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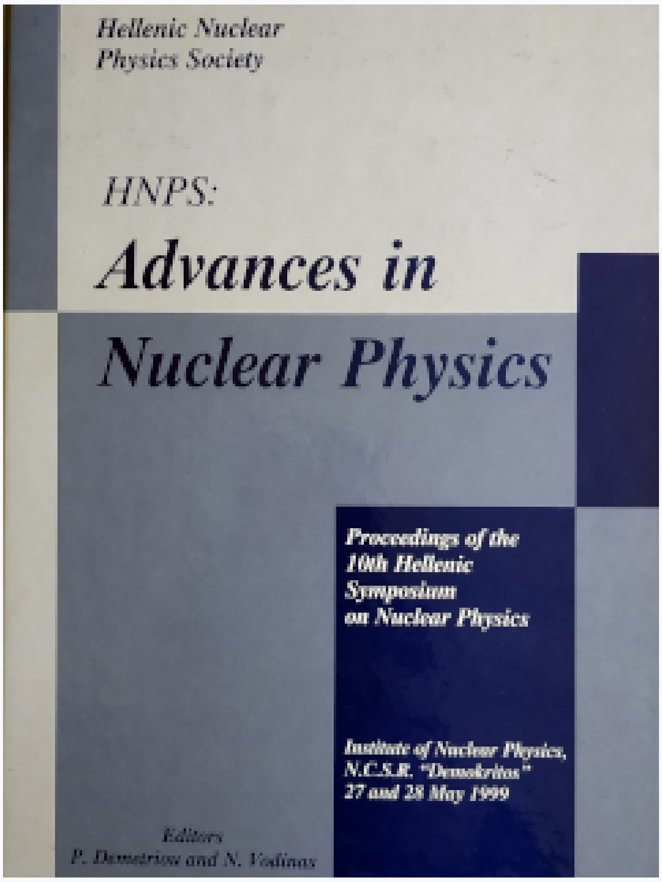

Weak magnetism and pseudoscalar coupling in neutrino mass mechanism of neutrinoless double beta-decay

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Weak magnetism and pseudoscalar coupling in neutrino mass mechanism of neutrinoless double beta-decay

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Abstract

In the calculations of the nuclear matrix element of Majorana neutrino mass mechanism of neutrinoless double beta decay so far only contributions from vector/axial-vector part and weak magnetism of the nuclear current have been included, while other contributions have been neglected. In the present work we are examining the effect of weak magnetism and induced pseudoscalar coupling. We have performed calculations within the proton-neutron renormalized quasiparticle random phase approximation and we have found that these additional contributions of the nucleon current, result in a considerable reduction of the nuclear matrix elements of all nuclei which we have considered. This reduction of the nuclear matrix element makes the extracted limits of the lepton number violating parameters ($< m_\nu >$, $< g >$ and $< \eta_N >$) less stringent yielding the best value for $< m_\nu >$ less than 0.62 eV for ^{76}Ge.

1 Introduction

Most calculations of double beta decay so far have not included higher order terms in the nucleon current, see e.g., Refs. [1–5]. Only the axial and the vector parts have been considered in great detail while an attempt has been made Ref. [6] to estimate weak magnetism. Pseudoscalar coupling and interference between pseudoscalar and axial parts for example have never been considered. While weak magnetism has been shown to be small, other higher order terms are expected to play a more important role. In the case of $2\nu\beta\beta$-decay the momentum transfer in the weak nucleon vertex is restricted by the Q-value for a given process, which is about few MeV. It allows us to neglect safely terms proportional to $|q|/m_\pi$ and $|q|/m_p$ (m_π- mass of pion, m_p- mass of proton and $q = |\vec{q}|$). These terms have also been ignored in the Majorana neutrino mass mechanism of $0\nu\beta\beta$-decay. However, there are reasons to believe that we should take them into account . In the case of $0\nu\beta\beta$-decay, the neutrino is emitted by one nucleon and absorbed by another. The average momentum $< q >$ of the exchange neutrino is expected to be 100 MeV for a mean nucleon-nucleon separation of 2 fm. The situation is even more clear in the case of heavy neutrino exchange. There the mean internucleon distance is considerably smaller and the average momentum $< q >$ is supposed to be considerably larger. Thus such large values of average momentum render it necessary to go beyond the usual approximation of the nucleon current at least for the $0\nu\beta\beta$-decay matrix element for which in general so far only terms of axial and vector contributions have been considered. It is the motivation of this work to include higher order terms of the nuclear current in our calculations. In addition we shall use the renormalised quasiparticle random phase approximation (RQRPA) which incorporates renormalisation effects due to Pauli principle corrections.

2 Theory and calculations

The relevant formalism for neutrinoless double beta decay have been discussed in many papers,see for example the recent review article [7] and references cited there. Here we shall only give the necessary formalae pertaining to this work. In the $0\nu\beta\beta$-decay process one assumes a Hamiltonian of the form,

$$\mathcal{H}^{\beta} = \frac{G_F}{\sqrt{2}} \left[\bar{e}\gamma_\mu(1-\gamma_5)\nu_{eL}\right] J_L^{\mu\dagger} + h.c., \tag{1}$$

where e and ν_{eL} are field operators representing the electron and the left handed electron neutrino, respectively.

Within the impulse approximation the nuclear current J_L^μ in Eq. (1) expressed with nucleon fields Ψ takes the form

$$J_L^{\mu\dagger} = \overline{\Psi}\tau^+ \left[g_V(q^2)\gamma^\mu - ig_M(q^2)\frac{\sigma^{\mu\nu}}{2M}q_\nu - g_A(q^2)\gamma^\mu\gamma_5 + g_P(q^2)q^\mu\gamma_5\right]\Psi, \tag{2}$$

where M is the nucleon mass, $q^\mu = (p-p')_\mu$ is the momentum transferred from hadrons to leptons (p and p' are four momenta of neutron and proton, respectively) and $\sigma^{\mu\nu} = (i/2)[\gamma^\mu, \gamma^\nu]$. $g_V(q^2)$, $g_M(q^2)$, $g_A(q^2)$ and $g_P(q^2)$ are real functions of a Lorenz scalar q^2. These form factors in the zero-momentum transfer limit are known as the vector, weak-magnetism, axial vector and induced pseudoscalar coupling constants, respectively. The induced pseudoscalar coupling constant is given by the Goldberger-Treinman relation

$$\frac{g_P(0)}{g_A(0)} = \frac{2M}{m_\pi^2}. \tag{3}$$

In the previous studies of the neutrino mass mechanism of $0\nu\beta\beta$-decay [1–5] the terms proportional to $g_M(q^2)$ and $g_P(q^2)$ of the nucleon current in Eq. (1) have been neglected and the q^2 dependence of $g_V(q^2)$ and $g_A(q^2)$ was taken to be of dipole shape $1/(1-q^2/\Lambda^2)^2$ with $\Lambda = 0.85GeV$. In this work we shall use the following parametrization of the form factors:

$$
\begin{aligned}
g_V(q^2) &= \frac{g_V}{(1-\frac{q^2}{\Lambda_V^2})^2}, \qquad \Lambda_V^2 = 0.71(GeV)^2, \\
g_M(q^2) &= \frac{\mu_p - \mu_n}{(1-\frac{q^2}{\Lambda_V^2})^2}, \\
g_A(q^2) &= \frac{g_A}{(1-\frac{q^2}{\Lambda_A^2})^2}, \qquad \Lambda_A = 1.086GeV,
\end{aligned} \tag{4}
$$

with

$$
g_V(0) \equiv g_V = 1.0, g_M(0) \equiv g_M = \mu_p - \mu_n, g_A(0) \equiv g_A = 1.254, \tag{5}
$$

where μ_p and μ_n are the magnetic moments of proton and neutron, respectively and Λ_V has been determined by Dumbrajs et al. [8] and is the best fit of the axial form factor for the neutrino reaction $\nu_\mu p \to \mu^+ n$ by Amaldi et al. [9]. The pseudoscalar form factor is determined by the pion pole and its form within the partially conserved axial-vector current hypothesis (PCAC) is given as follows [10]:

$$
g_P(q^2) = \frac{2M g_A(q^2)}{m_\pi^2 - q^2}(1 - \frac{m_\pi^2}{\Lambda_A^2}). \tag{6}
$$

We shall assume a similar relation for high q^2 as well.

For nuclear structure calculations it is necessary to reduce the nucleon current to non-relativistic form. If we neglect small energy transfer between nucleons in the non-relativistic expansion, then the form of the nucleon current coincides with that in Breit frame [11] and we get for the nuclear current,

$$
J^\mu(\vec{x}) = \sum_{n=1}^{A} \tau_n^+ [g^{\mu 0} J^0(q^2) + g^{\mu k} J_n^k(q^2)] \delta(\vec{x} - \vec{r}_n), \qquad k = 1,2,3, \tag{7}
$$

with

$$
\begin{aligned}
J^0(q^2) &= g_V(q^2) \\
\vec{J}_n(q^2) &= g_M(q^2) i \frac{\vec{\sigma}_n \times \vec{q}}{2M} + g_A(q^2)\vec{\sigma} - g_P(q^2) \frac{\vec{q}\, \vec{\sigma}_n \cdot \vec{q}}{2M}.
\end{aligned} \tag{8}
$$

$\vec{r}_i$ is a coordinate of the ith nucleon.

As we had mentioned in the previous section, in the case of $2\nu\beta\beta$-decay the momentum transfer in the weak nucleon vertex is restricted by the Q-value of the process, which is about a few MeV. It allows us to neglect safely terms proportional to $\vec{q}$ in Eq. (7). These terms have been ignored also in the Majorana neutrino mass mechanism of $0\nu\beta\beta$-decay. We shall show in this work that one should take such terms into account. If we assume both outgoing electrons being in the $s_{1/2}$-wave state and consider only the energetically most favored $0_i^+ \to 0_f^+$ transition we obtain for the $0\nu\beta\beta$-decay half-life,

$$[T_{1/2}^{0\nu}]^{-1} = G_{01}|\frac{< m_\nu >}{m_e} M^{light}_{<m_\nu>} + \eta_N M^{heavy}_{\eta_N}|^2. \tag{9}$$

Here $< m_\nu >$ and η_N are the neutrino mass lepton-number non-conserving parameters, m_e is the mass of electron and G_{01} is the integrated kinematical factor. Its numerical values can be found e.g. in Refs. [2,12]. A small difference of G_{01} in the above works arise from the slightly different values of the nuclear radius.

The nuclear matrix elements entering the half-life formula of $0\nu\beta\beta$-decay process in Eq. (9) take the form:

$$M^I_{<m_\nu>,\eta_N} = M^I_{VV} + M^I_{MM} + M^I_{AA} + M^I_{AP} + M^I_{PP}, \quad with I = light, heavy. \tag{10}$$

The partial nuclear matrix elements M^I_{VV}, M^I_{MM}, M^I_{AA}, M^I_{PP} and M^I_{AP} in Eq. (10) originate from the vector, weak magnetism, axial, pseudoscalar coupling and the interference of the axial and pseudoscalar coupling interaction, respectively. They can be expressed in relative coordinates by using the second quatization formalism. We end up at the formula

$$M^I_{type} = \left\langle H^I_{type-F}(r_{12})1 + H^I_{type-GT}(r_{12})\sigma_{12} + H^I_{type-T}(r_{12})\mathrm{S}_{12} \right\rangle \tag{11}$$

with $type = VV, MM, AA, PP, AP$.

The light neutrino-exchange potential $H^{light}_{type-K}(r_{12})$ and heavy neutrino-exchange potential $H^{heavy}_{type-K}(r_{12})$ $(K = F, GT, T)$ are of the following form

$$H^{light}_{type-K}(r_{12}) = \frac{2}{\pi g_A^2}\frac{R}{r_{12}}\int_0^\infty \frac{\sin(qr_{12})}{q + E^m(J) - (E^i + E^f)/2} h_{type-K}(q^2)\, dq, \tag{12}$$

$$H^{heavy}_{type-K}(r) = \frac{1}{m_p m_e}\frac{2}{\pi g_A^2}\frac{R}{r_{12}}\int_0^\infty \sin(qr_{12}) h_{type-K}(q^2)\; q\; dq \tag{13}$$

with

$$h_{VV-F}(q^2) = -g_V^2(q^2), \qquad h_{VV-GT}(q^2) = 0, \qquad h_{VV-T}(q^2) = 0,$$

$$h_{MM-F}(q^2)=0, \quad h_{MM-GT}(q^2)=\frac{2}{3}\frac{g_M^2(q^2)q^2}{4m_p^2}, \quad h_{MM-T}(q^2)=\frac{1}{3}\frac{g_M^2(q^2)q^2}{4m_p^2},$$
$$h_{AA-F}(q^2)=0, \quad h_{AA-GT}(q^2)=g_A^2(q^2), \quad h_{AA-T}(q^2)=0,$$
$$h_{PP-F}(q^2)=0, \quad h_{PP-GT}(q^2)=\frac{1}{3}\frac{g_P^2(q^2)q^4}{4m_p^2}, \quad h_{PP-T}(q^2)=-h_{PP-GT}(q^2),$$
$$h_{AP-F}(q^2)=0, h_{AP-GT}(q^2)=-\frac{2}{3}\frac{g_A(q^2)g_P(q^2)q^2}{2m_p}, h_{AP-T}(q^2)=-h_{AP-GT}(q^2) \quad (14)$$

Here, E^i, E^f and $E^m(J)$ are respectively the energies of the initial, final and intermediate nuclear state with angular momentum J. $R=r_0A^{1/3}$ is the mean nuclear radius, with $r_0=1.1\ fm$.

The short-range correlation function

$$f(r)=1-e^{-\alpha r^2}(1-br^2) \quad (\ \alpha=1.1\ \mathrm{fm}^2 \quad \text{and} \quad b=0.68\ \mathrm{fm}^2\ ), \qquad (15)$$

takes into account the short range repulsion of the nucleons. The exact form of the one-body transition densities to excited states $|J^\pi m_i>$ and $|J^\pi m_f>$ generated respectively from the intial (A,Z) and the final (A,Z+2) QRPA ground states $|0_i^+>$ and $0_f^+>$ within the pn-RQRPA can be found together with other details of the nuclear structure model in Refs. [5,12–14]. As it is expected if the additional current contributions are left out then we get the known formulae,

$$M^I_{<m_\nu>,\eta_N}=M^I_{VV}+M^I_{AA} \quad I=light, heavy, \qquad (16)$$

where $M^I_{VV}\equiv-\frac{g_V^2}{g_A^2}M_F^{0\nu-I}$ and $M^I_{AA}\equiv M_{GT}^{0\nu-I}$.

3 Discussion and conclusions

Our numerical values of the contributions of each term to the nuclear matrix elements for the light and heavy Majorana neutrino exchange modes are listed in Table 1 for all mass numbers A=76 up to A=150 which udergo double beta decay. For the A=76 and 128 systems they are also shown in a histogramm in Fig. 1. For the light neutrino exchange the weak magnetism is very small.The other contributions are significant. In fact the vector and the pseudoscalar parts together are almost equal to the interference term which however has oposite sign. As it is more clearly shown in the histogramm of Fig. 2 this brings a reduction to the nuclear matrix element of about 25 in all nuclear systems. The situation is even more pronounced in the heavy neutrino exchange. Here the weak magnetism contribution is much stronger and of oposite sign and this brings down the matrix element by factors 2 to 4.

In order to study the sensitivity of each nucleus to the light and heavy neutrino mass we have introduced sensitivity parameters for a given isotope which depend

Table 1
Nuclear matrix elements for the light and heavy Majorana neutrino exchange modes of $0\nu\beta\beta$-decay for the nuclei studied in this work. G_{01} is the integrated kinematical factor for $0^+ \to 0^+$ transition. $\zeta_{<m_\nu>}$ and ζ_{η_N} denote respectively the sensitivity of a given nucleus to the light and heavy neutrino mass.

	$(\beta\beta)_{0\nu}$ – *decay* : $0^+ \to 0^+$ transition								
M. E.	^{76}Ge	^{82}Se	^{96}Zr	^{100}Mo	^{116}Cd	^{128}Te	^{130}Te	^{136}Xe	^{150}Nd
	light Majorana neutrino (I=light)								
M^I_{VV}	0.80	0.74	0.45	0.82	0.50	0.75	0.66	0.32	1.14
M^I_{AA}	2.80	2.66	1.54	3.30	2.08	2.21	1.84	0.70	3.37
M^I_{PP}	0.23	0.22	0.15	0.26	0.15	0.24	0.21	0.11	0.35
M^I_{AP}	-1.04	-0.98	-0.65	-1.17	-0.69	-1.04	-0.91	-0.48	-1.53
$M^I_{VV} + M^I_{AA}$	3.60	3.40	1.99	4.12	2.58	2.96	2.50	1.02	4.51
M^I_{mass}	2.80	2.64	1.49	3.21	2.05	2.17	1.80	0.66	3.33
	heavy Majorana neutrino (I= heavy)								
M^I_{VV}	23.9	22.0	16.1	28.3	17.2	25.8	23.4	13.9	39.4
M^I_{MM}	-55.4	-51.6	-38.1	-67.3	-39.8	-60.4	-54.5	-31.3	-92.0
M^I_{AA}	106.	98.3	68.4	123.	74.0	111.	100.	58.3	167.
M^I_{PP}	13.0	12.0	9.3	16.1	9.1	14.9	13.6	7.9	23.0
M^I_{AP}	-55.1	-50.7	-41.1	-70.1	-39.0	-64.9	-59.4	-34.8	-101.
$M^I_{VV} + M^I_{AA}$	130.	120.	84.5	151.	91.1	137.	123.	72.3	206.
M^I_{mass}	32.6	30.0	14.7	29.7	21.5	26.6	23.1	14.1	35.6
	sensitivity to neutrino mass signal								
$G_{01} \times 10^{15} y$	7.93	35.2	73.6	57.3	62.3	2.21	55.4	59.1	269.
$\zeta_{<m_\nu>}(Y)$	2.49	4.95	4.04	7.69	5.11	1.02	4.24	1.60	17.3
$\zeta_{\eta_N}(Y)$	2.90	5.64	3.98	7.10	5.36	1.25	5.45	3.43	18.5

only on the corresponding nuclear matrix element and the kinematical phase-space factor. There are as follows:

$$\zeta_{<m_\nu>}(Y) = 10^7 \, |M^{light}_{<m_\nu>}| \, \sqrt{G_{01}} \tag{17}$$

$$\zeta_{\eta_N}(Y) = 10^6 \, |M^{heavy}_{<m_\nu>}| \, \sqrt{G_{01}} \tag{18}$$

The numerical values of $\zeta_{<m_\nu>}(Y)$ and $\zeta_{\eta_N}(Y)$ for the nuclear systems considered in this work are also listed in Table 1 and shown in Fig 2. Large numerical values of these parameters charachterise those $0\nu\beta\beta$-decay isotopes, which are the most

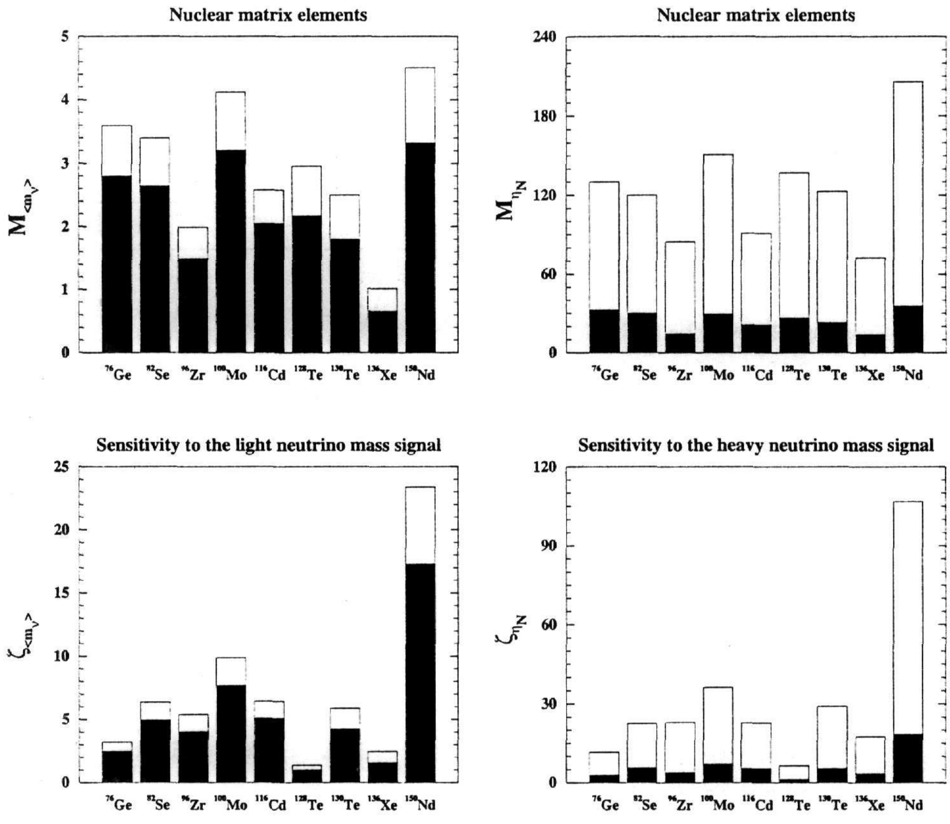


Fig. 1. Calculated light and heavy neutrino exchange $0\nu\beta\beta$-decay nuclear matrix elements for A=76 and 128 systems. The partial matrix elements M_{VV}, M_{AA}, M_{MM}, M_{PP} and M_{AP} originate from vector, axial, weak magnetism, pseudoscalar coupling and the interference of the axial and pseudoscalar coupling interaction, respectively. $M_{<m_\nu>}$ and M_{η_N} are $0\nu\beta\beta$-decay matrix elements associated with $< m_\nu >$ and η_N lepton number violating parameters, respectively.

promising candidates for searching the lepton number violating signal. These sensitivity parameters can be used as a guide by the experimentalists in planning the $0\nu\beta\beta$-decay experiments. Our results show that the A=150 system is the most sensitive for both the light and the heavy neutrino exchange. This should be taken into account together with other microscopic and macroscopic factors for building a $0\nu\beta\beta$-detectors. In Table 2 we list the upper limits of the lepton number non conserving parameters corresponding to the best presently availlable experimental values of the lower half-life limits for a given isotope. To see more clearly the influence of the additional terms we show a comparisson in a histogramm in Fig. 3. The best values are those of the A=76 isotope. With these best values we have derived the best lower half-life limits of the other isotopes listed in Table 2 and shown by open bars in Fig.3. By glancing at the histogramm it is obvious that most experiments have a long way to go to reach the A=76 target limit.

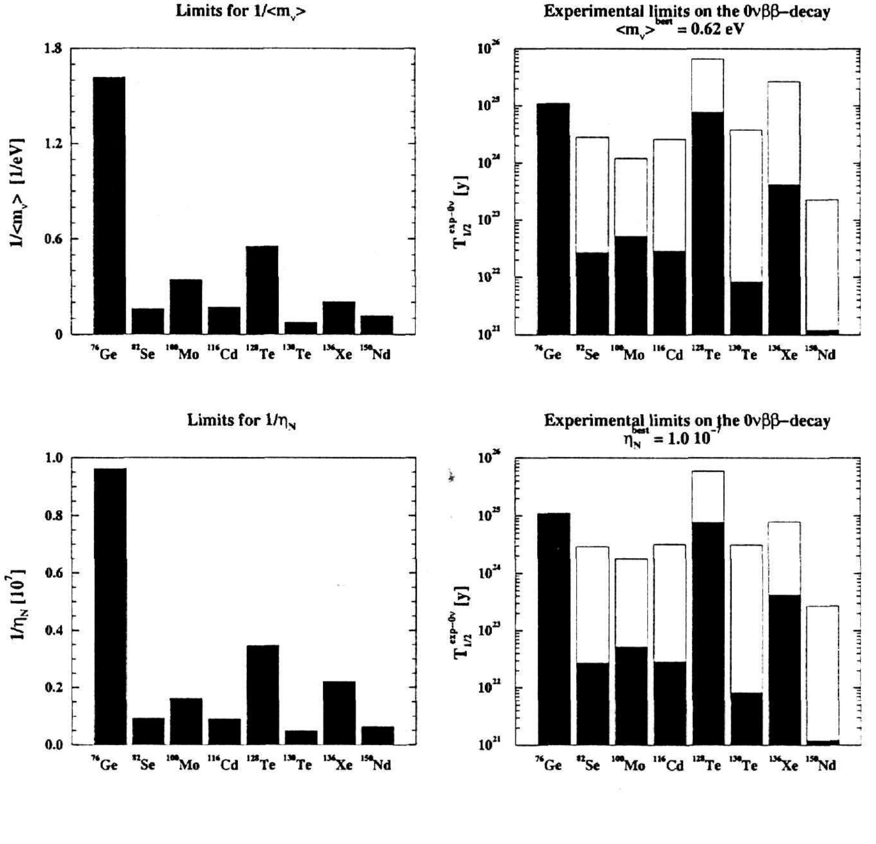


Fig. 2. Calculated nuclear matrix elements $M_{<m_\nu>}$ and M_{η_N} and sensitivities $\zeta_{<m_\nu>}$ and ζ_{η_N} for the experimentally interesting A=76, 82, 96, 100, 116, 128, 130, 136 and 150 nuclear systems. The black bars show the results for the total interaction the open bars the results without the pseudoscalar and weak-magnetism terms.

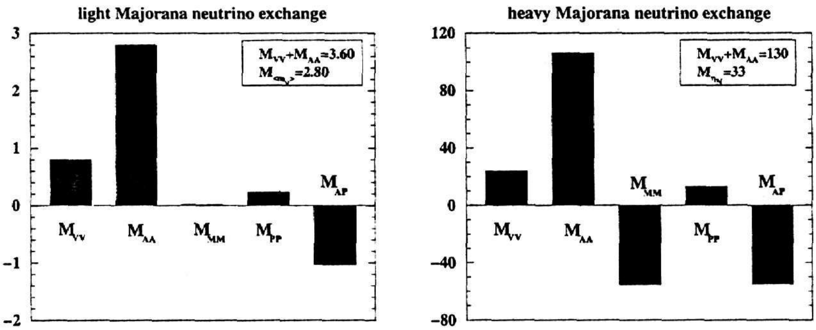
Table 2

The present state of the Majorana neutrino mass in $\beta\beta$-decay experiments. $T_{1/2}^{exp-0\nu}$(present) is the best presently available lower limit on the half-life of the $0\nu\beta\beta$-decay for a given isotope and $< m_\nu >$ and η_N are the corresponding upper limits on the lepton number non-conserving parameters. $T_{1/2}^{exp-0\nu}(< m_\nu >^{best})$, $T_{1/2}^{exp-0\nu}(\eta_N^{best})$ are calculated half-lifes of $0\nu\beta\beta$-decay assuming $< m_\nu >=< m_\nu >^{best}$ and $\eta_N = \eta_N^{best}$. Here $< m_\nu >^{best}= 0.62eV$, $\eta_N^{best} = 1.0 \times 10^{-7}$ are the best limits deduced from the ^{76}Ge experiments.

Nucleus	^{76}Ge	^{82}Se	^{96}Zr	^{100}Mo	^{116}Cd
$T_{1/2}^{exp-0\nu}$(present) [y]	1.1×10^{25}	2.7×10^{22}	3.9×10^{19}	5.2×10^{22}	2.9×10^{22}
Ref.	[15]	[16]	[17]	[18]	[19]
$< m_\nu >$ [eV]	0.62	6.3	203.	2.9	5.9
$T_{1/2}^{exp-0\nu}(< m_\nu >^{best})$ [y]	1.1×10^{25}	2.8×10^{24}	4.2×10^{24}	1.2×10^{24}	2.6×10^{24}
η_N	1.0×10^{-7}	1.1×10^{-6}	4.0×10^{-5}	6.2×10^{-7}	1.1×10^{-6}
$T_{1/2}^{exp-0\nu}(\eta_N^{best})$ [y]	1.1×10^{25}	2.9×10^{24}	5.8×10^{24}	1.8×10^{24}	3.2×10^{24}

Nucleus	^{128}Te	^{130}Te	^{136}Xe	^{150}Nd
$T_{1/2}^{exp-0\nu}$(present) [y]	7.7×10^{24}	8.2×10^{21}	4.2×10^{23}	1.2×10^{21}
Ref.	[20]	[21]	[22]	[23]
$< m_\nu >$ [eV]	1.8	13.	4.9	8.5
$T_{1/2}^{exp-0\nu}(< m_\nu >^{best})$ [y]	6.6×10^{25}	3.8×10^{24}	2.7×10^{25}	2.3×10^{23}
η_N	2.9×10^{-7}	2.0×10^{-6}	4.5×10^{-7}	1.6×10^{-6}
$T_{1/2}^{exp-0\nu}(\eta_N^{best})$ [y]	5.9×10^{25}	3.1×10^{24}	7.9×10^{24}	2.7×10^{23}

$^{76}Ge \rightarrow {}^{76}Se + 2e^-$



$^{128}Te \rightarrow {}^{128}Xe + 2e^-$

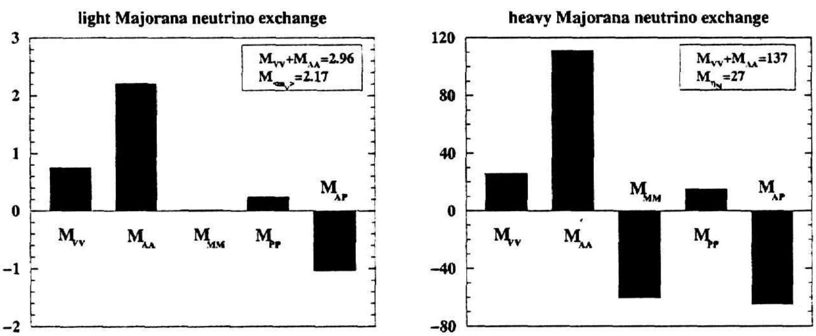


Fig. 3. The sensitivity of different experiments to the lepton-number violating parameters $< m_\nu >$ and η_N are illustrated by histograms on the left side. The best presently available lower limits on the $0\nu\beta\beta$-decay half-life $T_{1/2}^{exp-0\nu}$ are displayed by black bars on the histograms on the right side. The open bars in these histograms indicate the half-life limits $T_{1/2}^{exp-0\nu}(< m_\nu >^{best})$, $T_{1/2}^{exp-0\nu}(\eta_N^{best})$ required by a given experiment to reach the presently best limit of $< m_\nu >$ and η_N.

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