Physics Letters B 725 (2013) 153–157

Contents lists available at SciVerse ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Probing the quenching of g_A by single and double beta decays

Jouni Suhonen [a,*], Osvaldo Civitarese [b]

[a] Department of Physics, University of Jyväskylä, P.O. Box 35 (YFL), FI-40014, Finland
[b] Department of Physics, University of La Plata, 67 1900 La Plata, Argentina

ARTICLE INFO

Article history:
Received 14 May 2013
Accepted 19 June 2013
Available online 28 June 2013
Editor: W. Haxton

Keywords:
Quasiparticle random-phase approximation
Two-neutrino double beta decays
Beta decays
Electron-capture decays
Axial-vector coupling constant

ABSTRACT

Ground-state-to-ground-state two-neutrino double beta ($2\nu\beta\beta$) decays and single beta (EC and β^-) decays are studied for the $A = 100$ (^{100}Mo–^{100}Tc–^{100}Ru), $A = 116$ (^{116}Cd–^{116}In–^{116}Sn) and $A = 128$ (^{128}Te–^{128}I–^{128}Xe) nuclear systems by using the proton–neutron quasiparticle random-phase approximation exploiting realistic effective interactions in very large single-particle bases. The aim of this exercise is to see if both the single-beta and double-beta decay observables related to the ground states of the initial, intermediate and final nuclei participant in the decays can be described simultaneously by changing the value of the axial-vector coupling constant g_A. In spite of the very different responses to single and $2\nu\beta\beta$ decays of the considered nuclear systems, the obtained results point consistently to a quenched effective value of g_A that is (slightly) different for the single and $2\nu\beta\beta$ decays.

1. Introduction

Nuclear double beta decays constitute an important issue in the present-day nuclear and neutrino physics due to their connections to many fundamental issues of particle physics. The neutrino properties are closely tied with the neutrinoless ($0\nu\beta\beta$) modes of these decays [1–3]. A massive amount of experimental effort has been, and continues to be invested to determine the corresponding half-lives in many nuclear systems. The aim is a reliable prediction of e.g. the electron–neutrino masses once the nuclear properties, in the form of nuclear matrix elements (NMEs), are under control. Unfortunately, the situation with the NMEs is still rather confusing [4,5], but definite steps forward in this respect have been taken [5]. Central issues in the calculations of the NMEs are many: (a) the effects of the chosen single-particle valence space and orbital occupancies [6–8], (b) a proper account of the shell-closure effects [5,9], (c) deformation effects [10–12] and (d) the effective value of the axial-vector coupling constant g_A of weak interactions [13,14].

In the present Letter we want to address the issue (d) of the above list of important issues of nuclear-structure calculations by using the theoretical framework of the proton–neutron quasiparticle random-phase approximation (pnQRPA) with G-matrix based nuclear forces. The issue of renormalization of g_A is a long-standing one and the renormalization is believed to derive both from truncations in the nuclear-structure calculations and from the interference of non-nucleonic, mainly Δ degrees of freedom. The renormalization stemming from the truncations in the nuclear-structure calculations has been discussed widely in the community involved in the nuclear shell-model calculations [15]. The effects of non-nucleonic degrees of freedom were examined e.g. in Ref. [16], where it was found that the Δ degrees of freedom quench g_A roughly by a factor of $\sqrt{0.6} \approx 0.77$ in the case of the Gamow–Teller (p,n) strength of nuclei.

A consistent formalism to address the renormalization problem of the matrix elements of various operators is the nuclear field theory (NFT) [17]. In the NFT there are two channels through which the renormalization can be achieved, namely (a) by renormalization of the initial and final states on the single-(quasi)particle level by a series of interactions of the P-space states (in the calculational model space) with the Q-space states (states left out of the active model space) mediated by the interaction Hamiltonian [18], and (b) by the particle–hole (two-quasiparticle) channel by involving the (collective) phonons of the system [17]. In particular, following the steps outlined in [18], i.e. including both the quasiparticle and phonon degrees of freedom in the renormalization, one finds for the renormalized matrix element

$$\langle f|M_{\rm GT}|i\rangle_{\rm renormalized} = \langle f|M_{\rm GT}|i\rangle_{\rm bare}\big(1 - F(\omega_{\rm GT})\big), \tag{1}$$

where the function $F(\omega_{\rm GT})$ reduces to the usual polarization function [17] if the energy of the giant Gamow–Teller mode is much larger than the energy difference between initial ($|i\rangle$) and final ($|f\rangle$) state. In this limit the renormalized strength of the low-energy single beta decay amounts to approximately 66 per cent of the bare value [18]. In the present context, this

* Corresponding author.
E-mail addresses: jouni.suhonen@phys.jyu.fi (J. Suhonen), osvaldo.civitarese@fisica.unlp.edu.ar (O. Civitarese).

0370-2693/$ – see front matter
http://dx.doi.org/10.1016/j.physletb.2013.06.042

overall renormalization can be included in the effective value of g_A which we try to extract by comparison with the available data.

In the case of the two-neutrino double-beta ($2\nu\beta\beta$) decay we are dealing with a perturbative expression of the NME with a separable structure, unlike in the case of the $0\nu\beta\beta$ decay where the intermediate neutrino propagator makes the NME non-separable. The separable structure means that the decay involves one nucleon at a time and the decay vertices can be separately renormalized by the above-described procedures. From this point of view the g_A of single beta and $2\nu\beta\beta$ decays renormalize in a similar fashion. However, we still leave open the question of the similar or different renormalizations of these two decay variants and discuss the renormalization of g_A without prejudices by the two methods outlined below.

For the present investigations of the effective value of g_A we have harnessed nuclear systems where both single-beta (EC and β^-) and $2\nu\beta\beta$ half-lives (see [19]) have been measured for transitions between ground states of the participant nuclei. These systems number precisely three, namely the $A = 100$ system containing the triple ^{100}Mo–^{100}Tc–^{100}Ru of nuclei, the $A = 116$ system with the triple ^{116}Cd–^{116}In–^{116}Sn and the $A = 128$ system of ^{128}Te–^{128}I–^{128}Xe nuclei. In the following we investigate how the changes in the effective value of g_A, $g_A^{\rm eff}$, affect the computed $2\nu\beta\beta$ half-lives and the comparative half-lives ($\log ft$ values) of single beta decays. This study will be accomplished by resorting to the following two methods:

- Method MI: In the first method we assume that the g_A of single β and $2\nu\beta\beta$ decays renormalize in the same way. We then go through the following steps; step (i): we start from a given value of $g_A^{\rm eff}$, say $g_A^{\rm eff} = 0.8$, and then extract, by using this value of $g_A^{\rm eff}$, the experimental value of the $2\nu\beta\beta$ NME, NME(2ν, exp) (the value of NME(2ν, exp) is proportional to $(g_A^{\rm eff})^{-2}$ as shown below, in Eq. (2)). Step (ii): we fix the value of the particle–particle strength parameter g_{pp} of the pnQRPA [20] by reproducing the value of NME(2ν, exp). Step (iii): we calculate, by using this extracted value of g_{pp} and the given $g_A^{\rm eff}$, the $\log ft$ values of the EC and β^- transitions corresponding to the decays from the first 1^+ state of the intermediate nucleus to the initial and final nuclei of the $2\nu\beta\beta$ decay, respectively. Step (iv): finally we compare the computed $\log ft$ values with the corresponding experimental ones to see how closely they correspond to each other. In an ideal case there would be one value of $g_A^{\rm eff}$ (and the corresponding g_{pp}) with which both the NME(2ν, exp) and the two $\log ft$ values can be simultaneously reproduced. Such a match for the three observables and for the three different isobaric chains would be highly non-trivial and would point to a common $g_A^{\rm eff}$ for both the single and double beta decays.
- Method MII: In the second method we relax the above-discussed idea of a common $g_A^{\rm eff}$ for both the single and $2\nu\beta\beta$ decays and proceed as follows: step (i): we take both g_{pp} and $g_A^{\rm eff}$ to be independent parameters and try to fix their values by reproducing the experimental $\log ft$ values of both the EC and β^- branches of decay. This is also quite non-trivial if it can be achieved for all the three chains of isobars. Step (ii): by using the g_{pp} value extracted in step (i) we compute the value of the theoretical $2\nu\beta\beta$ NME. Step (iii): we determine a new value of $g_A^{\rm eff}$ such that we can fit the experimental $2\nu\beta\beta$ half-life using the theoretical NME computed in step (ii). Now we can compare the two values of $g_A^{\rm eff}$, extracted in steps (i) and (iii), to see how close they are to each other, i.e. are the $g_A^{\rm eff}$ values of single and $2\nu\beta\beta$ decays the same? In an ideal case the two values of $g_A^{\rm eff}$ would be the same and we would regain the result of the method MI above.

In summary, we want to explore by the above-discussed two methods whether simultaneous description of single-beta and $2\nu\beta\beta$ observables is at all possible or can be achieved by one single value or two different values of $g_A^{\rm eff}$. It is an other matter whether the present results, obtained within the pnQRPA formalism, can be interpreted as generally valid or are just characteristic of the QRPA formalism. We make an attempt to elucidate also this point in this Letter by comparing our results with those extracted from other theory frameworks that are quite different from the pnQRPA.

2. Brief outline of the theoretical framework

Here we present briefly the formalism that we use to compute the double-beta nuclear matrix elements as well as the Gamow–Teller EC and β^- decay amplitudes and the associated $\log ft$ values.

2.1. The $2\nu\beta\beta$-decay amplitude

The $2\nu\beta\beta$-decay half-life, $t_{1/2}^{(2\nu)}$, for a transition from the initial ground state, 0_i^+, to the final ground state, 0_f^+, can be compactly written in the form

$$\left[t_{1/2}^{(2\nu)}\left(0_i^+ \to 0_f^+\right)\right]^{-1} = g_A^4 G_{2\nu} \left|M^{(2\nu)}\right|^2, \tag{2}$$

where g_A is the weak axial-vector coupling constant and $G_{2\nu}$ stands for the leptonic phase-space factor without including g_A in a way defined in [21]. The involved nuclear matrix element is written as

$$M^{(2\nu)} = \sum_{mn} \frac{M_F(1_m^+)\langle 1_m^+|1_n^+\rangle M_I(1_n^+)}{D_m}. \tag{3}$$

The amplitudes connecting the initial ground state 0_i^+ and the final ground state 0_f^+ to the intermediate 1^+ states are

$$M_I(1_n^+) = \left(1_n^+ \left\| \sum_k t_k^{\pm} \boldsymbol{\sigma}_k \right\| 0_i^+\right),$$

$$M_F(1_m^+) = \left(0_f^+ \left\| \sum_k t_k^{\pm} \boldsymbol{\sigma}_k \right\| 1_m^+\right), \tag{4}$$

respectively, where $t_k^{\pm}$ is the flavor-changing operator for the k-th nucleon in the EC or β^- direction. The quantity D_m is the energy denominator containing the average energy of the 1^+ states emerging from the two pnQRPA calculations, one for the initial nucleus and the other for the final nucleus. The denominator can thus be written as

$$D_m = \left(\frac{1}{2}\Delta + \frac{1}{2}\left[E(1_m^+) + \tilde{E}(1_m^+)\right] - M_i c^2\right)/m_e c^2, \tag{5}$$

where Δ is the nuclear mass difference of the $\beta\beta$ initial and final ground states, $\tilde{E}(1_m^+)$ is the energy of the m-th 1^+ state in a pnQRPA calculation based on the initial ground state, $E(1_m^+)$ the same for a calculation based on the final ground state and $M_i c^2$ is the mass energy of the initial nucleus. The quantity $\langle 1_m^+|1_n^+\rangle$ is the overlap between the two sets of 1^+ states [2].

2.2. Beta transition amplitudes

Here we discuss only the Gamow–Teller type of allowed beta decays. The allowed Gamow–Teller beta decay transitions of interest in this work are of the type $1^+ \to 0^+$. For them the $\log ft$ value (comparative half-life) is defined as [22]

$$\log ft = \log(f_0 t_{1/2}) = \log\left[\frac{6147}{B_{\rm GT}}\right],$$
$$B_{\rm GT} = \frac{g_{\rm A}^2}{3}\left|\left(0^+\left\|\sum_k t_k^{\pm}\boldsymbol{\sigma}_k\right\|1^+\right)\right|^2 \qquad (6)$$

for the EC or β^- type of transitions. Here f_0 is the leptonic phase-space factor for the allowed EC or β^- decays as defined in [22].

3. Calculation details

The calculations were done in large single-particle spaces using as starting point a spherical Coulomb-corrected Woods–Saxon (WS) potential with the standard parametrization of Bohr and Mottelson [23], optimized for nuclei near the line of beta stability. The calculations for all the three systems, $A = 100$, $A = 116$ and $A = 128$ were no-core, containing all bound states and, in addition, a number of resonant states with small decay widths. Both the proton and neutron single-particle orbitals numbered 25 for the $A = 100$ and $A = 128$ systems and 26 for the $A = 116$ system. In cases of need small modifications of the WS energies were done at the vicinity of the Fermi surfaces to allow for a better reproduction of the one-quasiparticle type of spectra of the neighboring odd-A nuclei. In particular, the basis set 'EXPWS' of Ref. [24] was adopted for ^{100}Mo.

The Bonn-A G-matrix has been used as the starting point for the nucleon–nucleon interaction and it has been renormalized in the standard way [25,26]: The quasiparticles are treated in the BCS formalism and the pairing matrix elements are scaled by a common factor, separately for protons and neutrons. In practice these factors are fitted such that the lowest quasiparticle energies obtained from the BCS match the experimentally deduced pairing gaps for protons and neutrons respectively. The wave functions of the 1^+ states of the intermediate nuclei have been produced by using the pnQRPA with the particle–hole and particle–particle degrees of freedom [20] included. The particle–hole and particle–particle parts of the proton–neutron two-body interaction are separately scaled by the particle–hole ($g_{\rm ph}$) and particle–particle ($g_{\rm pp}$) parameters. The particle–hole parameter affects the position of the Gamow–Teller giant resonance (GTGR) and its value was fixed by the available systematics [22] on the location of the resonance.

The $g_{\rm pp}$ parameter affects the β^--decay amplitudes of the first 1^+ state in the intermediate nucleus [27] and hence also the decay rates of the $2\nu\beta\beta$ decays. In e.g. [27] the value of this parameter was fixed by the data on β^- decays whereas in many other works, e.g. in [28–31], it was fixed by the data on $2\nu\beta^-\beta^-$-decay rates within the interval $g_{\rm A} = 1.00$–1.25 of the axial-vector coupling constant. In these calculations the uncertainty in the value of $g_{\rm A}$ then induces an interval of acceptable values of $g_{\rm pp}$, the minimum value of $g_{\rm pp}$ related to $g_{\rm A} = 1.00$ and the maximum value to $g_{\rm A} = 1.25$. This is so because the magnitude of the calculated $2\nu\beta\beta$ NME, $M^{(2\nu)}$, decreases with the increasing value of $g_{\rm pp}$ in a pnQRPA calculation [20,25,32] and this magnitude is compared with the magnitude of the experimental NME, NME$(2\nu, \exp) \propto (g_{\rm A})^{-2}$, deduced from the experimental $2\nu\beta\beta$ half-life through (2).

In the present work we adopt a different philosophy and treat both $g_{\rm pp}$ and the axial-vector coupling constant $g_{\rm A}$ as parameters of the calculations. This is done by using the methods MI and MII

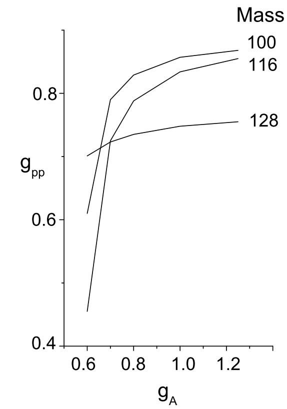

Fig. 1. Dependence of the values of the particle–particle parameter $g_{\rm pp}$ on the value of the axial-vector coupling constant $g_{\rm A}$ for the studied nuclear systems. The steps (i) and (ii) of method MI have been used.

explained in detail in the introduction. To have an idea of the relation between $g_{\rm A}^{\rm eff}$ and $g_{\rm pp}$ we show in Fig. 1 the value of $g_{\rm pp}$ as a function of $g_{\rm A}^{\rm eff}$ by using steps (i) and (ii) of method MI, i.e. for a given value of $g_{\rm A}$, between $g_{\rm A} = 0.60$ and $g_{\rm A} = 1.25$, we extract the NME$(2\nu, \exp)$ from the experimental $2\nu\beta\beta$ half-life (taken from [19]) via the relation (2) and then reproduce the value of NME$(2\nu, \exp)$ by changing the value of $g_{\rm pp}$ in pnQRPA calculations. Due to the above-described dependence of the calculated NME on $g_{\rm pp}$ this parameter is an increasing function of $g_{\rm A}$, as shown in Fig. 1 for the studied nuclear systems. From the figure one notices that for the $A = 128$ system the $g_{\rm pp}$ vs $g_{\rm A}$ curve is rather flat whereas for the other systems the curves are steeper. These qualitatively different behaviors stem from the different responses of the magnitude of the $2\nu\beta\beta$ NME to the increase in the value of $g_{\rm pp}$ in the region of interest. For $A = 128$ the $2\nu\beta\beta$ NME is a fast decreasing function of $g_{\rm pp}$ whereas for the other systems the NME is a relatively flat function of $g_{\rm pp}$.

4. Results and discussion

To have an idea of the differences or similarities in the $2\nu\beta\beta$ decays of the studied three isobaric chains of nuclei one can investigate how large are the contributions to the corresponding NME and from what energy regions of 1^+ states they are coming from. For this purpose we present in Fig. 2 the corresponding cumulative $2\nu\beta\beta$ NMEs for the value $g_{\rm A} = 0.80$ of the axial-vector coupling constant. The cumulative NMEs are very instructive with respect to pinning down contributions appearing in various energy regions of the 1^+ spectrum of the intermediate odd–odd nucleus. It should be noted that while the computed excitation energies of the 1^+ states go beyond 50 MeV, the energy scale of the figure goes only up to 30 MeV since all the NMEs saturate already between 20 MeV and 30 MeV.

From Fig. 2 one perceives that the $2\nu\beta\beta$ NMEs related to the three different nuclear systems are of three distinctive types. Type T1: The NME for the $A = 116$ system is a parade example of *single-state dominance* [33–35]; the final value of the NME is solely determined by the lowest 1^+ contribution to the NME. Type T2: the NME of $A = 128$ behaves quite differently; there are strong, partly canceling contributions between 5 and 10 MeV of excitation. The final NME is only slightly larger than the contribution stemming from the lowest 1^+ state, so the situation appears to be close to single-state dominance, but not due to a sole contribution from the first 1^+ state but rather due to large mutual cancellations at higher energies. The large contributions to the NME in the 5–10 MeV mark the activation of the spin–orbit pairs $0g_{9/2}$–$0g_{7/2}$

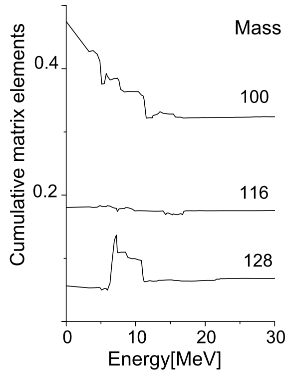

Fig. 2. Cumulative values of the computed $2\nu\beta\beta$ NMEs (3) for the studied nuclear systems with $g_A = 0.80$. The abscissa gives the excitation energy of the 1^+ states in the intermediate nuclei of $2\nu\beta\beta$ decays.

and $0h_{11/2}$–$0h_{9/2}$. Type T3: for the $A = 100$ system the lowest 1^+ contribution is by far the dominant one but beyond that there is a legion of counter-acting cancellations caused by the higher-lying 1^+ states, in accordance with the observations in [34,35] where smaller single-particle valence spaces were used for the calculations.

It should be noted that the above-discussed features of the cumulative NMEs persist for all values of g_A between $g_A = 0.60$ and $g_A = 1.25$. The point in showing the three types (T1–T3) of $2\nu\beta\beta$ NMEs is that the nuclear structures in the three discussed mass regions are different and thus the obtained results for the renormalization of g_A are not tied to only a certain class of nuclei. Hence, it is plausible that the results pertain to weak decay processes of nuclei in general.

Let us now turn to the application of the methods MI and MII (see the end of the introduction for a detailed elaboration on the methods) to access the renormalization of g_A in the discussed three types of nuclear system. First we present the results obtained by using the method MI and its steps (i)–(iv): for a given $0.60 \leqslant g_A \leqslant 1.25$ we determine the experimental $2\nu\beta\beta$ NME and extract the corresponding g_{pp} in a pnQRPA calculation. By using this g_{pp} the $\log ft$ values corresponding to the EC and β^- transitions from the 1_1^+ state of the intermediate nucleus are computed and compared with the available data. In Fig. 3 we show the dependence of the resulting $\log ft$ values on g_A. The dashed lines refer to the EC transitions $^{100}\mathrm{Tc} \rightarrow {}^{100}\mathrm{Mo}$ (the $A = 100$ system), $^{116}\mathrm{In} \rightarrow {}^{116}\mathrm{Cd}$ (the $A = 116$ system) and $^{128}\mathrm{I} \rightarrow {}^{128}\mathrm{Te}$ (the $A = 128$ system), and the solid lines to the β^--decay transition $^{100}\mathrm{Tc} \rightarrow {}^{100}\mathrm{Ru}$ (the $A = 100$ system), $^{116}\mathrm{In} \rightarrow {}^{116}\mathrm{Sn}$ (the $A = 116$ system) and $^{128}\mathrm{I} \rightarrow {}^{128}\mathrm{Xe}$ (the $A = 128$ system). From the figure one can see a clear improvement of the theoretical description of the beta decay $\log ft$ values with decreasing value of g_A. The best values of these quantities are obtained when $g_A \leqslant 0.80$, suggesting that such values of g_A should be used in the calculations of $2\nu\beta\beta$-decay and beta decay rates, at least in the framework of the pnQRPA.

From Fig. 3 one also notices that the correspondence between the computed and experimental $\log ft$ values is not perfect. It may suggest that the g_A^{eff} of the single and double beta decays are in fact different. To study this interesting option of renormalization we resort to the method MII and its steps (i)–(iii), described at the end of the introduction. Our calculations indicate that the step (i) is indeed possible to accomplish in the framework of the pnQRPA, i.e. the $\log ft$ values of both the EC and β^- branches of decay can be reproduced by unique values of g_{pp} and $g_A(\beta)$, listed in columns 2 and 3 of Table 1. Next we take steps (ii) and (iii) and use the extracted g_{pp} to compute the value of the theoretical $2\nu\beta\beta$

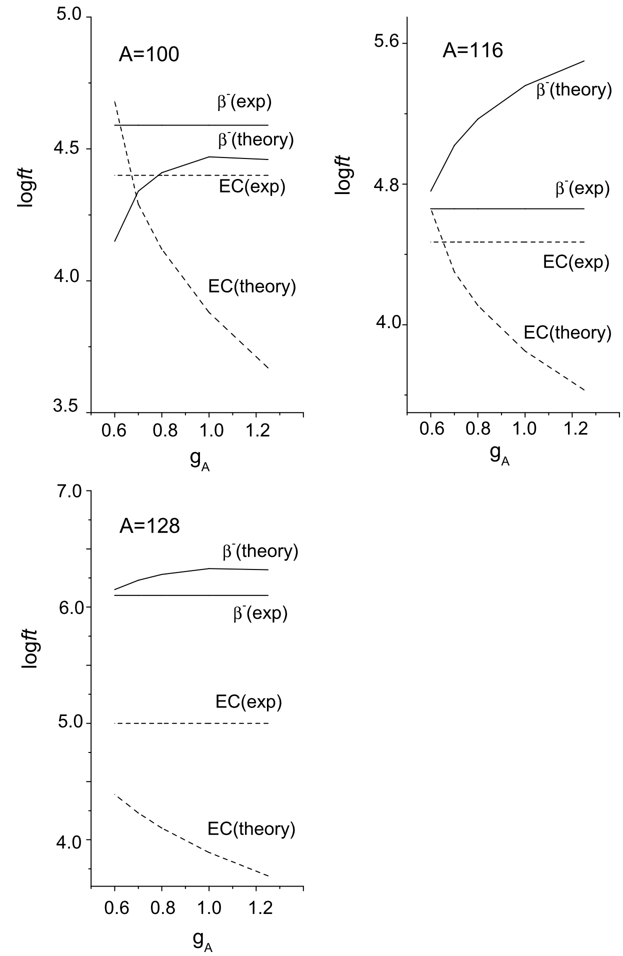

Fig. 3. Dependence on the axial-vector coupling g_A of the computed $\log ft$ values of the electron-capture (EC) (dashed lines) and β^--decay (solid lines) ground-state-to-ground-state transitions in the discussed nuclear systems.

NME and subsequently determine a new value $g_A(\beta\beta)$ (listed in the fourth column of Table 1) by using this NME to fit the experimental $2\nu\beta\beta$ half-life.

From Table 1 one can see that the single-beta and $2\nu\beta\beta$ observables can be exactly fitted with two slightly different values of g_A^{eff} and, by the method MII, with one single value of g_{pp}. For $A = 100, 128$ the $g_A(\beta)$, extracted from the beta decays, is smaller than the one extracted from the double beta decay, whereas for $A = 116$ the reverse is true. This shows that the considered nuclear systems behave differently in terms of nuclear structure, as also suggested by the running-sum analysis of the $2\nu\beta\beta$ NMEs performed in Fig. 2. In all the three cases (types T1–T3 discussed in the context of the running sums) the relative difference between the two g_A values is less than some 20 per cent. It is also interesting to note that contrary to $g_A(\beta)$ of the single beta decay the $g_A(\beta\beta)$ of the $2\nu\beta\beta$ decay is a monotonously decreasing function of the mass number A.

Concerning other calculations, a strong quenching of $g_A^{\mathrm{eff}} \sim 0.6$ was reported in the shell-model calculations of the Gamow–Teller single-beta decay strength in the mass $A = 90$–97 region in Ref. [36]. This is very close to the presently obtained value $g_A(\beta) = 0.59$ for the $A = 100$ case in Table 1. There are also other works where the renormalization of g_A has been discussed for double beta decays: in [13] an attempt was made to fit simultaneously both the single and double beta observables by a single value of g_A^{eff}, i.e. the method MI of the present work was used.

Table 1
Extracted values of g_{pp} and g_A for the three discussed isobaric chains by using the method MII. For comparison are shown the effective values $g_A(\beta\beta)$ obtained in other model frameworks.

A	g_{pp}	$g_A(\beta)$	$g_A(\beta\beta)$			
			Present	pnQRPA [13]	ISM [14]	IBA-2 [14]
100	0.815	0.59	0.75	0.74	0.73	0.55
116	0.515	0.71	0.61	0.84	0.72	0.54
128	0.530	0.33	0.40	–	0.71	0.53

The system $A = 128$ was not included in that work. Results of these calculations are presented in column 5 of Table 1. In [14] the quenching of $g_A(\beta\beta)$ was studied in the framework of IBM-2 and the interacting shell model (ISM) (the shell-model results of [37] were used). There the effective values $g_A^{eff} = 1.269A^{-0.18}$ (IBM-2) and $g_A^{eff} = 1.269A^{-0.12}$ (ISM) were deduced, giving for the specific mass chains the numbers cited in columns 6 and 7 of Table 1. A decreasing trend of the value of g_A with the mass number A prevails in these calculations, though the decrease is not as fast as in the present calculations.

Finally, it is striking that a similar type of quenching of g_A is obtained in many calculations pertaining to nuclear theory frameworks that are apparently very different from each other: the pnQRPA, the IBA-2 and the ISM. It is yet unclear whether the primary reason for this quenching is different in different models or is there something universal and model-independent behind the quenching. Also the relative share of the model-dependent and model-independent contributions to this quenching is unknown and necessitates further investigation in the future.

5. Summary and conclusions

The present investigation about the simultaneous theoretical description of the $2\nu\beta\beta$-decay, EC-decay and β^--decay half-lives for heavy nuclei ($A = 100$–128) is complementary to the earlier work of [27] which discussed the incompatibility of the pnQRPA-computed $2\nu\beta\beta$-decay and single-beta decay half-lives for the same nuclei and for the typical values $g_A = 1.00$–1.25 of the axial-vector coupling constant. The present calculations have been done in very large single-particle models spaces with G-matrix based two-nucleon interactions. By examining the running sums of the $2\nu\beta\beta$ NMEs we conclude that the three discussed isobaric systems behave differently with respect to the $2\nu\beta\beta$ decay (types T1–T3). They also behave differently with respect to single beta decays as visible from Table 1. In spite of the differences in nuclear structure, by letting the value of g_A vary between $g_A = 0.60$–1.25 (method MI) we noticed that the mentioned experimental decay observables can be brought to close correspondence with the computed ones for all the three isobaric chains with the effective values $g_A^{eff} \leqslant 0.80$. This rather striking result is in keeping with other recent calculations and their analyzes which yield values in the range $g_A^{eff} = 0.50$–0.80. Our analysis method MII also suggests that the g_A^{eff} of the single beta decays could be (slightly) different from that of the $2\nu\beta\beta$ decays. How all this affects the very interesting $0\nu\beta\beta$ NMEs is an open question and is beyond the scope of the present Letter.

Acknowledgements

This work was supported by the Academy of Finland under the Finnish Center of Excellence Program 2012–2017 (Nuclear and Accelerator Based Program at JYFL), and by the CONICET of Argentina. O.C. thanks with pleasure the JYFL for the hospitality extended to him.

References

[1] M. Doi, T. Kotani, E. Takasugi, Prog. Theor. Phys. Suppl. 83 (1985) 1–175.
[2] J. Suhonen, O. Civitarese, Phys. Rep. 300 (1998) 123–214.
[3] F.T. Avignone III, S.R. Elliott, J. Engel, Rev. Mod. Phys. 80 (2008) 481–516.
[4] J. Suhonen, O. Civitarese, J. Phys. G: Nucl. Part. Phys. 39 (2012) 085105.
[5] J. Suhonen, O. Civitarese, J. Phys. G: Nucl. Part. Phys. 39 (2012) 124005.
[6] J. Suhonen, O. Civitarese, Phys. Lett. B 668 (2008) 277–281.
[7] J. Suhonen, O. Civitarese, Nucl. Phys. A 847 (2010) 207–232.
[8] J. Suhonen, Nucl. Phys. A 853 (2011) 36–60.
[9] J. Barea, F. Iachello, Phys. Rev. C 79 (2009) 044301.
[10] R. Álvarez-Rodríguez, P. Sarriguren, E. Moya de Guerra, L. Pacearescu, A. Faessler, F. Šimkovic, Phys. Rev. C 70 (2004) 064309.
[11] E. Caurier, F. Nowacki, A. Poves, Eur. Phys. J. A 36 (2008) 195–200.
[12] J. Menéndez, A. Poves, E. Caurier, F. Nowacki, Nucl. Phys. A 818 (2009) 139–151.
[13] A. Faessler, G.L. Fogli, E. Lisi, V. Rodin, A.M. Rotunno, F. Šimkovic, J. Phys. G 35 (2008) 075104.
[14] J. Barea, J. Kotila, F. Iachello, Phys. Rev. C 87 (2013) 014315.
[15] B.A. Brown, B.H. Wildenthal, Ann. Rev. Nucl. Part. Sci. 38 (1988) 29–66.
[16] A. Bohr, B.R. Mottelson, Phys. Lett. 100B (1981) 10–12.
[17] A. Bohr, B.R. Mottelson, Nuclear Structure, vol. II, Benjamin, New York, 1975.
[18] D.R. Bes, O. Civitarese, Phys. Rev. C 81 (2010) 014315.
[19] A.S. Barabash, Phys. Rev. C 81 (2010) 035501.
[20] P. Vogel, M.R. Zirnbauer, Phys. Rev. Lett. 57 (1986) 3148–3151.
[21] J. Kotila, F. Iachello, Phys. Rev. C 85 (2012) 034316.
[22] J. Suhonen, From Nucleons to Nucleus: Concepts of Microscopic Nuclear Theory, Springer, Berlin, 2007.
[23] A. Bohr, B.R. Mottelson, Nuclear Structure, vol. I, Benjamin, New York, 1969.
[24] J. Suhonen, Nucl. Phys. A 700 (2002) 649–665.
[25] J. Suhonen, T. Taigel, A. Faessler, Nucl. Phys. A 486 (1988) 91–117.
[26] J. Suhonen, Nucl. Phys. A 563 (1993) 205–224.
[27] J. Suhonen, Phys. Lett. B 607 (2005) 87–95.
[28] M. Kortelainen, O. Civitarese, J. Suhonen, J. Toivanen, Phys. Lett. B 647 (2007) 128–132.
[29] M. Kortelainen, J. Suhonen, Phys. Rev. C 75 (2007) 051303(R).
[30] M. Kortelainen, J. Suhonen, Phys. Rev. C 76 (2007) 024315.
[31] F. Šimkovic, A. Faessler, V. Rodin, P. Vogel, J. Engel, Phys. Rev. C 77 (2008) 045503.
[32] O. Civitarese, A. Faessler, T. Tomoda, Phys. Lett. B 194 (1987) 11–14.
[33] M. Bhattacharya, et al., Phys. Rev. C 58 (1998) 1247–1256.
[34] O. Civitarese, J. Suhonen, Phys. Rev. C 58 (1998) 1535–1538.
[35] O. Civitarese, J. Suhonen, Nucl. Phys. A 653 (1999) 321–337.
[36] A. Juodagalvis, D.J. Dean, Phys. Rev. C 72 (2005) 024306.
[37] E. Caurier, F. Nowacki, A. Poves, Int. J. Mod. Phys. E 16 (2007) 552–560.