pointed out.



The enhancement of the axial-charge nuclear matrix element (NME) $\gamma_5$ due to nuclear medium effects in the form of meson-exchange currents was first suggested nearly four decades ago [1–3]. An enhancement of 40–70% over the impulse-approximation value was predicted based on chiral-symmetry arguments and soft-pion theorems. This enhancement is fundamental in nature and insensitive to nuclear-structure aspects [4,5]. Systematic shell-model studies of the $\gamma_5$ matrix elements in the $A \approx 16$, $A \approx 40$, and $A \approx 208$ regions indicated enhancements of 60–100% [6–8]. In [9] the exceptionally large enhancement of the $\gamma_5$ NME in heavy nuclei, witnessed in the shell-model studies of Warburton [8], was reproduced by introducing an effective Lagrangian incorporating approximate chiral and scale invariance of the QCD.

The non-trivial dependence of the spectral shapes of the fourth-forbidden non-unique decays of $^{113}$Cd and $^{115}$In on the effective value of the axial-vector coupling constant $g_{\rm A}$ was first pointed out in Ref. [10]. A new method, coined the spectrum-shape method (SSM), where theoretical and experimental spectra are compared was proposed as a complementary way to the half-life comparison method for extraction of the effective value of $g_{\rm A}$. Further investigations in [11] and [12] found that several other non-unique decays exhibit a similar dependence. It was also pointed out in [11–13] that the spectral shapes of many studied decays are practically indifferent to the fine details of the wave functions, making the spectrum-shape method a potentially much more robust tool than the often used half-life method.

Since the $\gamma_5$ NME is one of the two rank-zero matrix elements contributing to first-forbidden $\Delta J = 0$ transitions it plays quite an important role in the decay rates of many of these transitions. Therefore, a significant enhancement of this matrix element can also affect the shapes of the corresponding beta spectra. In the present Letter we investigate the impact of the $\gamma_5$ enhancement on half-lives and shapes of beta spectra for several first-forbidden $J^+ \leftrightarrow J^-$ decay transitions in the $A \approx 95$, $A \approx 135$, and $A \approx 208$ regions. Decays with beta spectra which have a significant dependence on the enhancement of the axial-charge matrix element could potentially be used to extract the enhancement in a similar way to the extraction of the effective value of $g_{\rm A}$. The importance of the consideration of the meson-exchange-current effects on beta spectra used in the reactor-antineutrino analyses and for characterization of the background radiation in rare-event searches is also pointed out.

The half-life of a forbidden non-unique beta decay can be written as

$$t_{1/2} = \kappa/\tilde{C}, \tag{1}$$

where $\kappa = 6147$ s [14] and $\tilde{C}$ is the dimensionless integrated shape function, given by

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$$\tilde{C} = \int_1^{w_0} C(w_e) p w_e (w_0 - w_e)^2 F_0(Z, w_e) dw_e. \quad (2)$$

The shape factor $C(w_e)$ of Eq. (2) contains complicated combinations of both (universal) kinematic factors and nuclear form factors. The nuclear form factors can be related to the corresponding NMEs using the impulse approximation. For the first-forbidden non-unique decays with $J_i = J_f$, considered in this work, the relevant NMEs are those of the transition operators denoted here by $\mathcal{O}(0^-)$, $\mathcal{O}(1^-)$, and $\mathcal{O}(2^-)$. We adopt the expansion of Behrens and Bühring [15], where in the leading order in the non-relativistic reduction there are six matrix elements corresponding to the operators

$$\mathcal{O}(0^-) : g_A(\boldsymbol{\sigma} \cdot \mathbf{p}_e), \quad g_A(\boldsymbol{\sigma} \cdot \mathbf{r}) \quad (3)$$

$$\mathcal{O}(1^-) : g_V \mathbf{p}_e, \quad g_A(\boldsymbol{\sigma} \times \mathbf{r}), \quad g_V \mathbf{r} \quad (4)$$

$$\mathcal{O}(2^-) : g_A[\boldsymbol{\sigma} \mathbf{r}]_2, \quad (5)$$

where $\mathbf{r}$ is the radial coordinate and $\mathbf{p}_e$ is the electron momentum. In this work the enhancement factor of the $\gamma_5$ NME ($\boldsymbol{\sigma} \cdot \mathbf{p}_e$ in the non-relativistic limit) is denoted by $\epsilon_{\rm MEC}$. In addition to these six NMEs, there are three NMEs corresponding to the operators $g_A(\boldsymbol{\sigma} \cdot \mathbf{r})$, $g_A(\boldsymbol{\sigma} \times \mathbf{r})$, and $g_V \mathbf{r}$ with the Coulomb factor $\mathcal{I}(1,1,1,1;r)$ included in the radial integral. The Coulomb factor is given by [15]

$$\frac{2}{3}\mathcal{I}(1,1,1,1;r) = \begin{cases} 1 - \frac{1}{5}\frac{r}{R}, & 0 < r < R \\ \frac{R}{r} - \frac{1}{5}\left(\frac{R}{r}\right)^3, & r > R, \end{cases} \quad (6)$$

where $R$ is the nuclear radius. In this work we include also the next-to-leading-order terms in the Behrens–Bühring expansion [15], which increases the number of NMEs involved in transitions up to 21. The NMEs involved in the transitions can be evaluated using the relation

$$^{V/A}\mathcal{M}^{(N)}_{KLS}(k_e, m, n, \rho) = \frac{\sqrt{4\pi}}{\widehat{J}_i} \sum_{pn} {}^{V/A}m^{(N)}_{KLS}(pn)(k_e, m, n, \rho)(\Psi_f || [c_p^\dagger \tilde{c}_n]_K || \Psi_i), \quad (7)$$

where $^{V/A}m^{(N)}_{KLS}(pn)(k_e, m, n, \rho)$ is the single-particle matrix element, and $(\Psi_f || [c_p^\dagger \tilde{c}_n]_K || \Psi_i)$ is the one-body transition density (OBTD), which contains the nuclear-structure information. The atomic screening effects and radiative corrections are also included in the shape factor. The details of the scope of the formalism can be found from Ref. [13].

In the present work the electron spectra of 16 first-forbidden $\Delta J = 0$ $\beta^-$ transitions were calculated using the NMEs produced by the use of the nuclear shell model. The spectra were calculated using nine different scenarios including all combinations of $\epsilon_{\rm MEC} = 1.40$, 1.70, 2.00 corresponding to 40%, 70% and 100% enhancements of the axial-charge matrix element and $g_A = 0.6$, 1.00, and 1.27 corresponding to the heavily quenched, quenched, and free-nucleon values of the axial-vector coupling constant. The nuclear-structure calculations were done using the shell-model code NuShellX@MSU [16], with appropriate model spaces and Hamiltonians chosen for the three mass regions separately.

For the decay transitions in the mass range $A = 92$–97 a model space including the proton orbitals $0f_{5/2}$, $1p_{3/2}$, $1p_{1/2}$, and $0g_{9/2}$ and the neutron orbitals $1d_{5/2}$, $1d_{3/2}$, and $0s_{1/2}$ were used together with the interaction glbepn [17]. The interaction glbepn is a bare G-matrix interaction which has an adjusted version glepn, where two-body matrix elements from Gloeckner [18] and Ji and Wildenthal [19] have been adopted. The half-lives calculated using the interactions glbepn and glepn agreed for $^{96}$Y to 2%, but for example for $^{93}$Y the modified interaction was unable to reproduce the experimental half-life with any physically meaningful value of $g_A$ and $\epsilon_{\rm MEC}$. The bare G-matrix interaction, on the other hand, was able to reproduce the experimental half-lives, ranging from less than a second to several days. Therefore the half-life analysis was done using this interaction. The decay transitions in the mass range $A = 133$–139 were calculated using a model space spanned by the proton orbitals $0g_{7/2}$, $1d_{5/2}$, $1d_{3/2}$, $2s_{1/2}$, and $0h_{11/2}$ and the neutron orbitals $0h_{9/2}$, $1f_{7/2}$, $1f_{5/2}$, $2p_{3/2}$, $2p_{1/2}$, and $0i_{13/2}$ with the effective interactions jj56pnb [20] and jj56cdb [21]. For the decays of $^{142}$Pr and $^{143}$Pr the dimensions of the problem were so large that truncations became necessary, and no nucleons were allowed on the $\pi 0h_{11/2}$ and $\nu 0i_{13/2}$ orbitals. This decay was not included in the half-life analysis since rather severe truncations were used. For the $A = 211$–213 nuclei the model space consisted of the proton orbitals $0h_{9/2}$, $1f_{7/2}$, $1f_{5/2}$, $2p_{3/2}$, $2p_{1/2}$, and $0i_{13/2}$ and the neutron orbitals $0i_{11/2}$, $1g_{9/2}$, $1g_{7/2}$, $2d_{5/2}$, $2d_{3/2}$, $3s_{1/2}$, and $0j_{15/2}$ with the corresponding Hamiltonian khpe [22]. Half-lives of the decays in the $A \approx 95$ and $A \approx 135$ regions were also calculated and compared with the experimental data. The effects of the mesonic enhancement on the half-lives of the first-forbidden decays in the $^{208}$Pb region have been studied in detail in Ref. [8] and thus we refer to these results for the half-life part. However, the beta spectra of these decays have not been previously published.

**Table 1**
Decays considered in this study. The references used for the half-lives and branching ratios are given in the fourth column.

| Transition | $t_{1/2}$ | BR(%) | Ref. |
|---|---|---|---|
| $^{92}$Rb$(0^-_{g.s.}) \to {}^{92}$Sr$(0^+_{g.s.})$ | 4.492(20) s | 95.2(7) | [23] |
| $^{93}$Y$(1/2^-_{g.s.}) \to {}^{93}$Zr$(1/2^+_1)$ | 10.18(8) h | 2.7(5) | [24] |
| $^{95}$Sr$(1/2^+_{g.s.}) \to {}^{95}$Y$(1/2^-_{g.s.})$ | 23.90(14) s | 55.7(25) | [25] |
| $^{96}$Y$(0^-_{g.s.}) \to {}^{96}$Zr$(0^+_{g.s.})$ | 5.34(5) s | 95.5(5) | [26] |
| $^{97}$Y$(1/2^+_{g.s.}) \to {}^{97}$Zr$(1/2^-_{g.s.})$ | 3.75(3) s | 40(10) | [27] |
| $^{133}$Sn$(7/2^-_{g.s.}) \to {}^{133}$Sb$(7/2^+_{g.s.})$ | 1.46(3) s | 85(3) | [28] |
| $^{134}$Sb$(0^-_{g.s.}) \to {}^{134}$Te$(0^+_{g.s.})$ | 0.78(6) s | 97.6(5) | [29] |
| $^{135}$Te$(7/2^-_{g.s.}) \to {}^{135}$I$(7/2^+_{g.s.})$ | 19.0(2) s | 62(3) | [30] |
| $^{137}$Xe$(7/2^-_{g.s.}) \to {}^{137}$Cs$(7/2^+_{g.s.})$ | 3.818(13) min | 67(3) | [31] |
| $^{138}$Cs$(3^-_{g.s.}) \to {}^{138}$Ba$(3^+_1)$ | 32.5(2) min | 44.0(10) | [32] |
| $^{139}$Ba$(7/2^-_{g.s.}) \to {}^{139}$La$(7/2^+_{g.s.})$ | 83.06(28) min | 70.0(3) | [33] |
| $^{139}$Cs$(7/2^+_{g.s.}) \to {}^{139}$Ba$(7/2^-_{g.s.})$ | 9.27(5) min | 85(3) | [33] |
| $^{142}$Pr$(2^-_{g.s.}) \to {}^{142}$Nb$(2^+_1)$ | 19.12(4) h | 3.7(5) | [34] |
| $^{143}$Pr$(7/2^+_{g.s.}) \to {}^{143}$Nb$(7/2^-_{g.s.})$ | 13.57(2) d | 100 | [35] |
| $^{211}$Pb$(9/2^+) \to {}^{211}$Bi$(9/2^-)$ | 36.1(2) min | 91.32(12) | [36] |
| $^{213}$Bi$(9/2^-) \to {}^{213}$Po$(9/2^+)$ | 45.61(6) min | 65.9(4) | [37] |

The transitions studied in this work, the corresponding branching ratios, and the half-lives of the mother nuclei are listed in Table 1. In Table 2 are given the values of $g_A$ needed to reproduce the experimental half-life with $\epsilon_{\rm MEC} = 1.4$, 1.7, and 2.0 for decay transitions in the mass range $A = 91$–97. The uncertainties in the obtained $g_A$ values stem from the uncertainty in the branching ratio and the half-life, and is given by the standard deviation $\sigma/\sqrt{N}$. The uncertainty coming from the nuclear-structure calculations is at least of the same order as the largest half-life errors, so the total uncertainty is of the order of 30% uniformly for all the decays.

The results for decays in the $A \approx 95$ region are listed in Table 2. When an enhancement of 40% for the axial-charge matrix element is assumed, an effective value close to unity is obtained