Neutrinoless double beta decay mediated by the neutrino magnetic moment

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Abstract. We present a new channel of the neutrinoless double beta decay. In this scenario neutrinos not only oscillate inside the nucleus but also interact with an external non-uniform magnetic field. We assume that the field rotates about the direction of motion of the neutrino and show, that for a certain speed of rotation the half-life of the $0\nu\beta\beta$ decay may be significantly lowered.

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1. Introduction

Neutrinos, although weakly, interact with other particles, and therefore propagation and oscillation of these particles in vacuum differs from that in matter. This is known as the Mikheyev – Smirnov – Wolfenstein effect (MSW) [1] and has recently been observed by the Super-Kamiokande Collaboration as an asymmetry in the oscillation rate between zenith and nadir neutrinos [2]. This effect is based on the fact, that the components of ‘ordinary’ matter, i.e., electrons, protons, and neutrons, interact with electron neutrinos via charged as well as neutral currents. Muon and tau neutrinos, on the other hand, cannot interact with the electrons, thus participate in the neutral current processes only. This results in an asymmetry in the forward scattering amplitude of different neutrino flavours, effectively shifting the neutrino oscillation parameters. So regular matter distinguishes between electron and other neutrino flavours.

Weak interactions are not the only factors that may affect neutrino propagation and oscillations. Despite being electrically neutral, neutrinos, according to the Standard Model, should exhibit electromagnetic properties. In the second order 1-loop process in which $\nu \leftrightarrow W^\pm \ell^\mp$ neutrino magnetic moment has been estimated by Fujikawa and Shrock to be $3.2 \times 10^{-19} (m_\nu) \mu_B$ [3], which yields for $m_\nu = 0.05$ eV the value $1.6 \times 10^{-20} \mu_B$. 
The same quantity has been given by Kayser as $\mu \sim 10^{-18} \mu_B$ [5], $\mu_B$ being the Bohr magneton. This value can be higher in various scenarios of physics beyond the Standard Model [6] and for certain ranges of non-standard parameters can reach the experimental limit of roughly $10^{-11} \mu_B$ [7]. Due to the CPT theorem Majorana neutrinos can have only transition (in the flavour basis) magnetic moments, while Dirac neutrinos can have also diagonal magnetic moments. The transition magnetic moments change neutrinos into antineutrinos of opposite helicity and different flavour. Diagonal magnetic moments change neutrinos into antineutrinos of the same flavour. It has also been pointed out [8, 9] that an external non-uniform magnetic field acts differently on neutrinos and antineutrinos, which is due to different helicities of these particles. This observation has been used to show that under special conditions Pontecorvo oscillations $\nu_\alpha \to \bar{\nu}_\alpha$ are possible [11, 12].

The neutrinoless double beta decay (0\textnu 2\beta) is a hypothetical second order process in which some lepton number violating non-standard mechanism accounts for the neutrinos not being released. This process is of the most importance because, if observed, will define neutrinos as Majorana particles, which is the contents of the famous Schechter–Valle black-box theorem [13]. It is also a good laboratory in which various models of non-standard physics and nuclear physics are tested. The simplest and most discussed mechanism of the 0\textnu 2\beta decay is the so-called mass mechanism, in which Majorana neutrinos of non-zero mass are produced in the beta vertex in left-handed states with small right-handed admixture. This R-handed admixture is responsible for the possibility of the neutrino being absorbed in the second beta vertex. Of course the intermediate electron neutrino propagates between the beta vertices as a superposition of three mass eigenstates and therefore the inverse half-life of the decay depends on the so-called effective neutrino mass $\langle m \rangle_{0\nu}$. Assuming same chirality in both beta vertices and light neutrinos, the whole process is described by

$$\sum_{i=1,2,3} U_{ei} \frac{m_i}{p^2 - m_i^2} U_{ei} \approx \frac{1}{p^2} \sum_{i=1,2,3} U_{ei}^2 m_i = \frac{1}{p^2} \langle m \rangle_{0\nu},$$

(1)

where $p$ is the neutrino momentum, $U_{ei}$ are the elements of the first row of the neutrino mixing matrix, and we have used the approximation of small, comparing to $p$, neutrino mass. The factor $1/p^2$ is then absorbed by the nuclear matrix element as a part of the energy denominator. Other mechanisms involve different intermediate particles, like the pions, supersymmetric particles, the Majoron, and others. In fact, the only true condition needed for the 0\textnu 2\beta decay to occur is the one that the electron neutrino somehow changes to its antineutrino. In this paper we propose a new channel of the neutrinoless double beta decay based on the two-step Pontecorvo oscillations. It has been shown [11, 12] that this mechanism has a resonant-like behaviour and its application to the solution of the solar neutrino puzzle has been discussed. In [14] the very same mechanism has been discussed in the context of the 0\textnu 2\beta decay for a simplified two-neutrino case. In this paper we present the realistic three neutrino case together with a numerical analysis to show, that the new mechanism may significantly lower the half-life of the decay in the resonance region.
2. Neutrinos in nuclear matter

Our goal is to describe the nuclear process of the neutrinoless double beta decay. We start therefore with the discussion of the behaviour of neutrinos in the nuclear matter. This problem — neutrino interactions and oscillations inside the nucleus — is omitted in the usual approach to the $0\nu 2\beta$ decay [15], as it is argued that neutrinos travel a very short path between the nucleons. However, as will be shown below, the process of flavour oscillations is vital for the proposed mechanism.

Neutrinos travelling through matter undergo a phase shift due to their interactions with electrons, neutrons and protons via neutral and charged weak currents. In the most typical case of the MSW effect matter is electrically neutral and contains equal amount of electrons and protons, plus some neutrons. The charged current reactions occur between charged leptons and the corresponding neutrinos, so typically electrons and electron neutrinos. Since electrons are absent in the nuclear medium, this interaction is not present inside the nucleus. The neutral current contributions coming from the electrons and protons cancel each other exactly due to the opposite charges of these particles. In the case of nuclear matter, however, there are no electrons and the proton contribution will not be canceled. Let us check, what is the nature of this contribution.

We closely follow the textbook approach presented in [16]. We start by writing the Hamiltonian of neutrinos in vacuum in the mass basis $(\nu_1, \nu_2, \nu_3)^T$

$$H = \text{diag}(E_1, E_2, E_3),$$

which for relativistic neutrinos can be cast in the form

$$H = E + \frac{1}{2E} \text{diag}(m_1^2, m_2^2, m_3^2).$$

The symbol diag represents the diagonal matrix. The neutrino mass eigenstates evolve in time according to the Schrödinger-like equation

$$\frac{1}{i} \frac{d}{dt} \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \end{pmatrix} = H \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \end{pmatrix}. \tag{4}$$

The transformation to the flavour basis $(\nu_e, \nu_\mu, \nu_\tau)^T$ is defined by the unitary Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrix $U$ as

$$H \rightarrow UHU^\dagger = E + \frac{1}{2E} \left[ U \text{diag}(m_1^2, m_2^2, m_3^2) U^\dagger \right]$$

$$= E + \frac{1}{2E} \mathcal{M}^2, \tag{5}$$

where we have denoted by $\mathcal{M}^2$ the square of the neutrino mass matrix in the flavour basis.

The energy levels of the flavour states are corrected in the nuclear matter by their possible interactions via the neutral currents with neutrons and protons [16],

$$V_{nc} = \sqrt{2} G_F \sum_{f=n,p} n_f \left( I_3^{(f)} - 2 q^{(f)} \sin^2 \theta_W \right), \tag{6}$$
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$I_3$ being the third component of the weak isospin, and $q$ the electric charge. Explicitly the proton and neutron contributions read

\begin{align}
V_{\text{nc}}^{(n)} &= \sqrt{2} G_F \left( -\frac{1}{2} \right) n_n, \\
V_{\text{nc}}^{(p)} &= \sqrt{2} G_F \left( \frac{1}{2} - 2 \sin^2 \theta_W \right) n_p,
\end{align}

where $G_F$ is the Fermi constant, $\theta_W$ the Weinberg mixing angle, and $n_{n,p}$ are the neutron and proton number densities. We recall once again that due to the absence of electrons inside the nucleus, the charged-current contribution to the energy of electron neutrino is zero. Therefore all flavour eigenstates are affected by the presence of nuclear matter in the same way. This leads to the conclusion, that in the specific case of nuclear matter the contribution is in the form of a constant shift of the neutrino energy levels. As the neutrino oscillations are sensitive to the differences of masses squared, this will not affect the oscillation rate.

3. Neutrinos in an external magnetic field

As it was already mentioned in the Introduction, neutrinos even within the Standard Model possess a non-zero magnetic moment generated in second-order processes. In the case of Majorana neutrinos this gives a tiny probability of the transition between a neutrino and an antineutrino of different flavour which is triggered by the effective interaction with an external photon.

The main mechanism which leads to the $0\nu2\beta$ decay is the conversion between an electron neutrino and an electron antineutrino. So, if one combines the interaction via the magnetic moment with the flavour oscillations, one gets two possible chains which will satisfy the required conditions ($\alpha = \mu, \tau$):

\begin{align}
\nu_{eL} &\rightarrow \bar{\nu}_\alpha R \rightarrow \bar{\nu}_{eR}, \\
\nu_{eL} &\rightarrow \nu_{\alpha R} \rightarrow \bar{\nu}_{eR}.
\end{align}

Both of these chains lead to the $0\nu2\beta$ decay, but their amplitudes cancel each other exactly. It is due to the fact, that the main part of the amplitudes are the propagators of $\bar{\nu}_\alpha$ and $\nu_\alpha$, respectively. In normal circumstances these particles have the same masses, so the propagators are equal. The opposite signs come from the antisymmetry of the magnetic moment, $\mu_{ea} = -\mu_{ae}$. However, we have assumed, that there is an external magnetic field with which neutrinos interact. Let us examine, how neutrinos behave in these conditions.

Neutrinos and antineutrinos differ in helicities in such a way, that a neutrino has its spin aligned opposite to its momentum, while an antineutrino has $\vec{s}$ aligned in the direction of $\vec{p}$. The presence of the magnetic field has a two-fold effect. Firstly, transitions between neutrino states of different helicities are possible via the neutrino magnetic moment. For Majorana neutrinos, and only those can be discussed in the
context of the neutrinoless double beta decay, the magnetic moment is antisymmetric \( \mu_{\alpha \beta} = -\mu_{\beta \alpha} \), \( \alpha, \beta = \{e, \mu, \tau\} \), i.e., only transitions of the form
\[
\nu_{\alpha L} \rightarrow \bar{\nu}_{\beta R}, \quad \alpha \neq \beta
\]
are possible. The corresponding energy correction will have the form \( B\mu_{\alpha \beta} \). The second effect is, that if the magnetic field changes along the neutrino path, the neutrinos and antineutrinos will obtain different corrections to their effective masses. This follows from the fact that a frame of reference with spin \( \vec{s} \) which is rotating with the angular velocity \( \vec{\omega} \) gains energy \( -\vec{s} \cdot \vec{\omega} \). Since neutrinos \( (s = -1/2) \) and antineutrinos \( (s = +1/2) \) have opposite helicities, the degeneracy of their energy levels is lifted in the presence of an external rotating magnetic field. The component of the rotating field that is parallel to \( \vec{s} \) does not contribute to this effect, so we denote by \( B = |\vec{B}_\perp| \) the magnitude of the perpendicular component of the magnetic field. Let us denote by the angle \( \phi = \phi(t) \) the direction of \( \vec{B}_\perp \) and switch to a reference frame that is rotating together with the field. In such a situation \( \omega = d\phi(t)/dt = \dot{\phi}(t) \). The corrections to the energy coming from the rotation of the magnetic field are thus \( +\dot{\phi}(t)/2 \) for neutrinos, and \( -\dot{\phi}(t)/2 \) for antineutrinos.

The immediate consequence of the lifting of the degeneracy of neutrino and antineutrino masses is, that the amplitudes of the chains (9)-(10) do not cancel each other.

3.1. The two-flavour case

In the simplified case of only two neutrinos \[11,12,13\] the results can be presented in an concise analytical form. They exhibit all the key features of the realistic three-flavour case, which will be described in the next section. Our basis is
\[
(\nu_e, \nu_\mu, \bar{\nu}_e, \bar{\nu}_\mu)^T
\]
and the mixing matrix depends on only one vacuum mixing angle \( \theta \):
\[
U = \begin{pmatrix}
\cos \theta & \sin \theta \\
-\sin \theta & \cos \theta
\end{pmatrix}.
\]
(13)
The Hamiltonian, which is diagonal in the mass basis, becomes non-diagonal in the flavour basis and taking into account matter and magnetic field corrections it takes the block form:
\[
H = \begin{pmatrix}
H_{\nu} & [B\mu] \\
-[B\mu] & H_{\bar{\nu}}
\end{pmatrix},
\]
(14)
where
\[
H_{\nu,\bar{\nu}} = E + \left( V_{nc}^{(n)} + V_{nc}^{(p)} \pm \frac{1}{2} \dot{\phi} \right) + \frac{1}{2E}\mathcal{M},
\]
(15)
with
\[
\mathcal{M} = \begin{pmatrix}
m_1^2 \cos^2 \theta + m_2^2 \sin^2 \theta & \frac{\Delta m^2}{2} \sin 2\theta \\
\frac{\Delta m^2}{2} \sin 2\theta & m_1^2 \sin^2 \theta + m_2^2 \cos^2 \theta
\end{pmatrix}.
\]
(16)
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\[ \Delta m^2 = m_2^2 - m_1^2, \]
\[ V_{\text{nc}}^{(n)} \text{ and } V_{\text{nc}}^{(p)} \text{ being given by (7) and (8), and} \]
\[ [B\mu] = \begin{pmatrix} 0 & B\mu_{e\mu} \\ -B\mu_{e\mu} & 0 \end{pmatrix}. \]

Here \( \mu_{e\mu} \) is the neutrino transition magnetic moment.

The Hamiltonian (14) can be diagonalized and its eigenvalues take the form

\[ E + V_{\text{nc}}^{(n)} + V_{\text{nc}}^{(p)} + \frac{1}{2E} \begin{pmatrix} m_1^2 \\ m_2^2 \\ \bar{m}_1^2 \\ \bar{m}_2^2 \end{pmatrix}, \]

where the overbar indicates masses of the antiparticles. Explicitly the masses squared are given by

\[ m_{1,2}^2 = \frac{1}{2} \left( m_1^2 + m_2^2 \pm \sqrt{(4EB\mu_{e\mu})^2 + (\Delta m^2 + \dot{\phi})^2} \right), \]

\[ \bar{m}_{1,2}^2 = \frac{1}{2} \left( m_1^2 + m_2^2 \pm \sqrt{(4EB\mu_{e\mu})^2 + (\Delta m^2 - \dot{\phi})^2} \right). \]

We notice that the term \( \dot{\phi} \) lifts the degeneracy between the mass eigenstates of neutrinos and antineutrinos, \( m_{1,2}^2 \neq \bar{m}_{1,2}^2 \). Also, in the absence of the magnetic field, \( B = 0, \dot{\phi} = 0 \), we arrive at the expected result \( m_{1,2}^2 = \bar{m}_{1,2}^2 = \bar{m}_{1,2}^2 \).

In the general case, the mixing angles of neutrino mass eigenstates in vacuum and in matter differ, and the source of this difference lies in the interaction of electron neutrinos with electrons via the charged current \( V_{cc} \). The corrected mixing angle is given by

\[ \tan 2\theta' = \frac{\Delta m^2 \sin 2\theta}{\Delta m^2 \sin 2\theta - 2EV_{cc}}, \]

where \( V_{cc} \sim n_e \), the electron number density in matter. In our case, however, \( n_e = 0 \) and therefore the \( U \) matrix remains unchanged.

The nuclear matter effect on neutrino propagation manifests itself as a constant shift of the mass eigenstates. This shift induces a constant phase factor which does not affect the oscillation probabilities, but has an influence on the neutrino propagator.

3.2. The three-flavour case and the neutrinoless double beta decay

The simplified case which assumes only two neutrinos has already been discussed in the literature. In this section we present the generalization to the realistic three-neutrino case.

We expand our flavour basis of neutrinos and antineutrinos to the form

\[ (\nu_e, \nu_\mu, \nu_\tau, \bar{\nu}_e, \bar{\nu}_\mu, \bar{\nu}_\tau)^T. \]

As previously, we take into account the presence of matter and an external non-constant magnetic field. We assume that the field rotates about the direction of motion of the
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We change the reference frame to a co-rotating one \cite{11,12}. In such a situation the full three-flavour effective Hamiltonian takes the form:

\[ H = E + V^{(n)}_{\text{nc}} + V^{(p)}_{\text{nc}} + \frac{\dot{\phi}}{2} \text{diag}(1,-1) + \frac{1}{2} \begin{pmatrix} M^2 & 2[B\mu]_3 \\ -2[B\mu]_3 & M^2 \end{pmatrix}, \]  

(23)

where the 1's are three-dimensional unit matrices, and now

\[ [B\mu]_3 = B \begin{pmatrix} 0 & \mu_{e\mu} & \mu_{e\tau} \\ -\mu_{e\mu} & 0 & \mu_{\mu\tau} \\ -\mu_{e\tau} & -\mu_{\mu\tau} & 0 \end{pmatrix}. \]  

(24)

The three-neutrino mixing matrix which enters \( M^2 \), see (5), is given in the standard trigonometric parametrization by

\[ U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}s_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}s_{13} \end{pmatrix} \times \text{diag}(1,e^{i\phi_2},e^{i\phi_3}), \]  

(25)

where \( s_{ij} \equiv \sin \theta_{ij}, c_{ij} \equiv \cos \theta_{ij} \), and \( \theta_{ij} \) is the mixing angle between the mass eigenstates \( m_i \) and \( m_j \). The \( \delta \) is the CP violating Dirac phase and \( \phi_2, \phi_3 \) are CP violating Majorana phases. In this paper we assume full CP conservation, i.e., we set all the possible phases to zero. This renders the \( U \) matrix real.

For the neutrinoless double beta decay the essential neutrino transition is that of \( \nu_eL \leftrightarrow \bar{\nu}_eR \). The off-diagonal block which describes the mixing of neutrinos with antineutrinos in our setup has the form (24). We see therefore, that a direct transition between \( \nu_eL \) and \( \bar{\nu}_eR \) is forbidden. This transition, however, is possible, as it was pointed out in \cite{11,12}, in a two-step process given by (9) and (10). In the neutrinoless double beta decay the chains of transitions (9)-(10) have to be realized between two beta vertices and may be described by a factor, call it \( \chi \), which is related to the half-life of the \( 0\nu2\beta \) decay by

\[ (T_{1/2}^{0\nu})^{-1} = G^{0\nu}|M^{0\nu}|^2|B\chi|^2, \]  

(26)

where \( G^{0\nu} \) is the exactly computable phase-space factor and \( M^{0\nu} \) the nuclear matrix element obtained in an appropriate (approximate) nuclear model.

Let us be more specific and discuss the case of a neutrinoless double beta \( \beta^- \) decay. The amplitude of this mechanism will be proportional to two propagators of the intermediate neutrino states and two transition probabilities, i.e., the magnetic moment and a product of the \( U \) matrices. We therefore obtain for the analog of the effective neutrino mass the expression

\[ \chi = \sum_{i,j} \sum_{\alpha,\beta} U_{ei} \frac{1}{p - \tilde{m}_i} U_{ai}^* U_{\alpha\beta} U_{\gamma j} \frac{1}{p - \tilde{m}_j} U_{ej}^*, \]  

(27)

where \( i, j = 1, 2, 3 \) number the mass eigenstates and \( \alpha, \beta = e, \mu, \tau \) denote the flavour eigenstates. The masses \( m' \) and \( \tilde{m}' \) can be obtained after diagonalization of the
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Hamiltonian (23). All possible choices of the intermediate flavour states are depicted in figure 1.

We notice first that we cannot use the light neutrino approximation because $m'$ are functions of $\dot{\phi}$ and this parameter can be tuned. If the masses of the neutrino and the corresponding antineutrino are the same, the expression (27) yields zero and this contribution to the $0\nu2\beta$ vanishes. However, the degeneracy is removed by different signs of the $\dot{\phi}$ term in the Hamiltonian (23). What is more, by having the possibility of changing this term, we can arrive at the resonance $p^2 \approx m'^2_\alpha$ or $\bar{p}^2 \approx \bar{m}'^2_\alpha$ boosting the $\chi$ significantly. The typical neutrino momentum $p$ in the $0\nu2\beta$ decay can be assessed from the nuclear radius and is of the order of 100 MeV, which means that for the effective neutrino masses of this order the resonance condition will be met.

Another interesting observation is that in the case of CP-violation non-zero phases in the matrix $U$, c.f. (25), appear and the expression (27) will not be zero even if there is degeneracy among the masses. Similar situation occurs also when there is mixing between the standard model neutrinos and neutrinos from a fourth generation, in which case the matrix $U$ will no longer be unitary.
4. Numerical example

In this section we study the analog of the effective neutrino mass \((27)\) numerically. One of the newest compilation of neutrino oscillation parameters have been given in \([17]\). These are:

\[
\begin{align*}
\sin^2 \theta_{12} &= 0.320, & \sin^2 \theta_{13} &= 0.026, & \sin^2 \theta_{23} &= 0.490, \\
\Delta m^2_{21} &= 7.62 \times 10^{-5} \text{ eV}^2, & \Delta m^2_{31} &= 2.53 \times 10^{-3} \text{ eV}^2, 
\end{align*}
\]

where the normal hierarchy of neutrino masses \((m_1 < m_2 \ll m_3)\) and full CP-conservation have been assumed. The mass of the lightest neutrino has been set arbitrarily to \(m_1 = 0.05 \text{ eV}\). We have estimated the neutrino magnetic moments to be of the order of \(10^{-15} \mu_B\), the field \(B = 1 \text{ T}\), and the neutron and proton number densities to be \(\sim 10^{-31} \text{ eV}^3\), which corresponds to the \(^{76}\text{Ge}\) nucleus.

It is important to note at this point, that the presented above choices are marginal when compared to the average momentum (energy), that neutrino has in the \(0\nu 2\beta\) decay, which is estimated to be roughly \(p \sim E \sim 10^8 \text{ eV}\). It is this value that must be compensated by the \(\dot{\phi}\) parameter in order to reach the resonance region. In the simplified case when due to small mass splitting \((m_1 \approx m_2)\) one discusses only two neutrino oscillations, it is possible to assess analytically the resonance condition, which reads \([14]\)

\[
\dot{\phi} \approx 2E. \tag{29}
\]

In the three-neutrino case the exact formula would be much more involved, but retaining the leading terms only will result in a very similar relation. This is illustrated in figure \([2]\) in which a clear resonance peak in the function \(\chi(\dot{\phi})\) is visible around the value \(2E = 2 \times 10^8 \text{ eV}\). The figure is symmetric around zero, as the field \(B\) may rotate clockwise or anticlockwise. The vertical axis has been rescaled for clearer presentation, therefore no units are given.

One would expect, that the resonance peaks should appear for each mass eigenstate separately. This is indeed the case, but in figure \([2]\) the detailed structure of the peaks

Figure 2. The shape of the function \(\chi(\dot{\phi})\) for the typical neutrino energy \(E = 10^8 \text{ eV}\). A clear resonance peak is visible around the value \(\dot{\phi} \sim 2E\).
is masked by the huge scale difference between the masses and the typical neutrino momentum. On the other hand, the neutrino in a $0\nu2\beta$ decay is a virtual particle represented by an internal line of the Feynman diagram. In principle one should integrate its propagator over the whole momentum space. The structure of the resonance is clearly visible in figure 3, where we have set the energy and the lightest neutrino mass to values comparable with the neutrino mass splitting:

$$E = 5 \times 10^{-5} \text{ eV}, \quad m_1 = 5 \times 10^{-3} \text{ eV}.$$  

One sees clearly that what appeared in figure 2 as a single peak, has a double structure corresponding to $m_{1,2}$ and $m_3$ respectively.

5. Conclusions

We have presented a new channel of the hypothetical neutrinoless double beta decay. In this scenario electron neutrino emitted in one beta vertex reacts through its induced magnetic moment with an external non-uniform field, which results in a helicity-flip and flavour change. A subsequent oscillation back to the electron neutrino flavour allows for an absorption of the resulting antineutrino in the second beta vertex. This process may drive the $0\nu2\beta$ decay.

An interesting feature of this scenario is that the ‘effective mass’ $\chi$ of the propagating neutrino depends on the change of the direction of the external magnetic field, which was represented in our discussion by the parameter $\dot{\phi}$. The function $\chi(\dot{\phi})$ has a pole and for certain values of the argument, which is roughly equal to twice the neutrino energy, becomes large. Since the half-life of the discussed decay is proportional to $\chi^{-1}$ it becomes short in the resonance region. This allows, at least in theory, to induce the $0\nu2\beta$ decay by tuning the external magnetic field. We hope that this scenario may be realized if not now, then in the future experiments.
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