The light neutrino exchange mechanism of the $0\nu\beta\beta$-decay with left- and right-handed lepton and hadronic currents revisited.

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The extension of Majorana neutrino mass mechanism of the neutrinoless double-beta decay ($0\nu\beta\beta$) with the inclusion of right-handed lepton and hadronic currents is revisited. While only the exchange of light neutrinos is assumed, the $s_{1/2}$ and $p_{1/2}$-states of emitted electrons as well as recoil corrections to the nucleon currents are taken into account. Within the standard approximations the decay rate is factorized into a sum of products of kinematical phase-space factors, nuclear matrix elements and the fundamental parameters that characterize the lepton number violation. Unlike in the previous treatments the induced pseudoscalar term of hadron current is included, resulting in additional nuclear matrix elements. An improved numerical computation of the phase-space factors is presented, based on the exact Dirac wave functions of the $s_{1/2}$ and $p_{1/2}$ electrons with finite nuclear size and electron screening taken into account. The dependence of values of these phase-space factors on the different approximation schemes used in evaluation of electron wave functions is discussed. The upper limits for effective neutrino mass and the parameters $(\lambda)$ and $(\eta)$ characterizing the right-handed current mechanism are deduced from data on the $0\nu\beta\beta$-decay of $^{76}$Ge and $^{136}$Xe using nuclear matrix elements calculated within the nuclear shell model and quasiparticle random phase approximation. The differential decay rates, i.e. the angular correlations and the single electron energy distributions for various combinations of the total lepton number violating parameters that can help to disentangle the possible mechanism are described and discussed.

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I. INTRODUCTION

The neutrinoless double-beta ($0\nu\beta\beta$) decay is a process in which an atomic nucleus with Z protons decays to another one with two more protons and the same mass number A, by emitting two electrons and nothing else

$$(A, Z) \rightarrow (A, Z + 2) + 2e^-.$$ (1)

Observing the $0\nu\beta\beta$-decay guaranties that neutrinos are massive Majorana particles - it means that neutrino is identical to its own antiparticle [1]. $0\nu\beta\beta$-decay violates total lepton number conservation and is forbidden in the Standard Model.

When the light-neutrino exchange produced by left-handed currents is the driving mechanism for $0\nu\beta\beta$-decay, the relation between the effective Majorana neutrino mass and the inverse half-life of the $0\nu\beta\beta$-decay can be written as [2]

$$\left[ T_{1/2}^{0\nu} \right]^{-1} = G^{\nu\nu}(Q, Z) g_A^4 |M^{\nu\nu}|^2 \frac{|m_{\beta\beta}|^2}{m_e^2},$$ (2)

where $G^{\nu\nu}(Q, Z)$, $g_A$ and $M^{\nu\nu}$ represent the phase-space factor, the axial-vector coupling constant and the nuclear matrix element of the process, respectively. In that case the ultimate goal of the search for $0\nu\beta\beta$-decay is the determination of the effective Majorana neutrino mass,

$$m_{\beta\beta} = U_{e1}^2 m_1 + U_{e2}^2 m_2 + U_{e3}^2 m_3.$$ (3)

Here, $U_{ei}$ and $m_i$ (i=1,2,3) are elements of Pontecorvo-Maki-Nakagawa-Sakata neutrino mixing matrix and masses of neutrinos, respectively.

An improved calculation of $G^{\nu\nu}$ in that case, taking into account the electron Dirac wave functions with finite nuclear size and electron screening, was performed in [3]. The main theoretical uncertainty is represented in the computed values of the nuclear matrix elements. There is a factor of 2-3 difference between the different methods of calculations of the $M^{\nu\nu}$.

The left-right symmetric theories [4, 5] provide a natural framework to understand the origin of neutrino Majorana masses. In general one cannot predict the scale where the left-right symmetry is realized, but it is natural to assume that it is as low as ~ a few TeV which can affect the $0\nu\beta\beta$ decay rate significantly.

In the left-right symmetric theories in addition to the left-handed V-A weak currents also leptonic and hadronic right-handed V+A weak currents are present. In that case new mechanism of the $0\nu\beta\beta$-decay need to be considered. In the past the $0\nu\beta\beta$-decay rate in the presence of the right-handed leptonic and hadronic currents was discussed in [6, 7]. Recently, contributions to the $0\nu\beta\beta$-decay in a TeV-scale left-right symmetric models for type-I seesaw dominance were revisited [8-11]. By making a qualitative analysis without considering relevant phase-space factors and nuclear matrix elements it was found that $W_L-W_R$ exchange ($\lambda$ mechanism) and $W_L-W_R$ mixing ($\eta$ mechanism) could give dominant con-
tribution to the $0\nu\beta\beta$-decay amplitude by assuming a wide particle physics available parameter space including left-right neutrino mixing \[10 \text{ and } 11\]. We note that the discovery of a non-standard $0\nu\beta\beta$-decay mechanism such as a right-handed current would rule out most models of baryogenesis at scales above 40 TeV \[12\].

The goal of this paper is to revisit the $0\nu\beta\beta$-decay mechanism due to the right-handed currents by considering exact Dirac wave function of electrons and the higher order terms of nucleon current. We note that in \[13\] the effect of induced pseudoscalar term of nucleon current was neglected and phase space factors were expressed using approximate electron wave functions for a uniform charge distribution in a nucleus by keeping only the lowest terms in the power expansion in $r$. In that context the subject of interest is the comparison of the power expansion versus the exact treatment and the finite nuclear size effects. In this work the newly derived decay rate will be written in a compact form and the corresponding nuclear matrix elements will be presented by assuming the usual closure approximation for intermediate nuclear states. The phase-space factors will be evaluated by using the exact Dirac wave functions with finite nuclear size and electron screening. The differential characteristics, i.e. the angular correlations and the single electron energy distributions will be described and discussed, and the decay rates and updated constraints on the lepton number violating parameters for different combinations of the total lepton number violating parameters will be recalculated. This will make it possible to judge the importance of the $\lambda$ and $\eta$ mechanism.

II. ELECTRON WAVE FUNCTIONS

An important ingredient in the calculation of the electron energy spectrum is the radial electron wave function distorted by the Coulomb field. We adopt the Dirac wave functions in a central field,

$$\Psi(\varepsilon, r) = \sum_{\kappa n\mu} \left( \frac{g_{k}(\varepsilon, r)\chi_{k\mu}(\hat{r})}{f_{k}(\varepsilon, r)\chi_{-k\mu}(\hat{r})} \right), \quad (4)$$

given, e.g. by Rose \[13\]. Here, $\varepsilon$ and $r$ stand for energy and position vector of the electron, respectively, $r = |r|$ and $\hat{r} = r/r$. The index $\kappa$ takes positive and negative integer values ($\kappa = \pm k; \ k = 1, 2, 3, \cdots$). Total angular momentum is given as $f_{k} = |\kappa| - 1/2$ while orbital angular momentum takes values

$$l_{\kappa} = |\kappa| - 1 \quad \text{for} \quad \kappa < 0$$

$$= \kappa \quad \text{for} \quad \kappa > 0 \quad (5)$$

Radial wave functions $g_{k}(\varepsilon, r)$ and $f_{k}(\varepsilon, r)$ obey the radial Dirac equations,

$$\frac{dg_{k}}{dr} = \frac{-\kappa + 1}{r}g_{k} + (\varepsilon - V(r) + m_{e})f_{k},$$

$$\frac{df_{k}}{dr} = -(\varepsilon - V(r) - m_{e})g_{k} + \frac{\kappa - 1}{r}f_{k}, \quad (6)$$

where $V(r)$ is the central Coulomb potential. The natural units $\hbar = c = 1$ are used.

The electron wave function expressed in terms of spherical waves is given by

$$\Psi(\varepsilon, r) = \Psi^{(s_{1/2})}(\varepsilon, r) + \Psi^{(p_{1/2})}(\varepsilon, r) + \cdots \quad (7)$$

Here, superscript displays the orbital angular momentum ($l_{\kappa} = 0, 1, 2, \cdots$) in a spectroscopic notation ($l_{\kappa} = s, p, d, \cdots$) and the total angular momentum $j_{\kappa}$. The states of particular interest in our calculations are:

$$\Psi^{(s_{1/2})}(\varepsilon, r) = \left( \begin{array}{c} g_{-1}(\varepsilon, r)\chi_{s} \\ f_{+1}(\varepsilon, r)(\sigma \cdot \hat{p})\chi_{s} \end{array} \right),$$

$$\Psi^{(p_{1/2})}(\varepsilon, r) = \left( i \begin{array}{c} g_{+1}(\varepsilon, r) (\sigma \cdot \hat{r})\chi_{s} \\ -f_{-1}(\varepsilon, r)(\sigma \cdot \hat{r})\chi_{s} \end{array} \right), \quad (8)$$

where $\hat{p} = p/p$ and $p$ is the electron momentum. The asymptotic behavior of the radial wave functions for large values of $pr$ is given by

$$\left( \begin{array}{c} g_{k}(\varepsilon, r) \\ f_{k}(\varepsilon, r) \end{array} \right) \approx \frac{1}{pr} \left( \begin{array}{c} \sqrt{\varepsilon + m_{e}} \sin(pr - l_{\kappa}^{2} + \alpha Z_{f}^{2} \log 2pr) \\ \sqrt{\varepsilon - m_{e}} \cos(pr - l_{\kappa}^{2} + \delta_{k} + \alpha Z_{f}^{2} \log 2pr) \end{array} \right). \quad (9)$$

$Z_{f}$ is the charge of the final system which generates the potential $V(r)$.

In what follows, different approximation schemes for the calculation of radial wave functions $g_{k1}$ and $f_{k1}$ associated with emitted electron in the $s_{1/2}$ and $p_{1/2}$ wave states are briefly presented.

The approximation scheme A): The relativistic electron wave function in a uniform charge distribution in nucleus is considered. The lowest terms in the power expansion in $r$ are taken into account. The radial wave functions take the form

$$\left( \begin{array}{c} g_{-1}(\varepsilon, r) \\ f_{+1}(\varepsilon, r) \end{array} \right) \approx \sqrt{F_{0}(Z_{f}, \varepsilon)} \left( \begin{array}{c} \sqrt{\varepsilon + m_{e}} \\ \sqrt{\varepsilon - m_{e}} \end{array} \right),$$

$$\left( \begin{array}{c} g_{+1}(\varepsilon, r) \\ f_{-1}(\varepsilon, r) \end{array} \right) \approx \sqrt{F_{0}(Z_{f}, \varepsilon)} \times \left( \begin{array}{c} \sqrt{\varepsilon + m_{e}}\left[ \alpha Z_{f}^{2} + (\varepsilon + m_{e})r/3 \right] \\ -\sqrt{\varepsilon - m_{e}}\left[ \alpha Z_{f}^{2} + (\varepsilon - m_{e})r/3 \right] \end{array} \right). \quad (10)$$

Here, $F_{k-1}$ (for $k = 1, 2, \cdots$) is given by

$$F_{k-1} = \left[ \frac{\Gamma(2k + 1)}{\Gamma(k)(1 + 2\gamma_{k})} \right]^{2} (2pr)^{2(\gamma_{k} - k)}e^{\gamma_{k}y} \times \Gamma(\gamma_{k} + iy)^{2}, \quad (11)$$
where
\[ \gamma_k = \sqrt{k^2 - (\alpha Z_f)^2} \]
\[ y = \frac{\alpha Z_f}{\gamma}. \]

This approximation scheme was commonly used in the past calculations of the phase-space integrals of double beta decay processes \[4\].

The approximation scheme B): The analytical solution of the Dirac equation for the point-like nucleus is considered \[14\]. Radial wave functions take then the form
\[ g_\kappa(\varepsilon, r) = \frac{\kappa}{k_p r} \sqrt{\frac{\varepsilon + m_e}{2\varepsilon}} \frac{\Gamma(1 + \kappa + iy)}{\Gamma(1 + 2\kappa)} (2pr)^\kappa e^{\gamma y/2} \]
\[ \Im \left\{ e^{i(pr+\xi)} F_1(\gamma_k - iy, 1 + 2\kappa, -2ipr) \right\}, \]
\[ f_\kappa(\varepsilon, r) = \frac{\kappa}{k_p r} \sqrt{\frac{\varepsilon - m_e}{2\varepsilon}} \frac{\Gamma(1 + \kappa + iy)}{\Gamma(1 + 2\kappa)} (2pr)^\kappa e^{\gamma y/2} \]
\[ \Re \left\{ e^{i(pr+\xi)} F_1(\gamma_k - iy, 1 + 2\kappa, -2ipr) \right\}, \]
\[ (13) \]

with
\[ e^{-2i\xi} = \frac{\gamma_k - iy}{\kappa - iy m_e/\varepsilon}. \]

Here, \( F_1(a, b, z) \) is the confluent hypergeometric function.

The approximation scheme C): The exact Dirac wave functions with finite nuclear size, which is taken into account by a uniform charge distribution in a sphere of nuclear radius \( R \), are assumed \[3\]. The numerical calculation can be accomplished by the subroutine package RADIAL \[15\], where the input central potential is given by
\[ V(r) = \begin{cases} \frac{\alpha Z_f}{2\pi r} \left( 3 - \left( \frac{r}{R} \right)^2 \right) & \text{for } r \leq R \\ \frac{\alpha Z_f}{2\pi} & \text{for } r > R \end{cases} \]
\[ (15) \]

Here, \( R = r_0 A^{1/3} \) with \( r_0 = 1.2 \) fm. In the code the radial Dirac equations are solved by using piecewise exact power series expansion, which are summed up to a prescribed accuracy so that truncation errors can be avoided completely.

The approximation scheme D): The exact Dirac wave functions with finite nuclear size and electron screening are used \[3\]. The effect of screening of atomic electrons is taken into account by the Thomas-Fermi approximation. It uses the solution of the Thomas-Fermi equation,
\[ \frac{d^2 \varphi}{dx^2} = \frac{\varphi^{3/2}}{\sqrt{x}}, \]
\[ (16) \]
\(f_{+1}(e)\). These differences are apparent especially at lower energies. We notice also a rather good agreement between results for wave functions of B and C in general. The screening of atomic electrons affects mostly the \(p_{1/2}\) wave functions, but is essentially negligible for the \(s_{1/2}\) states.

III. DECAY RATE FOR THE NEUTRINOLESS DOUBLE-BETA DECAY

One of the most prominent new physics model that incorporates the LNV and which leads to potentially observable rates for the \(0\nu\beta\beta\)-decay is the minimal left-right symmetric model (LRSM) \([4,5]\) which extends the standard model gauge symmetry to the group \(SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}\). The right-handed neutrinos necessary appear here as a part of the \(SU(2)_R\) doublets. The lepton multiplets \(L_i = (\nu_i, l_i)\) are characterized by the quantum numbers \(Q_{LL} = (1/2, 0, -1)\) and \(Q_{LR} = (0, 1/2, -1)\) under \(SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}\). The Higgs sector contains a bidoublet \(\phi\) and two triplets \(\Delta_L\) and \(\Delta_R\) with vacuum expectation values (VEV) \(v_L\) and \(v_R\), respectively. The VEVs fulfill the condition \(v_L v_R = v^2\). The VEV \(v_R\) breaks \(SU(2)_R \otimes U(1)_{B-L}\) to \(U(1)_Y\) and generates masses for the right-handed \(W_R\) and \(Z_R\) gauge bosons, and the heavy neutrinos. The \(W_L\) and \(W_R\) are in general not mass eigenstates and are related to the mass eigenstates \(W_1\) and \(W_2\) with masses \(M_1\) and \(M_2\) \((M_1 < M_2)\) as

\[
\begin{pmatrix}
  W_L^- \\
  W_R^-
\end{pmatrix} = \begin{pmatrix}
  \cos \zeta & \sin \zeta \\
  -\sin \zeta & \cos \zeta
\end{pmatrix}
\begin{pmatrix}
  W_1^- \\
  W_2^-
\end{pmatrix}.
\]

(19)

Then, the effective current-current interaction which can trigger the \(0\nu\beta\beta\)-decay can be written as \([6]\)

\[
H^\beta = \frac{G_\beta}{\sqrt{2}} \left[ j^\rho_L J^\rho_L + \chi j^\rho_L J^\rho_R \\
+ \eta j^\rho_R J^\rho_L + \lambda j^\rho_R J^\rho_R + h.c. \right].
\]

(20)

Here, \(G_\beta = G_F \cos \theta_C\), where \(G_F\) and \(\theta_C\) are Fermi constant and Cabibbo angle, respectively. The coupling constants \(\lambda, \eta\) and \(\chi\) are chosen to be real. We have

\[
\eta \simeq -\tan \zeta, \quad \chi = \eta, \\
\lambda \simeq (M_{W_1}/M_{W_2})^2.
\]

(21)

The left- and right-handed leptonic currents are given by

\[
j_L^\rho = \bar{\nu}_L (1 - \gamma_5) \nu_L, \quad j_R^\rho = \bar{\nu}_R (1 + \gamma_5) \nu_R.
\]

(22)

The \(\nu_{eL}\) and \(\nu_{eR}\) are the weak eigenstate electron neutrinos, which are expressed as superpositions of the light and heavy mass eigenstate Majorana neutrinos \(\nu_j\) and \(N_j\), respectively. The electron neutrinos eigenstates can be expressed as

\[
\nu_{eL} = \sum_{j=1}^{3} (U_{ej} \nu_{jL} + S_{ej} (N_j R)^C) ,
\]

\[
\nu_{eR} = \sum_{j=1}^{3} (T_{ej}^* (\nu_{jL})^C + V_{ej} N_j R).
\]

(23)

The \((3 + 3)\) scenario is assumed. The \(3 \times 3\) block matrices in flavor space \(U, S, T, V\), generalizations of the Pontecorvo-Maki-Nakagawa-Sakata matrix, constitute the \(6 \times 6\) unitary neutrino mixing matrix \([17]\)

\[
U = \begin{pmatrix}
  U & S \\
  T & V
\end{pmatrix},
\]

(24)

which diagonalizes the general \(6 \times 6\) neutrino mass matrix in the basis \((\nu_L, (N_R)^C)^T\):

\[
M = \begin{pmatrix}
  M_L & M_D \\
  M_D^T & M_R
\end{pmatrix}
\]

(25)

with Majorana and Dirac mass terms, which are proportional to \(M_L \approx y_M v_L\), \(M_R \approx y_M v_R\) and \(M_D \approx y_D v\), where \(y_M\) and \(y_D\) are the Yukawa couplings. The full parametrization of matrix \(U\) includes 15 rotational angles and 10 Dirac and 5 Majorana CP violating phases. It is possible to decompose \(U\) as follows \([17]\)

\[
U = \begin{pmatrix}
  1 & 0 & U_0 \\
  0 & A & R \\
  U_0 & R & B
\end{pmatrix},
\]

(26)

where \(0\) and \(1\) are the \(3 \times 3\) zero and identity matrices, respectively. The parametrization of matrices \(A, B, R\) and \(S\) and corresponding orthogonality relations are given in \([17]\). In the limit case \(A = 1\), \(B = 1\), \(R = 0\) and \(S = 0\) there would be a separate mixing of heavy and light neutrinos, which would participate only in left and right-handed currents, respectively. In that case only the neutrino mass mechanism of the \(0\nu\beta\beta\)-decay would be allowed and exchange neutrino momentum dependent mechanisms associated with the \(W_L-W_R\) exchange and \(W_L-W_R\) mixing would be forbidden. If masses of heavy neutrinos are above the TeV scale, the mixing angles responsible for mixing of light and heavy neutrinos are small. By neglecting the mixing between different generations of light and heavy neutrinos \(A, B, R, S\) matrices can be approximated as follows:

\[
A \approx 1, B \approx 1, R \approx \frac{m_D}{m_{LNV}}, S \approx - \frac{m_D}{m_{LNV}}.
\]

(27)

Here, \(m_D\) represents energy scale of charge leptons and \(m_{LNV}\) is the total lepton number violating scale, which corresponds to masses of heavy neutrinos. For sake of simplicity the same mixing angle is assumed for each generation of mixing of light and heavy neutrinos. We see that \(U_0\) can be identified to a good approximation...
with the PMNS matrix and $V_0$ is its analogue for heavy neutrino sector. Since $V_0$ is unknown, it is common to assume that the structure of $V_0$ is the same one as $U_0$.

Assuming the non-relativistic impulse approximation the left and right hadronic currents $J^p_L$ and $J^p_R$ become [6]

\[
J^p_L(x) = \sum_n \tau_n^p \delta(x - r_n) \left[ \left( g_V - g_A C_n \right) g^0 + g^\rho \left( g_A \sigma_n^p - g_V D_n^p - g_P q^p k^p \frac{\sigma_n \cdot q_n}{2 m_N} \right) \right],
\]

\[
J^p_R(x) = \sum_n \tau_n^p \delta(x - r_n) \left[ \left( g'_V - g'_A C_n \right) g^0 + g^\rho \left( -g'_A \sigma_n^p - g'_V D_n^p + g'_P q^p k^p \frac{\sigma_n \cdot q_n}{2 m_N} \right) \right].
\]

(28)

Here, $q_n = p_n - p'_n$ is the momentum transfer between the nucleons. The final proton (initial neutron) possesses energy $E_n' (E_n)$ and momentum $p_n' (p_n)$. $\sigma_n, \tau_n^p$ and $r_n$ are the Pauli matrix, the isospin raising operator and the position operator, respectively. These operators act on the $n$-th nucleon.

The nucleon recoil operators $C_n$ and $D_n$ are given by

\[
C_n = \frac{\vec{\sigma} \cdot (p_n + p'_n)}{2 m_N} - g_P \frac{E_n - E_n'}{2 m_N} \vec{\sigma} \cdot q_n,
\]

\[
D_n = \frac{(p_n + p'_n)}{2 m_N} - i \left( 1 + \frac{g_M}{g_V} \right) \vec{\sigma} \times q_n.
\]

(29)

Here, $m_N$ is the nucleon mass, $q_V = g_V q^2$, $q_M = g_M q^2$, $g_A = g_A q^2$ and $g_P = g_P q^2$ are, respectively, the vector, weak-magnetism, axial-vector and induced pseudoscalar form-factors in the case of left-handed hadronic currents. As the strong and electromagnetic interactions conserve parity there are relations among form-factors entering the left-handed and right-handed hadronic currents [6]:

\[
\frac{g_A}{g_V} = \frac{g'_A}{g'_V}, \quad \frac{g_M}{g_V} = \frac{g'_M}{g'_V}, \quad \frac{g_P}{g_V} = \frac{g'_P}{g'_V}.
\]

(30)

We note that the induced pseudoscalar term of the space component of hadronic currents was not taken into account in derivation of the $0\nu\beta\beta$-decay rate presented in [6]. This simplification is avoided here.

Due to helicity matching of the propagating neutrino the decay amplitude can be divided into two parts:

a) If both vertices are $V - A$ or $V + A$, the amplitude of the process is proportional to the neutrino mass $m_j$. We shall denote the corresponding parts of the $0\nu\beta\beta$-decay amplitude $L-L$ and $R-R$ terms, respectively.

b) If one vertex is $V - A$ and the other is $V + A$, the four momentum of propagating neutrino $q^\mu = (\omega, \vec{q})$ contributes. The corresponding part of the amplitude, which is denoted as $L-R$, is further separated into two terms, the $\omega$-term and the $q$-term.

In the case of $L-L$ and $R-R$-terms the dominant contribution is associated with the emission of electrons in the $s_{1/2}$-wave state [15]. However, the $q$-term changes the parity and therefore it requires that one of the final electrons to be in $s_{1/2}$-wave while the other to be in $p_{1/2}$-wave, or both electrons are in the $s_{1/2}$-wave and nucleon recoil operator is taken into account. Nevertheless, the $q$-term is not negligible since the $\omega$-term is suppressed by a factor $\varepsilon_{12}/q = 1/40$ [6], that makes the $q$-term comparable or even larger in comparison with the $\omega$-term.

The standard approximations of Ref. [6] are adopted: i) Only mechanisms with the exchange of light neutrinos are considered and contributions from heavier neutrinos are neglected. Recently, it was concluded in Ref. [10] that mechanisms with the exchange of light neutrinos could give dominant contributions to the $0\nu\beta\beta$ amplitude in a wide range of the LRSM parameter space.

ii) Closure approximation. Within this approximation energies of intermediate nuclear states $E_n - (E_i + E_f)/2$ are replaced by an average of $E_n - (E_i + E_f)/2 \sim 10$ MeV and the sum over intermediate states is taken by closure, $\sum_n |n\rangle \langle n| = 1$.

iii) The $R-R$-part of the amplitude, that is multiplied by factor $\left| \frac{\Lambda^2 \sum m_j T_{ij}}{m_\beta} \right|$, becomes negligible in comparison with $m_\beta$. Thus it is neglected.

iv) The terms proportional to the square of the nucleon recoil operators are also neglected.

v) For $L-L$-part of amplitude only electrons in the $s_{1/2}$ wave state are included. The inclusion of the $p_{1/2}$ electrons leads only to negligible contribution to the $0\nu\beta\beta$ standard decay rate [15].

vi) In the case of the $L-R$ term, two-nucleon potentials associated with the spatial $q$ and time $\omega$ components of neutrino exchange potentials are simplified as follows:

\[
H_q^L(x) = \int \frac{dq}{2\pi^2} \left( \frac{q^2}{q + \Delta - \varepsilon_{12}} + \frac{q^2}{q + \Delta + \varepsilon_{12}} \right) e^{i\vec{q} \cdot \vec{x}},
\]

\[
\approx \int \frac{dq}{\pi^2} \frac{q^3}{q + \Delta} e^{i\vec{q} \cdot \vec{x}},
\]

\[
H_\omega^L(x) = \int \frac{dq}{2\pi^2} \left( \frac{1}{q + \Delta - \varepsilon_{12}} - \frac{1}{q + \Delta + \varepsilon_{12}} \right) e^{i\vec{q} \cdot \vec{x}}.
\]

(31)

\[
\approx \varepsilon_{12} \int \frac{dq}{\pi^2} \frac{1}{(q + \Delta)^2} e^{i\vec{q} \cdot \vec{x}},
\]

where $\Delta = E_n - (E_i + E_f)/2$ and $\varepsilon_{12} = \varepsilon_1 - \varepsilon_2$. Here $\varepsilon_1$ and $\varepsilon_2$ represent the energies of the final electrons. Furthermore, contribution of the $p_{1/2}$-wave electrons and term in which the nucleon recoil is multiplied by the small $\omega$-term are also neglected.

vii) Since $|\chi U_{ej}^L q_e^L/g_V| \ll |U_{ej}|$, the coupling constant $\chi$ in Hamiltonian [20] is neglected.

viii) A factorization of phase-space factors and nuclear matrix elements is achieved by the approximation in which electron wave functions $g_{\pm 1}(\varepsilon, r), f_{\pm 1}(\varepsilon, r)$ are replaced by their values at the nuclear radius $R$. The no-
$g_{\pm 1}(\epsilon) \equiv g_{\pm 1}(\epsilon, R), \quad f_{\pm 1}(\epsilon) \equiv f_{\pm 1}(\epsilon, R)$ (32)

is used.

Within above approximations the $0\nu\beta\beta$-decay half-life can be written as

$$
\left[ T_{1/2}^{0\nu} \right]^{-1} = \frac{\Gamma^{0\nu}}{\ln 2} = g_A^2 |M_{GT}|^2 \left\{ C_{mm} \left( \frac{|m_{\beta\beta}|}{m_e} \right)^2 + C_{m\lambda} \left( \frac{|m_{\beta\beta}|}{m_e} \right) (\lambda) \cos \psi_1 + C_{m\eta} \left( \frac{|m_{\beta\beta}|}{m_e} \right) (\eta) \cos \psi_2 + C_{\lambda\lambda} (\lambda)^2 + C_{\eta\eta} (\eta)^2 + C_{\lambda\eta} (\lambda) (\eta) \cos (\psi_1 - \psi_2) \right\}.
$$

The effective lepton number violating parameters associated with right-handed currents and their relative phases are given by

$$
\langle \lambda \rangle = \lambda \sum_{j=1}^{3} U_{ej} T_{ej} (g'_v / g_v),
$$

$$
\langle \eta \rangle = \eta \sum_{j=1}^{3} U_{ej} T_{ej},
$$

$$
\psi_1 = \arg \left( \sum_{j=1}^{3} m_j U_{ej} \right) \left( \sum_{j=1}^{3} U_{ej} T_{ej}^* (g'_v / g_v)^* \right),
$$

$$
\psi_2 = \arg \left( \sum_{j=1}^{3} m_j U_{ej} \right) \left( \sum_{j=1}^{3} U_{ej} T_{ej}^* \right).
$$

(33)

The explicit form of nuclear matrix elements $M_{GT}$ and their ratios $\chi_l$ is presented in section III.B. The integrated kinematical factors are defined as

$$
G_{0k} = \frac{G_A^2 m_e^2}{64\pi^5 \ln 2 R^2} \int \delta(\epsilon_1 + \epsilon_2 + M_f - M_i) \times (h_{ok}(\epsilon_1, \epsilon_2, R) \cos \theta + g_{ok}(\epsilon_1, \epsilon_2, R)) \times p_1 p_2 \epsilon_1 \epsilon_2 d\epsilon_1 d\epsilon_2 d(\cos \theta),
$$

(38)

where $k = 1, 2, \ldots, 11$. $p_1$ and $p_2$ are momenta of electrons and $\theta$ is the angle between emitted electrons. The functions $h_{ok}(\epsilon_1, \epsilon_2, R)$ and $g_{ok}(\epsilon_1, \epsilon_2, R)$ are defined in section III.A. These definitions are independent of the weak axial-vector coupling constant $g_A$. The quantities $G_{0k}$ are given in units of inverse years. We note that if the standard wave functions of electron (w.f. A) are assumed $G_{010} = G_{03}$ and $G_{011} = G_{04}$. If in addition contributions from the induced pseudoscalar term of nucleon current are neglected the decay rate in Eq. (33) reduces to that given in [10]. Quantity $G_{0k}^0$ is relevant for the angular correlation between the two electrons. We note that $G_{013}^0 = G_{05}^0 = 0$.

A. Components due to electron wave functions in the phase-space factors

The $s_{1/2}$ and $p_{1/2}$ electron wave functions at nuclear surface associated with emission of both electrons enter into the phase-space factors through the functions presented below.

$M_{WZ} > 2.9$ TeV [11]. In the LRSM there could be additional contributions to $0\nu\beta\beta$-decay due to the double charged Higgs triplet. However, as pointed in Ref. [11], in the considered case of type-I seesaw dominance, these contributions can be neglected.
For phase-space factors $G^{\rho}_{0k}$ related with the angular distribution of emitted electrons the quantities $h_{0k}(\epsilon_1, \epsilon_2, R)$ are:

\[
\begin{align*}
    h_{01} &= -4C_{ss}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2), \\
    h_{02} &= \frac{2\epsilon_2^2}{m_e^2}C_{ss}(\epsilon_1)C_{ss}(\epsilon_2), \\
    h_{03} &= 0, \\
    h_{04} &= -\frac{2}{3m_eR} \left( C_{sp}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) + C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) + C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) + C_{sp}(\epsilon_1)C_{ss}(\epsilon_2) \right), \\
    h_{05} &= \frac{4}{m_eR} \left( C_{sp}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) + C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) + C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) + C_{sp}(\epsilon_1)C_{ss}(\epsilon_2) \right), \\
    h_{06} &= 0, \\
    h_{07} &= -\frac{16}{(m_eR)^2} \left( C_{sp}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) + C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) - C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) - C_{sp}(\epsilon_1)C_{ss}(\epsilon_2) \right), \\
    h_{08} &= -\frac{8}{(m_eR)^2} \left( C_{sp}^{\dagger}(\epsilon_1)C_{sp}(\epsilon_2) + C_{sp}(\epsilon_2)C_{sp}(\epsilon_1) + C_{ss}(\epsilon_1)C_{pp}(\epsilon_2) + C_{ss}(\epsilon_2)C_{pp}(\epsilon_1) \right), \\
    h_{09} &= \frac{32}{(m_eR)^2} \left( C_{ss}(\epsilon_1)C_{ss}(\epsilon_2) \right), \\
    h_{10} &= -\frac{9}{2}h_{101} - h_{02}, \\
    h_{11} &= 9h_{011} + \frac{1}{9}h_{02} + h_{1010}. \quad (39)
\end{align*}
\]

with

\[
\begin{align*}
    \tilde{h}_{101} &= \frac{\epsilon_1\epsilon_2}{3m_e^2R} \left( C_{sp}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) - C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) - C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) - C_{sp}(\epsilon_1)C_{ss}(\epsilon_2) \right), \\
    \tilde{h}_{11} &= -\frac{2}{(3m_eR)^2} \left[ C_{sp}^{\dagger}(\epsilon_1)C_{sp}(\epsilon_2) + C_{sp}(\epsilon_2)C_{sp}(\epsilon_1) + C_{ss}(\epsilon_1)C_{pp}(\epsilon_2) + C_{ss}(\epsilon_2)C_{pp}(\epsilon_1) \right]. \quad (40)
\end{align*}
\]

In addition, the components $g_{0k}(\epsilon_1, \epsilon_2, R)$ of the phase-space factors (38) are:

\[
\begin{align*}
    g_{01} &= g_{11} = C_{ss}^{\dagger}(\epsilon_1)C_{ss}^{\dagger}(\epsilon_2), \\
    g_{02} &= \frac{\epsilon_1^2}{2m_e^2} \left( C_{ss}^{\dagger}(\epsilon_1)C_{ss}^{\dagger}(\epsilon_2) - C_{ss}^{\dagger}(\epsilon_1)C_{ss}^{\dagger}(\epsilon_2) \right), \\
    g_{03} &= \frac{\epsilon_1^2}{m_e} \left( C_{ss}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) - C_{ss}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) \right), \\
    g_{04} &= \frac{1}{3m_eR} \left( -C_{ss}^{\dagger}(\epsilon_1)C_{sp}(\epsilon_2) - C_{ss}(\epsilon_2)C_{sp}(\epsilon_1) + C_{ss}(\epsilon_1)C_{sp}(\epsilon_2) + C_{sp}(\epsilon_2)C_{sp}(\epsilon_1) \right) - g_{03}/9, \\
    g_{05} &= \frac{\epsilon_2^2}{m_eR} \left( C_{ss}^{\dagger}(\epsilon_1)C_{sp}(\epsilon_2) + C_{ss}(\epsilon_2)C_{sp}(\epsilon_1) + C_{ss}(\epsilon_1)C_{sp}(\epsilon_2) + C_{sp}(\epsilon_2)C_{sp}(\epsilon_1) \right), \\
    g_{06} &= \frac{4}{m_eR} \left( C_{ss}^{\dagger}(\epsilon_1)C_{ss}^{\dagger}(\epsilon_2) + C_{ss}^{\dagger}(\epsilon_2)C_{ss}^{\dagger}(\epsilon_1) \right),
\end{align*}
\]

\[
\begin{align*}
    g_{07} &= \frac{-8}{(m_eR)^2} \left( C_{ss}^{\dagger}(\epsilon_1)C_{sp}(\epsilon_2) + C_{ss}(\epsilon_2)C_{sp}(\epsilon_1) \\
    &+ C_{ss}(\epsilon_1)C_{sp}(\epsilon_2) + C_{ss}(\epsilon_2)C_{sp}(\epsilon_1) \right), \\
    g_{08} &= \frac{2}{(m_eR)^2} \left( -C_{sp}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) - C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) \right), \\
    g_{09} &= \frac{8}{(m_eR)^2} \left( C_{ss}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) + C_{ss}(\epsilon_1)C_{ss}(\epsilon_2) \right), \\
    g_{10} &= -\frac{9}{2}g_{010} - g_{02}, \\
    g_{11} &= 9g_{011} + \frac{1}{9}g_{02} + g_{1010}.
\end{align*}
\]

with

\[
\begin{align*}
    g_{010} &= \frac{\epsilon_1\epsilon_2}{3m_e^2R} \left( -C_{ss}^{\dagger}(\epsilon_1)C_{sp}(\epsilon_2) - C_{ss}(\epsilon_2)C_{sp}(\epsilon_1) \right), \\
    g_{011} &= \frac{1}{18m_e^2R^2} \left( C_{sp}^{\dagger}(\epsilon_1)C_{ss}(\epsilon_2) + C_{sp}(\epsilon_2)C_{ss}(\epsilon_1) \right), \\
    g_{012} &= -2C_{sp}(\epsilon_1)C_{sp}(\epsilon_2) + 2C_{sp}(\epsilon_1)C_{sp}(\epsilon_2). \quad (42)
\end{align*}
\]

Here, $C_{ss}$ are combinations of radial components of $s_{1/2}$ and $p_{1/2}$ wave functions,

\[
\begin{align*}
    C_{ss}(\epsilon) &= g_{11}(\epsilon)f_{11}(\epsilon), \\
    C_{sp}(\epsilon) &= g_{11}(\epsilon)f_{11}(\epsilon), \\
    C_{ss}^{\dagger}(\epsilon) &= f_{-1}(\epsilon)f_{11}(\epsilon), \\
    C_{sp}^{\dagger}(\epsilon) &= f_{-1}(\epsilon)f_{11}(\epsilon), \\
    C_{ss}^{\dagger}(\epsilon) &= g_{21}(\epsilon)f_{21}(\epsilon), \\
    C_{sp}^{\dagger}(\epsilon) &= g_{21}(\epsilon)f_{21}(\epsilon), \\
    C_{sp}^{\dagger}(\epsilon) &= g_{11}(\epsilon)f_{11}(\epsilon) \pm g_{21}(\epsilon)f_{21}(\epsilon). \quad (43)
\end{align*}
\]

### B. Nuclear matrix elements entering the decay rate

The expression for the 0νββ-decay half-life in (33) contains matrix element ratios $\chi_I$ and their linear combinations $\chi_{1\pm}$ and $\chi_{2\pm}$. The quantities $\chi_I$ are defined as

\[
\chi_I = M_I/M_{GT},
\]

where $I=F,T, \omega F, \omega GT, \omega T, qF, qGT, qT, R$ and $P$ and $M_{GT}$ is the dominant Gamow-Teller matrix element associated with mechanism due to the left-handed currents. The combinations $\chi_{1\pm}$ and $\chi_{2\pm}$ are given by

\[
\begin{align*}
    \chi_{1\pm} &= \chi_{qF} \pm 3\chi_{qF}, \\
    \chi_{2\pm} &= \chi_{GT} + \chi_{T} \pm \chi_{F} - \frac{1}{9}\chi_{1\pm}. \quad (45)
\end{align*}
\]

The nuclear matrix elements $M_I$ depend on the exchange potentials $h_I(r)$ through

\[
M_{F,G,T} = \sum_{rs} \langle A_f \| h_{F,G,T}(r) \| O_{F,G,T} \| A_i \rangle
\]
where 

\[ M_{\omega F,\omega GT,\omega T} = \sum_{rs} (A_f \langle |h_{\omega F,\omega GT,\omega T}(r-) |\mathcal{O}_{F,GT,T} | A_i \rangle) \]

\[ M_P = \sum_{rs} i \left( A_f \left\langle |h_P(r-) |\mathcal{O}_{F,GT,T} | A_i \rangle \right\rangle \right) \]

\[ M_{qF,qGT,qT} = \sum_{rs} (A_f \langle |h_{qF,qGT,qT}(r-) |\mathcal{O}_{F,GT,T} | A_i \rangle) \]

\[ M_R = \sum_{rs} (A_f \langle |h_{RG}(r-) |\mathcal{O}_{GT} + h_{RT}(r-) |\mathcal{O}_{T} | A_i \rangle) \]

where \( \mathcal{O}_{F,GT,T} \) are the familiar operators 1, \( \vec{\sigma}_1 \cdot \vec{\sigma}_2 \) and 3(\( \vec{\sigma}_1 \cdot \vec{\sigma}_2 \))\( \cdot (\vec{\sigma}_2 \cdot \vec{\sigma}_2 \)).

The two-nucleon exchange potentials \( h_1(r) \) with \( \omega = \omega F, \omega GT \) and \( \omega T \) take the form

\[ h_1(r) = \frac{2R}{\pi} \int f_1(q,r) \frac{q \ dq}{q + E_n - (E_i + E_f)/2} \tag{46} \]

where

\[ f_{GT} = \frac{j_0(q,r)}{g_A^2} \times \left( \frac{g_A^2 q^2}{m_N} - \frac{g_A^2 q^2 g_0^2(q^2)}{3} + \frac{g_A^2 q^2}{4m_N^2} \right) \]

\[ f_{F} = \frac{g_A^2 q^2}{g_A^2} j_0(qr), \]

\[ f_{T} = \frac{j_2(q,r)}{g_A^2} \left( \frac{g_A^2 q^2 g_0^2(q^2)}{2m_N^2} r j_1(qr), \right) \]

\[ f_{qF} = r \frac{g_A^2 q^2}{g_A^2} j_1(qr) q, \]

\[ f_{qGT} = \left( \frac{g_A^2 q^2}{g_A^2} q + 3\frac{g_A^2 q^2}{4m_N^2} \right) j_1(qr), \]

\[ f_{qT} = \frac{r}{3} \left( \left( \frac{g_A^2 q^2}{g_A^2} q - \frac{g_A^2 q^2 g_0^2(q^2)}{2g_A^2} \right) j_1(qr) \right) \]

\[ - \frac{g_A^2 q^2}{2g_A^2} \frac{q^5}{2m_N^2} \left[ 2 j_1(qr)/3 - j_3(qr) \right], \]

\[ f_{RG} = - \frac{R}{3m_N^2} \left( 1 + \frac{g_A^2 q^2}{g_0^2(q^2)} \right) \frac{g_A^2 q^2}{g_A^2} j_0(qr) q^2, \]

\[ f_{RT} = - \frac{R}{6m_N^2} \left( 1 + \frac{g_A^2 q^2}{g_0^2(q^2)} \right) \frac{g_A^2 q^2 g_0^2(q^2)}{g_A^2} j_2(qr) q^2, \]

\[ f_{P} = \frac{R^2}{r} \frac{g_0^2(q^2)}{g_A^2} j_1(qr) q \tag{47} \]

And the two-nucleon exchange potentials \( h_1(r) \) with \( \omega = \omega F, \omega GT \) and \( \omega T \) take the form

\[ h_1(r) = \frac{4R}{\pi} \int f_1(q,r) \frac{q^2 \ dq}{(q + E_n - (E_i + E_f)/2)^2} \tag{48} \]

where

\[ f_{\omega F} = f_F, \quad f_{\omega GT} = f_GT, \quad f_{\omega T} = f_T. \tag{49} \]

Here, \( r_+ = (r_r + r_s)/2, \quad r_- = (r_r - r_s). \quad r_{rs} \) is the coordinate of decaying nucleon and \( j_i(qr) \) \( (i=1,2,3) \) are the spherical Bessel functions. It is assumed that \( p_r + p_s \approx 0, \quad E_r - E_s \approx 0 \) and \( p_r - p_s \approx q, \) where \( q \) is the momentum exchange. The form factors \( g_0^2(q^2), \quad g_A^2(q^2), \quad g_M^2(q^2) \) and \( g_0^2(q^2) \) are defined in [23] and \( g_A = 1.269. \]

If right-handed currents are switched off, all terms in Eq. (33) except that proportional to \( C_1 \) vanish. The connection with the standard \( 0\nu\beta\beta \)-decay formula (33) is then \( G_{01} \equiv G_{0\nu} \) and \( M_{GT}(1 - \chi_F + \chi_T) \equiv M^{0\nu}. \]

**IV. PHASE-SPACE FACTORS WITH IMPROVED ACCURACY**

In Section II different ways of the treatment of radial wave functions \( g_{11}(\varepsilon, r) \) and \( f_{11}(\varepsilon, r) \) associated with emitted electrons in the \( s_1/2 \) and \( p_{1/2} \) wave states were presented. The derivation of the \( 0\nu\beta\beta \)-decay rate was accomplished by considering electron wave functions for the point-like nucleus (wave function B) or an extended one (wave functions A, C and D); that allowed to separate phase-space factors and nuclear matrix elements. The accuracy of the calculation of phase space factors will be discussed next and the improved phase-space factors associated with mechanisms due to right-handed currents obtained using screened exact finite-size Coulomb wave functions of emitted electrons (wave functions D) will be presented.

A numerical computation of all 11 phase-space factors entering the \( 0\nu\beta\beta \)-decay rate was performed by using previously described 4 types of wave functions (A, B, C and D) for a sample of 3 isotopes \( ^{76}\text{Ge}, \quad ^{130}\text{Te} \) and \( ^{150}\text{Nd}. \) Results are presented in Table I. We see that by using the standard treatment of electron wave functions, corresponding to leading finite-size Coulomb corrections (wave functions A), a significant difference with results of the other three approaches appear, especially for nuclei with large nuclear charge \( Z. \) Surprisingly, results obtained with wave functions B, corresponding to an analytical solution of Dirac equations for a pointlike nucleus, better agree with results corresponding to wave functions C and D (exact solution of Dirac equations for a uniform nuclear charge distribution with the radius \( R \)) than those obtained by the standard treatment of wave functions (wave functions A). It indicates that the exact treatment of the Coulomb field plays more important role than the position of decaying nucleon in the nucleus. From the Table I it is apparent that effect of the screening of atomic electrons on the wave functions of emitted electrons does not play an important role.

The phase-factors differ among themselves significantly in magnitude. This fact is manifested in Fig. 2. One can see that phase-factors obtained with standard wave functions (w.f. A) are always larger than those with phase...
from radial wave functions products of phase-space factors.

FIG. 2: The phase factors $G_{0k}$ (k=1, · · ·, 11) in units yr$^{-1}$ for the 0νββ-decay of $^{76}$Ge, $^{130}$Te, and $^{150}$Nd. Calculation was performed by assuming different approximations for the radial wave functions $g_{±1}$ and $f_{±1}$ of an electron: A) The standard approximation of Doi et al. 8; B) An analytical solution of Dirac equations for a pointlike nucleus is assumed; C) An exact solution of Dirac equations for a uniform charge distribution in nucleus is considered; D) The same as the previous case but the electron screening is taken into account 3.

| w.f. | $^{76}$Ge | $^{130}$Te | $^{150}$Nd |
|------|-----------|-----------|-----------|
|      | A | B | C | D | A | B | C | D | A | B | C | D |
| G_{01}.10^{14} | 0.261 | 0.244 | 0.240 | 0.237 | 1.807 | 1.535 | 1.453 | 1.425 | 8.827 | 6.986 | 6.432 | 6.316 |
| G_{02}.10^{14} | 0.428 | 0.404 | 0.397 | 0.391 | 4.683 | 4.064 | 3.851 | 3.761 | 40.190 | 32.401 | 29.869 | 29.187 |
| G_{03}.10^{15} | 1.478 | 1.340 | 1.316 | 1.305 | 12.237 | 9.566 | 9.065 | 8.967 | 70.032 | 49.465 | 45.593 | 45.130 |
| G_{04}.10^{15} | 0.501 | 0.489 | 0.477 | 0.470 | 3.625 | 3.315 | 3.086 | 3.021 | 18.343 | 16.000 | 14.348 | 14.066 |
| G_{05}.10^{13} | 0.791 | 0.727 | 0.572 | 0.566 | 6.390 | 5.185 | 3.842 | 3.790 | 28.537 | 21.183 | 15.061 | 14.873 |
| G_{06}.10^{12} | 0.605 | 0.547 | 0.536 | 0.531 | 3.091 | 2.398 | 2.258 | 2.227 | 11.922 | 8.323 | 7.591 | 7.497 |
| G_{07}.10^{10} | 0.365 | 0.345 | 0.274 | 0.270 | 2.713 | 2.383 | 1.788 | 1.755 | 13.625 | 11.362 | 8.233 | 8.085 |
| G_{08}.10^{11} | 0.245 | 0.236 | 0.151 | 0.149 | 2.877 | 2.653 | 1.579 | 1.549 | 16.833 | 14.996 | 8.564 | 8.405 |
| G_{09}.10^{10} | 1.360 | 1.263 | 1.238 | 1.223 | 6.398 | 5.354 | 5.063 | 4.972 | 27.582 | 21.530 | 19.790 | 19.454 |
| G_{10}.10^{15} | 1.478 | 1.531 | 1.423 | 1.410 | 12.237 | 14.602 | 11.616 | 11.455 | 70.032 | 105.415 | 72.249 | 71.154 |
| G_{11}.10^{15} | 0.501 | 0.500 | 0.484 | 0.476 | 3.625 | 3.564 | 3.220 | 3.148 | 18.343 | 18.334 | 15.376 | 15.055 |

TABLE II: The phase-space factor $G_{01}$ in units yr$^{-1}$ for the 0νββ-decay of $^{76}$Ge, $^{130}$Te, and $^{136}$Xe. Results are presented for exact Dirac wave functions with finite nuclear size and electron screening at nuclear radius R (Exact, approximation scheme D) and those averaged over the distribution deduced from the analysis of the dominant nuclear matrix element (Exact and averaged - see Eqs. (50)).

| w.f. | $^{76}$Ge | $^{130}$Te | $^{136}$Xe |
|------|-----------|-----------|-----------|
|      | A | B | C | D | A | B | C | D |
| G_{01}.10^{14} | Exact | 0.23681 | 1.42547 | 1.46187 |
|      | Exact and averaged | 0.23987 | 1.47396 | 1.52851 |

The distribution function $D(r_{1})$ as follows:

$$g_{±1}(ε,R) = \int_{0}^{∞} g_{±1}(ε,r_{1})D(r_{1})dr_{1},$$

$$f_{±1}(ε,R) = \int_{0}^{∞} f_{±1}(ε,r_{1})D(r_{1})dr_{1}$$

(50)

where in this particular case $D(r_{1}) = δ(r_{1} - R)$.

In Fig. 3 the distribution function $D(r_{1})$ (or equivalently $D(r_{2})$) is shown corresponding to the nuclear matrix element $M^{0ω}$ (associated with the $m_{ββ}$ mechanism) of the 0νββ-decay of $^{76}$Ge and $^{136}$Xe and calculated within quasiparticle random phase approximation with restoration of isospin symmetry 21. We see that β-decay of both nucleons happens mostly in the vicinity of nuclear surface. In Table II the phase space factor $G_{01}$ calculated with help of $D(r_{1}) = δ(r_{1} - R)$ and $D(r_{1})$ deduced from calculated nuclear matrix elements of $^{76}$Ge, $^{130}$Te and $^{136}$Xe are compared. We see that the corresponding effect is very small for $^{76}$Ge and is only about 4-5 % in the case of $^{130}$Te and $^{136}$Xe.
The improved phase-space factors $G_{0j}$ (j=1, · · · , 11) in units yr$^{-1}$ associated with left and right-handed mechanism of the 0νββ-decay for nuclei of experimental interest are presented in Table III. They were obtained using screened exact finite-size Coulomb wave functions for $s_{1/2}$ and $p_{1/2}$ states of electron (wave functions D). In Table IV the phase space factors $G_{0j}$ associated with angular distribution of emitted electrons are presented.

V. CONSTRAINTS ON THE EFFECTIVE TOTAL LEPTON NUMBER VIOLATING PARAMETERS

Experimental 0νββ-decay half-life limits may be used, in combination with the formula (33), to constrain the effective Majorana neutrino mass $m_{ββ}$ and the effective coupling constants ($λ$) and ($η$) of the right-handed currents. This can be done provided the values of phase-space factors and nuclear matrix elements are available. We shall use the quasiparticle random phase approximation (QRPA) [22] and interacting shell-model (ISM) [23] matrix elements for such analysis. Their values are presented in Table V. In the case of ISM the magnitude of matrix elements $M_{GT}$ calculated in [24] are assumed. We note that these matrix elements were evaluated when the contribution from the induced pseudoscalar term of hadron current was not taken into account. In analysis below the case of CP conservation ($ψ_1 = ψ_2 = 0$) is assumed.

Different contributions to the 0νββ-decay rate (33) are associated with different products of effective lepton number violating parameters $m_{ββ}$, $⟨λ⟩$ and $⟨η⟩$, which values are unknown. The importance of these contributions depends also on the value coefficients $C_I$ (I=mm, mλ, mη, λλ, λη and λη), which can be calculated. The quantity is a superposition of contributions $C_I^{0k}$ associated with phase-space factors $G_{0k}$ (k=1, · · · ,11). In Fig. 4 we show ratios $C_I^{0k}/C_I$ for the 0νββ-decay $^{76}$Ge and $^{136}$Xe and both sets of nuclear matrix elements. We note that coefficients $C_{mm}$, $C_{mλ}$, $C_{mη}$, and $C_{mm}$ are dominated by a single contribution associated with a different phase-factor. In the case of $C_{mλ}$ and $C_{λη}$ there is a competition of mostly two contributions.

Using these nuclear matrix elements (Table V) and the phase-space factors calculated here (see Table III) we can deduce from the experimental data $T_{1/2}^{0ν} ≥ 3.0 \times 10^{25}$ for $^{76}$Ge decay [25] and $T_{1/2}^{0ν} ≥ 3.4 \times 10^{25}$ for $^{136}$Xe decay [20] (we use here the combined limit from the EXO and KamLAND-Zen experiments) the constraints on the effective right-handed current couplings ($λ$), ($η$) and the effective Majorana neutrino mass $m_{ββ}$ listed in Table VI. The constraints in Table VI are of a similar magnitude as those in Table I of Ref. [11]. However, they are based now on the exact treatment of the phase-space factors as well as on the more complete account of nuclear matrix elements. Fig. 5 shows the allowed regions for $m_{ββ}$ and $⟨λ⟩$ ($⟨η⟩$) for ($η$) = 0 ($⟨λ⟩ = 0$). Results are presented for the two sets of nuclear matrix elements (ISM [23, 24] and QRPA [22]) and the standard

![FIG. 3: The normalized $r_1$ ($r_2$) dependence of $M^{0ν}$ for $^{76}$Ge and $^{136}$Xe. $r_1$ and $r_2$ are absolute values of a position vector of β-decaying nucleon in a nucleus.](image)

![FIG. 4: The decomposition of coefficients $C_I$ (I=mm, mλ, mη, λλ, λη and λη), which can be calculated. The quantity is a superposition of contributions $C_I^{0k}$ associated with phase-space factors $G_{0k}$ (k=1, · · · ,11). In Fig. 4 we show ratios $C_I^{0k}/C_I$ for the 0νββ-decay $^{76}$Ge and $^{136}$Xe and both sets of nuclear matrix elements. We note that coefficients $C_{mm}$, $C_{mλ}$, $C_{mη}$, and $C_{mm}$ are dominated by a single contribution associated with a different phase-factor. In the case of $C_{mλ}$ and $C_{λη}$ there is a competition of mostly two contributions.

![FIG. 5: The allowed regions for $m_{ββ}$ and $⟨λ⟩$ ($⟨η⟩$) for ($η$) = 0 ($⟨λ⟩ = 0$). Results are presented for the two sets of nuclear matrix elements (ISM [23, 24] and QRPA [22]) and the standard...](image)
TABLE III: Phase-space factors $G_{0j}$ (j=1, · · ·, 11) in units yr$^{-1}$ obtained using screened exact finite-size Coulomb wave functions for $s_{1/2}$ and $p_{1/2}$ states of electron (wave functions D). The Q values were taken from experiment when available, or from tables of recommended value [3]. $G_{01}$ is associated with the mechanism generated just by $m_{3/2}$. In the case of dominance of the (λ) mechanism the decay rate includes phase factors $G_{02}, G_{010}$ and $G_{011}(G_{02}, G_{07}, G_{08}, G_{09}, G_{010}$ and $G_{011}$). The remaining phase factors are due to interference of these mechanisms (see Eq. (35)).

| Isotope  | $^{48}$Ca | $^{76}$Ge | $^{82}$Se | $^{96}$Zr | $^{100}$Mo | $^{106}$Pd | $^{116}$Cd | $^{124}$Sn | $^{130}$Te | $^{136}$Xe | $^{150}$Nd |
|----------|----------|----------|----------|----------|----------|----------|----------|----------|----------|----------|----------|
| $Q_{0j}$ [MeV] | 4.27226 2.03990 2.99512 3.35037 3.03440 2.01785 2.8135 2.28697 2.52697 2.45783 3.37138 |
| $G_{01,10^{14}}$ | 2.483 0.237 1.018 2.062 1.595 0.483 1.673 0.906 1.425 1.462 6.316 |
| $G_{02,10^{14}}$ | 16.229 0.391 3.529 8.959 5.787 0.814 5.349 1.967 3.761 3.679 29.187 |
| $G_{03,10^{15}}$ | 18.907 1.305 6.913 14.777 10.974 2.672 11.128 5.403 8.967 9.047 45.130 |
| $G_{04,10^{15}}$ | 5.327 0.470 2.141 4.429 3.400 0.978 3.569 1.886 3.021 3.099 14.066 |
| $G_{05,10^{15}}$ | 3.007 0.566 2.004 4.120 3.484 1.400 4.060 2.517 3.790 4.015 14.873 |
| $G_{06,10^{12}}$ | 3.984 0.531 1.733 3.043 2.478 0.934 2.563 1.543 2.227 2.275 7.497 |
| $G_{07,10^{15}}$ | 2.682 0.270 1.163 2.459 1.927 0.599 2.062 1.113 1.755 1.812 8.085 |
| $G_{08,10^{11}}$ | 1.109 0.149 0.708 1.755 1.420 0.462 1.703 0.939 1.549 1.657 8.405 |
| $G_{09,10^{10}}$ | 16.246 1.223 4.779 8.619 6.540 1.939 6.243 3.301 4.972 4.956 19.454 |
| $G_{10,10^{14}}$ | 2.116 0.141 0.801 1.855 1.359 0.309 1.418 0.660 1.146 1.165 7.115 |
| $G_{11,10^{15}}$ | 5.376 0.476 2.183 4.557 3.502 1.010 3.704 1.955 3.148 3.238 15.055 |

TABLE IV: Phase space factors $G_{0j}^{θ}$ associated with angular distribution of emitted electrons (see Eq. (38)) in units yr$^{-1}$ obtained using screened exact finite-size Coulomb wave functions for $s_{1/2}$ and $p_{1/2}$ states of electron (wave functions D). The Q-values of Tab. III are assumed.

| Isotope  | $^{48}$Ca | $^{76}$Ge | $^{82}$Se | $^{96}$Zr | $^{100}$Mo | $^{106}$Pd | $^{116}$Cd | $^{124}$Sn | $^{130}$Te | $^{136}$Xe | $^{150}$Nd |
|----------|----------|----------|----------|----------|----------|----------|----------|----------|----------|----------|----------|
| $-G_{01,10^{15}}$ | 8.010 0.679 3.141 6.484 4.951 1.397 5.153 2.699 4.328 4.426 20.101 |
| $-G_{02,10^{14}}$ | 5.144 0.113 1.075 2.769 1.770 0.239 1.629 0.587 1.137 1.111 9.138 |
| $-G_{04,10^{15}}$ | 1.786 0.152 0.703 1.456 1.112 0.314 1.161 0.608 0.977 1.000 4.566 |
| $-G_{05,10^{15}}$ | 10.714 0.910 4.219 8.734 6.674 1.884 6.965 3.650 5.862 6.002 27.397 |
| $G_{07,10^{11}}$ | 8.458 0.711 3.422 7.431 5.706 1.589 6.026 3.089 5.016 5.155 24.824 |
| $G_{08,10^{12}}$ | 3.553 0.402 2.121 5.383 4.271 1.251 5.054 2.651 4.498 4.787 26.100 |
| $G_{09,10^{10}}$ | 5.024 0.313 1.379 2.562 1.904 0.504 1.795 0.899 1.397 1.387 5.889 |
| $G_{10,10^{15}}$ | 0.695 0.028 0.317 1.118 0.764 0.114 0.884 0.334 0.706 0.741 7.816 |
| $-G_{11,10^{15}}$ | 1.790 0.152 0.707 1.466 1.120 0.317 1.172 0.615 0.988 1.012 4.594 |

(w.f. A) and improved (w.f. D) description of electron wave functions. Note that limits on lepton number violating parameters are softened a little when other lepton number violating parameters have non-vanishing values at the same time in comparison with the case when only a single parameter is non-zero. By assuming $\zeta = 0.013$ and 0.0025 mentioned earlier and the current limit $\langle \eta \rangle \leq 2.98 \times 10^{-9}$ ($^{136}$Xe, ISM, w.f. D) we end up with $m_D/m_{LN, LV} = 2.8 \times 10^{-7}$ and $1.5 \times 10^{-6}$, respectively. For $M_{W2} = 2.9$ TeV and $\langle \lambda \rangle \leq 3.34 \times 10^{-7}$ ($^{136}$Xe, ISM, w.f. D) we get $m_D/m_{LN, LV} = 5.0 \times 10^{-6}$. Thus, from the more stringent limits on $\langle \eta \rangle$ we obtain $m_{LN, LV}/TeV = 0.3 - 2 m_D/MeV$, in agreement with the assumption that the basic scale of LRSM is O(TeV). It is therefore obvious that already the present limits of $0\nu\beta\beta$-decay half-lives can be used to constrain meaningfully the allowed parameter space of LRSM, and that mechanism associated with right-handed currents can compete with the one based on $m_{3/2}$ that is so often used.

VI. DIFFERENTIAL DECAY RATES FOR LIMITING CASES

It is of interest to consider the angular correlations of the emitted electrons and the single electron energy spectrum for the three limiting cases of lepton number violating mechanism since with sufficient experimental accuracy one could distinguish between decays due to coupling to the left-handed and right-handed hadronic
TABLE V: Nuclear matrix elements and their ratios. The quasiparticle random phase approximation (QRPA) and interacting shell-model (ISM) matrix elements are from \cite{22} and \cite{23}, respectively. In the case of ISM matrix elements $M_{GT}$ calculated in \cite{24} is used.

|            | $^{76}$Ge  | $^{136}$Xe |
|------------|------------|------------|
| $M_{GT}$   | 2.350      | 1.770      |
| $\chi_F$  | -0.106     | -0.151     |
| $\chi_{1+}$| 0.686      | 0.782      |
| $\chi_{1-}$| 1.340      | 1.784      |
| $\chi_{2+}$| 0.633      | 0.556      |
| $\chi_{2-}$| 0.912      | 0.965      |
| $\chi_R$  | 0.684      | 0.955      |
| $\chi_P$  | -0.544     | 0.256      |

FIG. 5: Limits on the effective neutrino mass $m_{\beta\beta}$ and right handed parameters $\eta$ (left panels, $\langle \lambda \rangle = 0$) and $\lambda$ (right panels, $\langle \eta \rangle = 0$) implied by the constraints on the $0\nu\beta\beta$-decay of $^{76}$Ge (lower panels, $T_{1/2}^{0\nu} \geq 3.0 \times 10^{25}$ \cite{23}) and $^{136}$Xe (upper panels, $T_{1/2}^{0\nu} \geq 3.4 \times 10^{25}$ \cite{23}). To derive the bounds, the values of nuclear matrix elements calculated within the ISM \cite{23} and the QRPA \cite{22} are used. Results are presented for approximate electron wave functions (type A) and exact Dirac wave functions with finite nuclear size and electron screening (type D). Ellipsoids show the boundaries of the allowed domains.

continued...

The differential rate for the $0^+ \rightarrow 0^+ 0\nu\beta\beta$-decay with the energy of one of the emitted electrons $\tilde{\epsilon}_1$ ($\tilde{\epsilon}_1$ is the kinetic energy fraction with respect to $Q_{\beta\beta}$ of one electron, i.e. $\tilde{\epsilon}_1 = \tilde{\epsilon}_1 Q_{\beta\beta} + m_e$ and $\tilde{\epsilon}_2 = Q_{\beta\beta} + 2m_e - \tilde{\epsilon}_1$) and the angular distribution with the angle $\theta$ between the two electrons for three limiting cases can be written as follows:

i) Case $m_{\beta\beta} \neq 0$ ($\langle \lambda \rangle = 0$ and $\langle \eta \rangle = 0$),

$$d\Gamma = g_A^4 |M_{GT}|^2 \left( \frac{|m_{\beta\beta}|}{m_e} \right)^2 dC_{mm}. \tag{51}$$

ii) Case $\langle \lambda \rangle \neq 0$ ($m_{\beta\beta} = 0$ and $\langle \eta \rangle = 0$),

$$d\Gamma = g_A^4 |M_{GT}|^2 \langle \eta \rangle^2 dC_{\eta\eta}. \tag{52}$$

iii) Case $\langle \eta \rangle \neq 0$ ($m_{\beta\beta} = 0$ and $\langle \lambda \rangle = 0$),

$$d\Gamma = g_A^4 |M_{GT}|^2 \langle \eta \rangle^2 dC_{\eta\eta}. \tag{53}$$

where

$$dC_{mm} = (1 - \chi_P)^2 dG_{01},$$

$$dC_{\lambda\lambda} = \chi_2^2 dG_{02} + \frac{1}{9} \chi_{1+}^2 dG_{011} - \frac{2}{9} \chi_{1-} \chi_{2-} dG_{010},$$

$$dC_{\eta\eta} = \chi_2^2 dG_{02} + \frac{1}{9} \chi_{1-}^2 dG_{011} - \frac{2}{9} \chi_{1-} \chi_{2+} dG_{010}.$$
Here, \( a_1/a_0 \) is the energy-dependent angular correlation coefficient, which depends also on the chosen limiting case for lepton number violating parameters.

In Fig. 6 the single electron spectra normalized to the total decay rate are shown as function of the electron energy \( \tilde{\varepsilon} \) for the 0\( \nu \)\( \beta \beta \)-decay of \( ^{76}\text{Ge} \) and \( ^{136}\text{Xe} \) due to non-vanishing \( m_{\beta\beta} \). This quantity, ideally accessible experimentally, depends only very weakly on the chosen isotope, set of calculated nuclear matrix elements, and whether standard or improved description of electron wave functions is used. The different characteristics of these three limiting cases provide a possibility to identify which of the parameters is responsible for 0\( \nu \)\( \beta \beta \)-decay.

Here, \( G_{0k} = \ln 2 \ G_{0k} \).

The differential decay rate can be written as

\[
\frac{d\Gamma}{d\cos\theta d\tilde{\varepsilon}_1} = a_0 \left( 1 + \frac{a_1}{a_0} \cos \theta \right).
\]
It is often assumed that if and when the neutrinoless double beta decay ($0\nu\beta\beta$) is observed, it will be caused by the exchange of a virtual light Majorana neutrino and its decay rate will be proportional to the square of the effective Majorana mass $m_{\beta\beta}$. It would be possible, therefore, to constrain or determine the magnitude of this fundamental parameter based on experimental limits or values of the decay half-life. However, that is not the only possibility. Many other manifestations of the ‘Physics beyond the Standard Model’ that would cause the $0\nu\beta\beta$ decay were considered in the past. Among them the possibility of the right-handed lepton and/or hadron currents that could perhaps compete with the mass mechanism was often discussed, see e.g. [6,7]. Such possibility arises naturally in the left-right symmetric models. In that case the $0\nu\beta\beta$-decay half-life will depend not only on the $m_{\beta\beta}$ but also, perhaps dominantly, on the parameters that characterize the right-handed currents, denoted $\langle \lambda \rangle$ and $\langle \eta \rangle$ here.

When it is assumed that the right-handed currents exist, the $0\nu\beta\beta$-decay half-life can be expressed as a sum of products of the phase-space factors, nuclear matrix elements, and the fundamental parameters that characterize the new physics. In this work the particular emphasis is on the reformulation of this relation, on a careful derivation of all terms in that expression, and on the new and more general evaluation of the phase-space factors. The phase-space factors depend on the wave functions of the emitted electrons, and various approximations were used in the past in their calculation. We use here the exact solutions of the Dirac equation for the $s_{1/2}$ and $p_{1/2}$ electron states, solving the Dirac equation in the potential that includes the nuclear finite size and the electron screening. The possible approximations to this problem are analyzed and discussed, and in particular it is shown that using just the first order expansion in $r$ in order to include the nuclear finite size with the sufficient accuracy, is not enough. Complete table of accurate phase-space factors for nuclei of interest is given. Compared with the treatment that uses only the first terms in the $r$ expansion (denoted as approximation A in this work), the exact phase space factors (approximation D) are smaller, in particular in the heavier nuclei ($^{130}$Te and $^{150}$Nd) the reduction is $\sim 30 \%$ or even more.

It is also often assumed that the nucleons involved in the $0\nu\beta\beta$-decay are essentially at the nuclear surface, hence the phase space factors are evaluated with the electrons placed at $r = R$. The adequacy of that assumption was not tested until now. Here it is shown, see Fig. 3 and Table II that it is a reasonable assumption, even though an increase in the phase-space factors by a few percent in the heavier candidate nuclei is expected.

Having the full set of the phase-space factors, it is possible by combining them with the full set of nuclear matrix elements evaluated elsewhere, to obtain simultaneous or separate limits for the fundamental parameters $m_{\beta\beta}$ and those associated with the right-handed currents $\langle \lambda \rangle$ and $\langle \eta \rangle$. It turns out that again the difference between the previously used approximation A (just the first term in the expansion in $r$) and the more exact treatment (exact Dirac electron wave functions with the nuclear radius $R$ and electron screening) in the final limits is relatively benign in $^{76}$Ge, enlarging the limits on the fundamental parameters only by about 5%. However, in the heavier nucleus $^{136}$Xe the effect is larger, 10-15%.

It is well known that by convincingly determining the $0\nu\beta\beta$ half-life one would obviously show that the total lepton number is not a conserved quantity. However, that determination by itself will be insufficient to decide which of the possible mechanism is responsible for the decay. If, in addition, the single electron energy spectra, and the angular distribution of the emitted electrons, could be detected, it will help substantially in that task. If one of
the possible parameters, $m_{\beta\beta}$, \langle \lambda \rangle, or \langle \eta \rangle dominates, the single particle spectra and angular correlations will be a decisive tool to determine the mechanism. Formulas that determine these quantities and the corresponding phase space factors are shown here. In that case the exact treatment of nuclear size makes only little difference.

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