Neutrino Propagation in Nuclear Medium and Neutrinoless Double-β Decay

S. Kovalenko,[1] M. I. Krivoruchenko,[2,3] and F. Šimkovic[4,5,6]

[1]Universidad Técnica Federico Santa María, Centro-Cientifico-Tecnológico de Valparaiso, Casilla 110-V, Valparaiso, Chile
[2]Institute for Theoretical and Experimental Physics, B. Cheremushkinskaya 25, 117218 Moscow, Russia
[3]Department of Nano-, Bio-, Information, and Cognitive Technologies, Moscow Institute of Physics and Technology, 141700 Dolgoprudny, Russia
[4]Bogoliubov Laboratory of Theoretical Physics, JINR, 141980 Dubna, Russia
[5]Department of Nuclear Physics and Biophysics, Comenius University, Mlynská Dolina F1, SK-842 48 Bratislava, Slovakia
[6]Czech Technical University in Prague, CZ-12800 Prague, Czech Republic

(Received 25 November 2013; revised manuscript received 26 February 2014; published 11 April 2014)

We discuss a novel effect in neutrinoless double-β ($0\nu\beta\beta$) decay related with the fact that its underlying mechanisms take place in the nuclear matter environment. We study the neutrino exchange mechanism and demonstrate the possible impact of nuclear medium via lepton-number-violating (LNV) four-fermion interactions of neutrinos with quarks from a decaying nucleus. The net effect of these interactions is the generation of an effective in-medium Majorana neutrino mass matrix. The enhanced rate of the $0\nu\beta\beta$ decay can lead to the apparent incompatibility of observations of the $0\nu\beta\beta$ decay with the value of the neutrino mass determined or restricted by the β-decay and cosmological data. The effective neutrino masses and mixing are calculated for the complete set of the relevant four-fermion neutrino-quark operators. Using experimental data on the $0\nu\beta\beta$ decay in combination with the β-decay and cosmological data, we evaluate the characteristic scales of these operators: $\Lambda_{\rm LNV} \geq 2.4$ TeV.

DOI: 10.1103/PhysRevLett.112.142503 PACS numbers: 23.40.−s, 14.60.St, 21.65.Jk

Various mechanisms of neutrinoless double-β ($0\nu\beta\beta$) decay have been considered in the literature (for recent reviews see [1,2]). The mechanisms are conventionally constructed as lepton-number-violating (LNV) quark-lepton processes proceeding in a vacuum. Then, after an appropriate hadronization, the presence of the initial and final nuclei is taken into account as a smearing effect via convolution with the corresponding nuclear wave function. On the other hand, the nuclear matter may impact an underlying LNV process in a more direct way via the standard model (SM) or beyond the SM interactions. If this is relevant, an especially notable effect should be expected from the LNV interactions with the nuclear matter. In the present Letter, we consider the Majorana neutrino exchange mechanism and examine the possible impact of nuclear medium via LNV four-fermion neutral current interactions of neutrinos with quarks from a decaying nucleus. The nuclear matter effect on the $0\nu\beta\beta$-decay rate is calculated in the mean field approach. The mean field associated with the strong interaction is created in nuclei by the scalar and vector quark currents and described effectively in terms of the σ and ω mesons [3]. Here, we consider the scalar mean field associated with the LNV interaction. Then, an effective four-fermion neutrino-quark Lagrangian with the operators of the lowest dimension can be written in the form

$$\mathcal{L}_{\rm eff} = \frac{1}{\Lambda_{\rm LNV}^2}\sum_{i,j,q}(g^q_{ij}\overline{\nu^C_{Li}}\nu_{Lj}\cdot\bar{q}q + {\rm H.c.}) \quad (1)$$

$$+\frac{1}{\Lambda^3}\sum_{i,j,q}h^q_{ij}\overline{\nu_{Li}}i\gamma^\mu\overleftrightarrow{\partial}_\mu\nu_{Lj}\cdot\bar{q}q, \quad (2)$$

where the fields ν_{Li} are the active neutrino left-handed flavor states, g^q_{ij} and h^q_{ij} are their dimensionless couplings to the scalar quark currents with $i,\ j = e, \mu, \tau$ satisfying $g^q_{ij} = g^q_{ji}$ and $(h^q_{ij})^* = h^q_{ji}$. The first property follows from the identity $\overline{\nu^C_{Li}}\nu_{Lj} = \overline{\nu^C_{Lj}}\nu_{Li}$, the second one from the Hermiticity of the neutrino operator in the form of kinetic terms. Note that the first term in Eq. (1) violates the lepton number by two units $\Delta L = 2$ while the second one is lepton number conserving $\Delta L = 0$. We neglect all the surface terms, which could, in principle, be nontrivial due to the presence of a nuclear surface where the gradient of the nuclear matter density is large. Thus, we consider a simplified case of the infinite nuclear radius. The scales $\Lambda_{\rm LNV}$ and Λ of the $\Delta L = 2$ and $\Delta L = 0$ operators are, in general, different and are of the order of the masses M of virtual particles inducing these effective operators at tree level. These particles could be either scalars or vectors (vector leptoquarks) with the masses $M \gg p_F \sim 280$ MeV, where p_F is the Fermi momentum of nucleons in nuclei, which sets the momentum scale of $0\nu\beta\beta$ decay. The gauge invariant structure of the operators in Eq. (1) is briefly discussed later.

In the mean field approximation, we replace the operator $\bar{q}q$ in Eq. (1) with its average value $\langle\bar{q}q\rangle$ over the nuclear medium. Relying on the MIT bag model, we have for the light quarks $q = u,\ d$ an estimate $\langle\bar{q}q\rangle \approx \frac{1}{2}\langle q^\dagger q\rangle$ [4], which is equivalent to $\langle\bar{q}q\rangle \approx 0.25\ {\rm fm}^{-3}$ at the saturation. Thus, in the nuclear environment, the Lagrangian (1) is reduced to

$$\mathcal{L}_{\rm eff} = \frac{\langle \bar{q}q \rangle}{\Lambda_{\rm LNV}^2}(\overline{\nu_{Li}^C} g_{ij} \nu_{Lj} + {\rm H.c.}) + \frac{\langle \bar{q}q \rangle}{\Lambda^3} \overline{\nu_{Li}} h_{ij} i\gamma^\mu \overset{\leftrightarrow}{\partial}_\mu \nu_{Lj}, \quad (3)$$

where $g_{ij} = (g_{ij}^u + g_{ij}^d)/2$ and $h_{ij} = (h_{ij}^u + h_{ij}^d)/2$. We assume for simplicity the nuclear medium to be an isosinglet.

Let us recall the terms of the electroweak (EW) Lagrangian in vacuum relevant to the calculation of the amplitude of $0\nu\beta\beta$ decay via the Majorana neutrino exchange mechanism. They are

$$\mathcal{L}_{\rm EW}^{\rm vac} = \frac{1}{4}\overline{\nu_{Li}} i\gamma^\mu \overset{\leftrightarrow}{\partial}_\mu \nu_{Li} - \frac{1}{2}\overline{\nu_{Li}^C}\hat{M}_{ij}^L \nu_{Lj} + \frac{4G_F \cos\theta_C}{\sqrt{2}}\overline{l_{Li}}\gamma^\mu \nu_{Lj} \cdot \bar{u}_L \gamma_\mu d_L + {\rm H.c.}, \quad (4)$$

where $M_{ij}^L = M_{ji}^L$ is a Majorana mass matrix symmetric for the same reason as h_{ij}^q is a matrix in Eqs. (1) and (2). It can be diagonalized by a unitary transformation $\nu_i = U_{ij}^L \nu'_j$. In the basis where the charged lepton mass matrix is diagonal, the unitary matrix U^L coincides with the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) mixing matrix. Thus, in the vacuum, we have

$$\mathcal{L}_{\rm EW}^{\rm vac} = \frac{1}{4}\overline{\nu'_{Li}} i\gamma^\mu \overset{\leftrightarrow}{\partial}_\mu \nu'_{Li} - \frac{1}{2}m_i \overline{\nu'^C_{Li}} \nu'_{Li} + \frac{4G_F \cos\theta_C}{\sqrt{2}}\overline{l_{Li}}\gamma^\mu U_{ij}^L \nu'_{Lj} \cdot \bar{u}_L \gamma_\mu d_L + {\rm H.c.} \quad (5)$$

Here, m_i ($i = 1, 2, 3$) is the neutrino mass in the vacuum.

According to the conventional parametrization $U^L = V^L D$, where V^L is a matrix depending on the three mixing angles and one Dirac phase, $D = {\rm Diag}\{1, \exp(i\alpha_{21}/2), \exp(i\alpha_{31}/2)\}$ is the diagonal matrix of the Majorana phases, which are chosen so that $m_i^* = m_i \geq 0$ and the entry $V_{e3}^L = \sin^2\theta_{13}$ has no Dirac phase.

As is seen from Eq. (3), the neutrino interactions with the nuclear matter affect both the mass and kinetic terms of the vacuum Lagrangian (4), (5) so that the in-medium Lagrangian written in the vacuum mass eigenstate basis takes the form

$$\mathcal{L}_{\rm EW}^{\rm med} = \frac{1}{4}\overline{\nu'_{Li}}\hat{\mathcal{K}}_{ij} i\gamma^\mu \overset{\leftrightarrow}{\partial}_\mu \nu'_{Lj} - \frac{1}{2}\overline{\nu'^C_{Li}}\widehat{\mathcal{M}}_{ij}\nu'_{Lj} + \frac{4G_F \cos\theta_C}{\sqrt{2}}\overline{l_{Li}}\gamma^\mu U_{ij}^L \nu'_{Lj} \cdot \bar{u}_L \gamma_\mu d_L + {\rm H.c.}, \quad (6)$$

where

$$\hat{\mathcal{K}}_{ij} = \delta_{ij} + 4\frac{\langle \bar{q}q \rangle}{\Lambda^3}\hat{h}_{ij}, \qquad \widehat{\mathcal{M}}_{ij} = m_i \delta_{ij} - 2\frac{\langle \bar{q}q \rangle}{\Lambda_{\rm LNV}^2}\hat{g}_{ij}, \quad (7)$$

with $\hat{h} = U^{L\dagger} h U^L$, $\hat{g} = (U^L)^T g U^L$. Thus, we have $\hat{\mathcal{K}}^\dagger = \hat{\mathcal{K}}$, and $\widehat{\mathcal{M}}^T = \widehat{\mathcal{M}}$.

First, we bring the neutrino kinetic term in the Lagrangian (6) to the canonical form. Toward this end, we diagonalize it by a unitary transformation $\nu''_i = V_{ij}\nu''_j$, $V^\dagger \hat{\mathcal{K}} V = {\rm Diag}\{\lambda_k\} \equiv \Omega$, where $\lambda_k^* = \lambda_k \geq 0$. The positiveness of these eigenvalues is maintained as long as $4\langle \bar{q}q \rangle \hat{h} \leq \Lambda^3$, which is implied in our analysis. With this condition, a field rescaling $\nu''_i \to \lambda_i^{-1/2}\nu''_i$ allows us to arrive at the canonical kinetic term

$$\mathcal{L}_{\rm EW}^{\rm med} = \frac{1}{4}\bar{\nu}''_{Li} i\gamma^\mu \overset{\leftrightarrow}{\partial}_\mu \nu''_{Li} - \frac{1}{2}\overline{\nu''^C_{Li}}\lambda_i^{-1/2} V_{ji}\widehat{\mathcal{M}}_{jk} V_{kn}\lambda_j^{-1/2}\nu''_{Ln} + \frac{4G_F \cos\theta_C}{\sqrt{2}}\overline{l_{Li}}\gamma^\mu U_{ij}^L V_{jk}\lambda_k^{-1/2}\nu''_{Lk} \cdot \bar{u}_L \gamma_\mu d_L + {\rm H.c.} \quad (8)$$

Then, we diagonalize the effective Majorana mass term by a unitary transformation $\nu''_i = W_{ij}^L \tilde{\nu}_j$,

$$(W^L)^T(\Omega^{-1/2} V^T \widehat{\mathcal{M}} V \Omega^{-1/2}) W^L = {\rm Diag}\{\bar{\mu}_i\}, \quad (9)$$

where $\bar{\mu}_i = \mu_i \exp(-i\phi_i)$ with $|\bar{\mu}_i| = \mu_i$. These phases can be absorbed by the neutrino fields $\tilde{\nu}_{Li} \to \exp(i\phi_i/2)\tilde{\nu}_{Li}$. Only two of these phases are physical. One of $\phi_{1,2,3}$ can be erased by an overall phase rotation of the charged lepton fields: $l_{Li} \to l_{Li}\exp(-i\phi_1/2)$, where we conventionally selected the phase ϕ_1 to be eliminated. After all that, we finally arrive at the neutrino Lagrangian in the nuclear matter

$$\mathcal{L}_{\rm EW}^{\rm med} = \frac{1}{4}\overline{\tilde{\nu}_{Li}} i\gamma^\mu \overset{\leftrightarrow}{\partial}_\mu \tilde{\nu}_{Li} - \frac{1}{2}\mu_i \overline{\tilde{\nu}^C_{Li}}\tilde{\nu}_{Li} + \frac{4G_F \cos\theta_C}{\sqrt{2}}\overline{l_{Li}}\gamma^\mu U_{ij}^{\rm eff}\tilde{\nu}_{Lj} \cdot \bar{u}_L \gamma_\mu d_L + {\rm H.c.}, \quad (10)$$

in terms of an effective mass eigenstate neutrino field $\tilde{\nu}_{Li}$ in the nuclear environment related to the in-vacuum fields ν_i from Eq. (4) as $\nu_{Li} = U_{ij}^{\rm eff}\tilde{\nu}_{Lj}$ with $U^{\rm eff} = U^L V \Omega^{-1/2} W^L \mathcal{P}$, where $\mathcal{P} = {\rm Diag}\{1, \exp(i\phi_{21}/2), \exp(i\phi_{31}/2)\}$ is the diagonal matter generated Majorana phase matrix, with $\phi_{21} = \phi_2 - \phi_1$, $\phi_{31} = \phi_3 - \phi_1$. Note that the neutrino mixing matrix in medium $U^{\rm eff}$ is not unitary, contrasting to unitarity of the neutrino mixing matrix U^L in vacuum.

The amplitude of $0\nu\beta\beta$ decay for the Majorana neutrino exchange in nuclear medium is proportional to the quantity

$$m_{\beta\beta} = \sum_i (U_{ei}^{\rm eff})^2 \mu_i, \quad (11)$$

which should be compared with the corresponding quantity without nuclear matter effects

$$m_{\beta\beta}^{\rm vac} = \sum_i (U_{ei}^L)^2 m_i. \qquad (12)$$

The experimental searches for $0\nu\beta\beta$ decay provide information on the in-medium effective parameter $m_{\beta\beta}$ from Eq. (11). For various choices of nuclear matrix elements, the currently most stringent limit on this parameter derived by EXO-200 and KamLAND-Zen experiments with ^{136}Xe [5] and by the GERDA experiment with ^{76}Ge [6] is in the range $|m_{\beta\beta}| \leq 0.2$–$0.4$ eV. Discussion of the next-generation experiments aimed at improving the $0\nu\beta\beta$ limits can be found in Ref. [1].

The information on the in-vacuum neutrino masses and mixing is provided by neutrino oscillation experiments (for a review, see Ref. [7]). The quantities measured in these experiments are the neutrino mass squared differences $\Delta m_{ij}^2 = m_i^2 - m_j^2$ and mixing angles θ_{12}, θ_{23}, and θ_{13}. If the overall mass scale is fixed, e.g., by the mass of the lightest neutrino, $m_0 \equiv \min(m_i)$, all the other masses are determined. Two types of the neutrino mass spectra are possible: the normal one with $m_1 < m_2 < m_3$ (NS) and the inverted one with $m_3 < m_1 < m_2$ (IS).

The overall neutrino mass scale in vacuum can be constrained by tritium β decay measurements and cosmological data.

Presently, the best experimental limit on the neutrino parameter m_β observable in tritium β decay is [8] $m_\beta^2 = \sum_i |U_{ei}^L|^2 m_i^2 \leq (2.2 \text{ eV})^2$ at 95% C.L. The KATRIN experiment is expected to improve this limit by a factor of 10 in the near future [9].

Recently, the Planck collaboration [10] reported new limits on the sum of the neutrino masses: $\sum_i m_i \leq 0.23$–1.08 eV, derived from the measurements of the temperature of the cosmic microwave background and lensing-potential power spectra. The lowermost bound implies $m_0 \leq 0.07$ eV. An upper limit of 0.28–0.47 eV for the sum of neutrino masses was reported in Ref. [11].

From the constraints of Refs. [5,6] and [8,10,11], we derive limitations on the four-fermion effective neutrino-quark interactions introduced in Eq. (1). We consider a simplified case for the scalar couplings in Eqs. (1–3) such that $4\hat{h}_{ij}\Lambda^{-3} = \delta_{ij}h$, $2\hat{g}_{ij}\Lambda_{\rm LNV}^{-2} = \delta_{ij}g$, with h, g being real numbers, where $\hat{h}$, $\hat{g}$ are defined after Eq. (7). Then, we have $V_{ij} = \delta_{ij}$, $W_{ij}^L = \delta_{ij}$, $\Omega_{ij} = \delta_{ij}\lambda$, $\lambda = 1 + \langle\bar{q}q\rangle h$, $\mu_i = \lambda^{-1}|m_i - \langle\bar{q}q\rangle g|$. The effective Majorana mass (11) in this case is

$$m_{\beta\beta} = \sum_{i=1}^{n} (V_{ei}^L)^2 \xi_i \frac{|m_i - \langle\bar{q}q\rangle g|}{(1 - \langle\bar{q}q\rangle h)^2}. \qquad (13)$$

Here, V_{ij}^L is the PMNS mixing matrix in vacuum without Majorana phases. The Majorana phase factor is $\xi_i = \{1, \exp(i\alpha_1), \exp(i\alpha_2)\}$ with $\alpha_1 = (\alpha_{21} + \phi_{21})/2$, $\alpha_2 = (\alpha_{31} + \phi_{31})/2$, where α_{ij} are the Majorana phases in vacuum defined together with the matrix V^L after Eq. (5).

Within the simplified scheme, the quantity $m_{\beta\beta}$ in nuclear medium in comparison with the one in vacuum depends on the two new unknown parameters: h, g. In our numerical estimations, we assume that only one of them is different from zero at a time. The unknown phases in Eq. (13) are varied in the interval $[0, 2\pi]$. The vacuum mixing angles and the neutrino mass squared differences are taken from Ref. [7]. We illustrate our results in Fig. 1. The shaded areas display allowed values of $|m_{\beta\beta}|$ and m_0 for a set of sample values of g with $h = 0$. For both NS and IS, these results, being combined with the cosmological and tritium β-decay limits, suggest for the LNV scale

$$\Lambda_{\rm LNV} \geq 2.4 \text{ TeV (Planck)}, \qquad 1.1 \text{ TeV (tritium)}. \qquad (14)$$

With the future KATRIN data, the limit 1.1 TeV in Eq. (14) will be pushed up to ~2 TeV. For convenience, we also give our limits in terms of a dimensionless parameter ε_{ij}

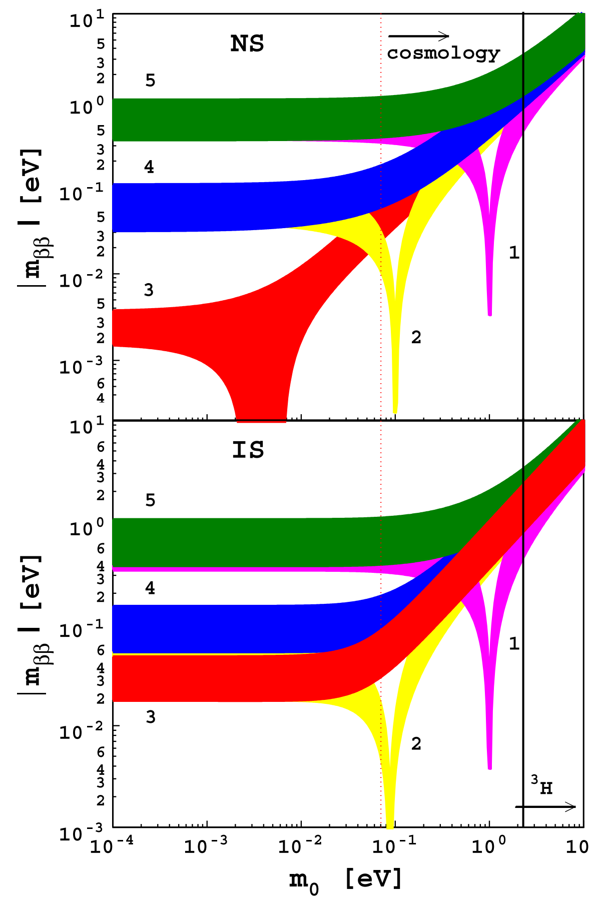

FIG. 1 (color online). The bands 1, 2, 3, 4, and 5 show admissible values of $|m_{\beta\beta}|$ and m_0 for $h = 0$ and $\langle\bar{q}q\rangle g = -1$, -0.1, 0, 0.1, and 1 eV, respectively. The upper and lower panels correspond to the normal (NS) and the inverted (IS) neutrino spectrums. The charge-parity-violating phases spread in the interval $[0, 2\pi]$. Regions to the right from the vertical solid and dotted lines are excluded by the tritium β decay [8] and by the cosmological data [10,11].

defined as $\varepsilon_{ij}G_F/\sqrt{2} = g_{ij}/\Lambda_{\rm LNV}^2$ and characterizing the relative strength of the four-fermion LNV operators in (1) with respect to the Fermi constant G_F. From (14) we have $\varepsilon_{ij} \leq 0.02$ (Planck), 0.1 (Tritium).

The effect of a nonzero value of the coupling constant $h \neq 0$ is particularly simple. Its variation results in shifting the plots in Fig. 1 along the vertical axis. For the case $g = 0$, corresponding to the domains 3 in Fig. 1, the limit $m_{\beta\beta} \leq 0.2$ eV implies very weak constraint, $\Lambda \geq 0.2$ GeV on the scale Λ of the lepton-number-conserving operator in Eq. (1).

Let us briefly comment on the gauge-invariant origin of the operators in Eq. (1). The lepton-number-conserving operator with the derivative stems after the electroweak symmetry breaking (EWSB) from the SM gauge-invariant operators of the type

$$\frac{1}{M^4}[z^d \bar{L} iDL \cdot \bar{Q} d_R \cdot H + z^u \bar{L} iDL \cdot \bar{Q} u_R \cdot \epsilon H^\dagger], \quad (15)$$

where L, Q, and H are the lepton, quark, and Higgs SU_{2L} doublets. The gauge-invariant contractions of their components are implied and involve the SM gauge-covariant derivative $D = \gamma_\mu D^\mu$. The LNV operators in Eq. (1) may have various origins. Some examples are

$$\frac{1}{M_{\rm LNV}^3}[\kappa_1 \overline{L_\alpha^C} L_\beta \cdot \overline{Q}_\alpha u_R \cdot \epsilon_{\beta\gamma} H_\gamma + \kappa_2 \overline{L_\alpha^C} \gamma_\mu u_R \cdot \bar{Q}_\alpha \gamma^\mu L_\beta \cdot \epsilon_{\beta\gamma} H_\gamma + \cdots]. \quad (16)$$

Here, the subscript Greek letters denote components of the SU_{2L} doublets. A complete list of the corresponding operators and their possible ultraviolet completions will be presented elsewhere. In Eqs. (15) and (16), we introduced common scales M, $M_{\rm LNV}$ of the operators and their dimensionless couplings $z^{u,d}$, κ_i, which are, in general, nondiagonal matrices in the flavor space. After the EWSB, due to $\langle H^0 \rangle = v$, these operators engender the corresponding operators in Eqs. (1) and (2) with the scales

$$\frac{z^q v}{M^4} = \frac{h^q}{\Lambda^3}, \qquad \frac{\kappa v}{M_{\rm LNV}^3} = \frac{g}{\Lambda_{\rm LNV}^2}. \quad (17)$$

As seen from Eq. (16), due to the gauge invariance, the terms with the scalar quark currents $\bar{q}q$ in Eq. (1), appearing after the EWSB, have to be accompanied with the pseudoscalar ones $\bar{q}\gamma_5 q$ having the same couplings so that

$$\frac{1}{\Lambda_{\rm LNV}^2} \overline{\nu_L^C} \nu_L [(g^u \bar{u}u + g^d \bar{d}d) + (g^u \bar{u}\gamma_5 u - g^d \bar{d}\gamma_5 d)].$$

Therefore, the scale $\Lambda_{\rm LNV}$ can also be evaluated from $\mathrm{BR}(\pi^0 \to \nu\nu) \leq 2.7 \times 10^{-7}$ [7]. Assuming $g_{ij} = 1$ as in Eq. (14), we have $\Lambda_{\rm LNV} \geq 560$ GeV, which is less stringent than those in Eq. (14).

Note that the scale (14) of the operators in Eq. (2) suggests underlying renormalizable mechanisms with heavy intermediate particles with masses at the TeV scale which is within the reach of the experiments at the LHC. As shown in Ref. [12], these experiments have great potential in distinguishing the underlying mechanisms and setting limits on the scales of the effective operators.

Nonstandard interactions affect neutrino propagation in matter. Thus, one may expect additional constraints on the energy scale of these interactions from astrophysical implications. The vector four-fermion interactions $\bar{\nu}\gamma\nu \cdot \bar{q}\gamma q$ are intensively discussed in the literature (for a review, see Ref. [13]). Their contribution to the in-medium neutrino Hamiltonian is independent of the neutrino energy in contrast to the scalar-type interactions whose effect reduces to renormalization of the neutrino mass matrix suppressed by the neutrino energy. Neutrino oscillations in matter are, therefore, much less sensitive to the interaction of Eqs. (1) and (2). On the other hand, our constraints (14) are comparable to the most stringent ones derived so far for the nonstandard interactions of the vector type.

Note that the Majorana neutrino mass m_ν in vacuum and the LNV operators in Eq. (1) should originate from the same underlying LNV physics at energy scales above $\Lambda_{\rm LNV}$. However, mechanisms generating these two effective Lagrangian terms may be very different. In this context, it is instructive to estimate the significance of the direct contribution δm_ν of the LNV operators in Eq. (1) to the Majorana neutrino mass. This contribution is given by the quark bubble attached to the neutrino line as it follows from the contraction of the quark fields in Eq. (1). The result is $\delta m_\nu \sim g^q/(4\pi\Lambda_{\rm LNV})^2 m_q^3 \log(\Lambda_{\rm LNV}/m_q)$ where m_q is the light quark $q = u, d$ mass in the loop. The usual $\overline{\rm MS}$ renormalization scheme is applied to obtain the finite result. This relies on the assumption that the complete underlying theory is renormalizable. For $\Lambda_{\rm LNV} \sim 2.4$ TeV, $m_d \sim 5$ MeV, and $g^q = 1$, we find $\delta m_\nu \sim 10^{-6}$ eV, which is very small and could represent only a subdominant contribution to neutrino mass. There must be another mechanism of the neutrino mass generation compatible with the neutrino oscillation data.

In the future, the gradually improving cosmological and single β-decay neutrino mass limits may come into conflict with the possible evidence of $0\nu\beta\beta$ decay. If so, the new physics would be mandatory. In particular it can be represented by the new effective TeV scale neutrino-quark interactions (1), (2) enhanced in $0\nu\beta\beta$ decay by the nuclear mean field. If the dominant mechanism of $0\nu\beta\beta$ decay is Majorana neutrino exchange, the scenario presented here will provide the most direct explanation for the above mentioned possible incompatibility between the experiments.

In conclusion, we revisited the Majorana neutrino exchange mechanism of $0\nu\beta\beta$ decay in the presence of nonstandard LNV interactions of neutrinos with nuclear matter of decaying nuclei. These interactions were parametrized with the effective lepton-number-violating and lepton-number-conserving four-fermion neutrino-quark operators of the lowest dimension. In terms of these operators, we calculated the in-medium Majorana neutrino mass, mixing matrix and the parameter $m_{\beta\beta}$ driving the $0\nu\beta\beta$ decay within the neutrino exchange mechanism. Combining experimental limits on this parameter with the cosmological and tritium β decay constraints on the neutrino overall mass scale we extracted a stringent limit on the scale of the LNV interactions of neutrinos with the quark scalar current. In a similar way, the nuclear matter may affect other underlying mechanisms of $0\nu\beta\beta$ decay.

M. I. K. and F. Š. acknowledge kind hospitality at the Universidad Técnica Federico Santa María, Valparaíso. This work was supported in part by RFBR Grant No. 13-02-01442, the VEGA Grant Agency of the Slovak Republic under Contract No. 1/0876/12, by the Ministry of Education, Youth, and Sports of the Czech Republic under Contract No. LM2011027, and by Fondecyt (Chile) under Grant No. 1100582.

[1] J. D. Vergados, H. Ejiri, and F. Šimkovic, Rep. Prog. Phys. **75**, 106301 (2012).
[2] F. F. Deppisch, M. Hirsch, and H. Pas, J. Phys. G **39**, 124007 (2012).
[3] S. A. Chin and J. D. Walecka, Phys. Lett. **52B**, 24 (1974).
[4] R. L. Jaffe, Phys. Rev. D **21**, 3215 (1980).
[5] M. Auger *et al.* (EXO Collaboration), Phys. Rev. Lett. **109**, 032505 (2012); A. Gando *et al.* (KamLAND-Zen Collaboration), Phys. Rev. Lett. **110**, 062502 (2013).
[6] M. Agostini *et al.* (GERDA Collaboration), Phys. Rev. Lett. **111**, 122503 (2013).
[7] J. Beringer *et al.* (Particle Data Group), Phys. Rev. D **86**, 010001 (2012).
[8] V. N. Aseev *et al.*, Phys. Rev. D **84**, 112003 (2011); Ch. Kraus *et al.*, Eur. Phys. J. C **40**, 447 (2005).
[9] G. Drexlin, V. Hannen, S. Mertens, and C. Weinheimer, Adv. High Energy Phys. **2013**, 293986 (2013).
[10] P. A. R. Ade *et al.* (Planck Collaboration), arXiv:1303.5076.
[11] S. A. Thomas, F. B. Abdalla, and O. Lahav, Phys. Rev. Lett. **105**, 031301 (2010).
[12] J. C. Helo, S. G. Kovalenko, M. Hirsch, and H. Pas, Phys. Rev. D **88**, 073011 (2013).
[13] T. Ohlsson, Rep. Prog. Phys. **76**, 044201 (2013).