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Overview of neutrino electromagnetic properties

Alexander Studenikin[a,b,c,*]

[a] *Department of Theoretical Physics, Moscow State University, 119992 Moscow, Russia*

[b] *National Centre for Physics and Mathematics, 607328 Satis, Nizhny Novgorod District, Russia*

[c] *Interdisciplinary Scientific and Educational School of Moscow University "Fundamental and Applied Space Research", Moscow State University, 119992 Moscow, Russia*

E-mail: studenik@srd.sinp.msu.ru, a-studenik@yandex.ru

A short introduction to the theory of neutrino electromagnetic properties is presented. Then the main issues related to studies of neutrino electromagnetic characteristics in the scattering experiments are considered. A set of basic neutrino electromagnetic processes that can occur in various astrophysical conditions and be fundamentally observable are discussed. The future prospects that open up in the study of the electromagnetic properties of neutrinos are also discussed, including those provided by the recently obtained results of the XENON1T experiment, new $GEMMA-\nu Gen$ experiment and the work begun on the preparation of a new experiment to detect coherent elastic ν-helium atom scattering with use of a high-intensity tritium antineutrino source in Sarov.

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1. The theory

The most well understood and studied among neutrino electromagnetic characteristics are the neutrino magnetic moments. However, in the Standard Model with massless neutrinos magnetic moments of neutrinos are zero. In a minimal extension of the Standard Model the diagonal magnetic moment of a massive Dirac neutrino is given [1] by

$$\mu_{ii}^D = \frac{3eG_F m_i}{8\sqrt{2}\pi^2} \approx 3.2 \times 10^{-19} \Big(\frac{m_i}{1\ \mathrm{eV}}\Big) \mu_B, \tag{1}$$

μ_B is the Bohr magneton. Therefore, it is believed that the studies of neutrino electromagnetic properties open a window to *new physics* beyond the Standard Model [2]. Note that the Majorana neutrinos can have only transition (off-diagonal) magnetic moments $\mu_{i\neq j}^M$.

It is usually assumed that the entirety of the electromagnetic properties of neutrinos are embodied by the structure of the amplitude corresponding to the Feynman diagramme shown in Fig. 1 that describes the interaction of a neutrino with a real photon. Two incoming and outgoing

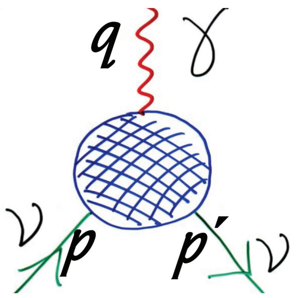

Figure 1: A neutrino effective one-photon coupling.

lines represent initial and final neutrino states and the third line stands for the real photon. These three lines are connected by an effective vertex, which in general contains the whole set of the neutrino electromagnetic characteristics. The diagramme in Fig. 1 corresponds to the one-photon approximation for the electromagnetic interactions of a neutrino field $\nu(x)$ that can be described by the effective interaction Hamiltonian

$$H_{em}(x) = j_\mu(x) A^\nu(x) = \bar{\nu}(x) \Lambda_\mu \nu(x) A^\mu(x), \tag{2}$$

where $j_\mu(x)$ is the neutrino effective electromagnetic current and Λ_μ is a matrix in the spinor space. Considering the neutrinos as free particles with the Fourier expansion of the free Dirac fields, for the amplitude corresponding to the diagramme in Fig. 1 it is possible to get (see [2] and [3] for the detailed derivations)

$$\langle \nu(p_f) | j_\mu(0) | \nu(p_i) \rangle = \bar{u}(p_f) \Lambda_\mu(q) u(p_i), \tag{3}$$

where $q = p_i - p_f$.

In the most general form the neutrino electromagnetic vertex function $\Lambda_\mu^{ij}(q)$ can be expressed [2] in terms of four form factors

$$\Lambda_\mu^{ij}(q) = \left(\gamma_\mu - q_\mu \not{q}/q^2\right) \left[f_Q^{ij}(q^2) + f_A^{ij}(q^2) q^2 \gamma_5 \right] - i\sigma_{\mu\nu} q^\nu \left[f_M^{ij}(q^2) + i f_E^{ij}(q^2) \gamma_5 \right], \tag{4}$$

PoS(CORFU2021)057

where $\Lambda_\mu(q)$ and form factors $f_{Q,A,M,E}(q^2)$ are 3×3 matrices in the space of massive neutrinos. Note that in the derivation of the decomposition (4) the demands followed from the Lorentz-invariance and electromagnetic gauge invariance are taken into account.

In the case of coupling with a real photon ($q^2 = 0$) the form factors $f(q^2)$ provide four sets of neutrino electromagnetic characteristics: 1) the electric millicharges $q_{ij} = f_Q^{ij}(0)$, 2) the dipole magnetic moments $\mu_{ij} = f_M^{ij}(0)$, 3) the dipole electric moments $\epsilon_{ij} = f_E^{ij}(0)$ and 4) the anapole moments $a_{ij} = f_A^{ij}(0)$.

The expression (4) for $\Lambda_\mu^{ij}(q)$ is applicable for Dirac and Majorana neutrinos. However, a Majorana neutrino does not have diagonal electric charge and dipole magnetic and electric form factors, only a diagonal anapole form factor can be nonzero. At the same time, a Majorana neutrino can also have nonzero off-diagonal (transition) form factors.

If one considers the case of equal masses for the initial and final neutrinos, $m_i = m_f$, then the decomposition of the neutrino electromagnetic vertex function reduces to

$$\Lambda_\mu^{ii}(q) = f_Q^{ii}(q^2)\gamma_\mu - i\sigma_{\mu\nu}q^\nu \left[f_M^{ii}(q^2) + i f_E^{ii}(q^2)\gamma_5\right] + f_A^{ii}(q^2)q^2 \left(\gamma_\mu - q_\mu \not{q}/q^2\right)\gamma_5. \tag{5}$$

From the demand that the form factors at zero momentum transfer, $q^2 = 0$, are elements of the scattering matrix, it follows that in any consistent theoretical model the form factors in the matrix element (3) should be gauge independent and finite. Then, the form factors values at $q^2 = 0$ determine the static electromagnetic properties of the neutrino that can be probed or measured in the direct interaction with external electromagnetic fields. This is the case for charge, dipole magnetic and electric neutrino form factors in the minimally extended Standard Model [4, 5].

In non-Abelian gauge theories, the form factors in the matrix element (3) at nonzero momentum transfer, $q^2 \neq 0$, can be non-invariant under gauge transformations. This happens because in general the off-shell photon propagator is gauge dependent. Therefore, the one-photon approximation is not enough to get physical quantities. In this case the form factors in the matrix element (3) cannot be directly measured in an experiment with an external electromagnetic field. However, they can contribute to higher-order diagrams describing some processes that are accessible for an experimental observation (see [6]).

It is also interesting to consider neutrino electromagnetic properties for the case of massless neutrinos of the Standard Model when neutrinos are described by two-component left-handed Weyl spinors. In this case neutrinos have only one form factor which is equal to the difference of the charge form factoer $f_Q(q^2)$ and anapole $f_A(q^2)$ form factor multiplied by q^2, and the electromagnetic vertex function is given by

$$\Lambda_\mu(q) = (\gamma_\mu - q_\mu \not{q}/q^2) f(q^2), \tag{6}$$

where

$$f(q^2) = f_Q(q^2) - f_A(q^2)q^2. \tag{7}$$

From these expressions one can expect that at least from the phenomenological point of view it is not possible to treat consequences of neutrino nonzero electric charge $f_Q(q^2)$ and anapole $f_A(q^2)$ form factors separately. This case can approximate the relation between these two form factors that arises for ultrarelativistic massive neutrinos. For a detailed discussion of the correspondence between the electric charge $f_Q(q^2)$ and anapole $f_A(q^2)$ form factors see in [2].

In the case of a massive neutrino, there is no simple relation between electric charge and anapole form factors since the $q_\mu \not{q}\gamma_5$ term in the anapole part of the vertex function (5) cannot be neglected. Moreover, a direct one-loop calculation of the massive neutrino electromagnetic vertex function (see [4, 5] and [7]), reveals that each of the corresponding Feynman diagrams gives nonzero contribution to the term proportional to $\gamma_\mu\gamma_5$. These contributions are not vanishing even at $q^2 = 0$. Therefore, in addition to the usual four terms in (5) an extra term proportional to $\gamma_\mu\gamma_5$ appears and a corresponding additional form factor $f_5(q^2)$ must be introduced. This problem is related to the decomposition of the massive neutrino electromagnetic vertex function. The calculation [4, 5] of the contributions of the proper vertex diagrams and $\gamma - Z$ self-energy diagrams for an arbitrary gauge fixing parameter $\alpha = 1/\xi$ in the general R_ξ gauge and arbitrary mass parameter $a = m_l^2/m_W^2$ shows that at least in the zeroth and first orders of the expansion over the small neutrino mass parameter $b = (m_\nu/m_W)^2$ the corresponding "charge" $f_5(q^2 = 0)$ is zero. The cancellation of contributions from the proper vertex and self-energy diagrams to the form factor $f_5(q^2)$ at $q^2 \neq 0$,

$$f_5(q^2) = f_5^{(\gamma-Z)}(q^2) + f_5^{(\text{prop.vert.})}(q^2) = 0, \tag{8}$$

was also shown [4, 5] for arbitrary mass parameters a and b in the 't Hooft-Feynman gauge $\alpha = 1$.

Hence, in the minimally extended Standard Model one can perform a direct calculation of the neutrino vertex function leading to the four terms in (4) with gauge-invariant electric charge, magnetic, electric and anapole moments.

Taking into account the above-mentioned feature of the anapole form factor $f_A(q^2)$ (the ambiguity of its allocation against the charge form factor $f_Q(q^2)$ which manifests itself even in the Standard Model), an alternative decomposition of the electromagnetic vertex function $\Lambda_\mu(q)$ has been proposed in [8]. In [8] the toroidal dipole form factor $f_T(q^2)$ is introduced as a characteristic of the neutrino instead of the anapole form factor $f_A(q^2)$. In this case, the neutrino electromagnetic vertex can be written in the so-called toroidal parametrization

$$\Lambda_\mu^{ij}(q) = f_Q^{ij}(q^2)\gamma_\mu - i\sigma_{\mu\nu}q^\nu \left[f_M^{ij}(q^2) + i f_E^{ij}(q^2)\gamma_5 \right] + i f_T^{ij}(q^2)\epsilon_{\mu\nu\lambda\rho}P^\nu q^\lambda \gamma^\rho, \tag{9}$$

where $P = p_i + p_f$ and $\epsilon_{\mu\nu\lambda\rho}$ is the Levi-Civita unit antisymmetric tensor. The toroidal parametrisation of the neutrino vertex has a more clear physical interpretation than the anapole one, because it provides a one-to-one correspondence between the form factors and the multipole moments in expansion of electromagnetic fields. The corresponding toroidal dipole moment $f_T(q^2 = 0)$ provides an alternative for the description of T-invariant interactions with nonconcervation of the P and C symmetries to the anapole moment $f_A(q^2 = 0)$, for the first time introduced in [9].

From the identity

$$\bar{u}(p_f)\left[(m_i - m_f)\sigma_{\mu\nu}g^\mu + (g^2\gamma_\mu - \not{q}q_\mu) - \epsilon_{\mu\nu\lambda\rho}P^\nu q^\lambda\gamma^\rho\gamma_5\right]\gamma_5 u_i(p_i) = 0 \tag{10}$$

it follows [8] that in the static limit the toroidal $f_T(q^2 = 0)$ and anapole $f_A(q^2 = 0)$ dipole moments coincide when the masses of the initial and final neutrino states are equal to each other, $m_i = m_f$.

2. Neutrino electromagnetic properties in scattering experiments

The possible electromagnetic characteristics of neutrinos can manifest themselves in astrophysical environments, where neutrinos propagate in strong magnetic fields and dense matter, and also in ground-based laboratory measurements of neutrinos fluxes from various sources. So far, there are no any evidenced in favour of neutrino nonzero electromagnetic properties either from laboratory measurements of neutrinos from ground-based sources or from observations of neutrinos from astrophysical sources [10, 11]. Only constraints (the upper bounds on neutrino electromagnetic characteristics) are obtained in different experiments. The available constraints are discussed in the review paper [2] (see also [11–15] for the latest progress and achievements in this field).

A widely used method to probe the neutrino electromagnetic properties is based on the direct measurements of low-energy elastic (anti)neutrino-electron scattering in reactor, accelerator, and solar neutrino experiments. Possible nonzero electromagnetic characteristics of neutrinos, such as the electric millicharges, charge radii, dipole magnetic and electric moments, and anapole (toroidal) moments can provide additional contributions to the scattering of neutrinos on a target which are measured in the corresponding experiments.

A general strategy of such experiments consists in determining deviations of the scattering cross section differential with respect to the energy transfer from the value predicted by the Standard Model of electroweak interaction. The desired goal of such experiments can be, firstly, to obtain information about the magnitude of the contributions of electromagnetic interactions of neutrinos to the cross sections. Then this information (a modification of the scattering cross-section confirmed in a specific experiment due to electromagnetic interactions of neutrinos) can be used to study the patterns of neutrino propagation in various media or to analyze data from various other experiments. Within the framework of this approach, we are not interested in the nature of the occurrence of nontrivial electromagnetic properties of neutrinos. In this case, only restrictions on the numerical values of specific electromagnetic characteristics of neutrinos extracted from the data of the neutrino scattering experiment are used. The question of the origin of the nonzero electromagnetic properties of neutrinos, that is, about a more fundamentally theoretical model predicting the nonzero properties of neutrinos, remains behind the scenes.

The second main purpose of neutrino scattering experiments can be to compare the obtained experimental data on the cross section (more precisely, from the obtained upper bounds of the possible electromagnetic contribution to the cross section) with the predictions of a more general theoretical model that provides nonzero electromagnetic characteristics of neutrinos from the first fundamental principles. In this context, we can say that the study of the electromagnetic properties of neutrinos opens a window into new physics.

The most consistent approach to the theoretical description of the electromagnetic properties of neutrinos involves the initial introduction of nonzero electromagnetic characteristics for the mass states of neutrinos. At the same time, since in neutrino scattering experiments flavour neutrinos are registered in detectors, it is necessary to translate electromagnetic properties from the mass into the flavour neutrino basis. Therefore, due to the neutrino mixing and oscillations along the neutrino path from the source to the detector the observed (constrained) neutrino electromagnetic characteristics depends on the neutrino flavour composition in the detector. The recent and most comprehensive study of neutrino electromagnetic properties in the neutrino electron scattering with account for neutrino mixing and oscillations can be found in [16].

Consider the most stringent constraints on the effective neutrino magnetic moments that are obtained with the reactor antineutrinos: $\mu_\nu \leq 2.9 \times 10^{-11} \mu_B$ (GEMMA Collaboration [17]), and solar neutrinos: $\mu_\nu \leq 2.8 \times 10^{-11} \mu_B$ (Borexino Collaboration [18]). Both these constraints are obtained with investigations of the elastic scattering of a flavour neutrino ν_l (or an antineutrino $\bar{\nu}_l$) on an electron at rest: $\nu_l + e^- \to \nu_l + e^-, \quad l = e, \mu, \tau$. There are two contributions, one from the Standard Model weak interaction and another one from the neutrino magnetic moment interaction, to the electron neutrino cross section [19],

$$\frac{d\sigma_{\nu_l e^-}}{dT_e} = \left(\frac{d\sigma_{\nu_l e^-}}{dT_e}\right)_{SM} + \left(\frac{d\sigma_{\nu_l e^-}}{dT_e}\right)_{\mu}. \tag{11}$$

The weak-interaction cross section is

$$\left(\frac{d\sigma_{\nu_l e^-}}{dT_e}\right)_{SM} = \frac{G_F^2 m_e}{2\pi}\left\{(g_V^{\nu_l} + g_A^{\nu_l})^2 + (g_V^{\nu_l} - g_A^{\nu_l})^2\left(1 - \frac{T_e}{E_\nu}\right)^2 + \left[(g_A^{\nu_l})^2 - (g_V^{\nu_l})^2\right]\frac{m_e T_e}{E_\nu^2}\right\}, \tag{12}$$

with the standard coupling constants g_V and g_A given by: $g_V^{\nu_e} = 2\sin^2\theta_W + 1/2, g_A^{\nu_e} = 1/2, \quad g_V^{\nu_{\mu,\tau}} = 2\sin^2\theta_W - 1/2, \quad g_A^{\nu_{\mu,\tau}} = -1/2$. The cross section depends on the initial neutrino energy E_ν and also contains the electron recoil energy T_e. For antineutrinos one must substitute $g_A \to -g_A$. If to account that the born in the source flavour neutrino $|\nu_l\rangle$ arrives to the detector in the flavour state given by

$$|\nu_l(L)\rangle = \sum_{k=1}^{3} U_{lk}^* e^{-\frac{m_k^2}{2E_\nu}L}|\nu_k\rangle, \tag{13}$$

then the neutrino magnetic-moment contribution to the cross section is given by

$$\left(\frac{d\sigma_{\nu_l e^-}}{dT_e}\right)_{\mu} = \frac{\pi\alpha^2}{m_e^2}\left(\frac{1}{T_e} - \frac{1}{E_\nu}\right)\left(\frac{\mu_{\nu_l}}{\mu_B}\right)^2, \tag{14}$$

here μ_B is the Bohr magneton. The cross section contains the effective magnetic moment μ_{ν_l} [2], [16]

$$\mu_{\nu_l}{}^2(L, E_\nu) = \sum_j \left|\sum_k U_{lk}^* e^{-i\Delta m_{kj}^2 L/2E_\nu}\left(\mu_{jk} - i\epsilon_{jk}\right)\right|^2 \tag{15}$$

that indeed receives equal contributions from the neutrino electric and magnetic dipole moments, both the diagonal and transition, which are given by the static values of the corresponding form factors, $\mu_{jk} = f_M^{jk}(q^2 = 0), \quad \epsilon_{jk} = f_E^{jk}(q^2 = 0)$.

It is just straightforward that scattering experiments based on detection of neutrinos arriving from different distances probe (or constrain) different combinations of the fundamental magnetic moments. Now for simplicity we omit possible contributions from the dipole electric moments [16]. In the reactor short-baseline experiments (for instance, in the GEMMA experiment) one studies (and constrains) the effective magnetic moments in the flavor basis. Whereas in the case of long-baseline experiments (such as the Borexino experiment with the solar neutrinos) a more convenient interpreting the results is based on effective magnetic moments in the fundamental mass basis, rather than in the flavour basis. These should be accounted for when one compares quite similar values for effective magnetic moments obtained form the GEMMA and Borexino experiments.

Note that with the inclusion [12] of the effect from possible nonzero neutrino millicharge to the analysis of the GEMMA collaboration data on the antineutrino-electron scattering provides the most severe reactor upper bound on the neutrino millicharge: $q_\nu \leq 1.5 \times 10^{-12} e_0$.

3. Neutrino electromagnetic processes

Neutrinos with nonzero electromagnetic characteristics, due to their couplings with real and virtual photons, generate processes that can occur in various astrophysical conditions and be the cause of important phenomena that are fundamentally observable. The most important neutrino electromagnetic processes are shown in Fig. 2 (see also [2] and [10]). They are the following: 1) a heavier neutrino decay to a lighter mass state in vacuum, 2) the Cherenkov radiation by a neutrino in matter or an external magnetic field, 3) the spin-light of neutrino in matter, 4) the plasmon decay to a neutrino-antineutrino pair in matter, 5) the neutrino scattering on an electron or a nuclei, and 6) the neutrino spin precession in an external magnetic field or the transversally moving (or the transversally polarized) matter. All of these processes can be of great interest in astrophysics, and registration of possible consequences of these processes in experiments allows us to obtain information about the values of the electromagnetic characteristics of neutrinos and also set appropriate limits. So, indeed, astrophysics can be considered as a laboratory for studying the electromagnetic properties of neutrinos (see [2], [10] and [11]).

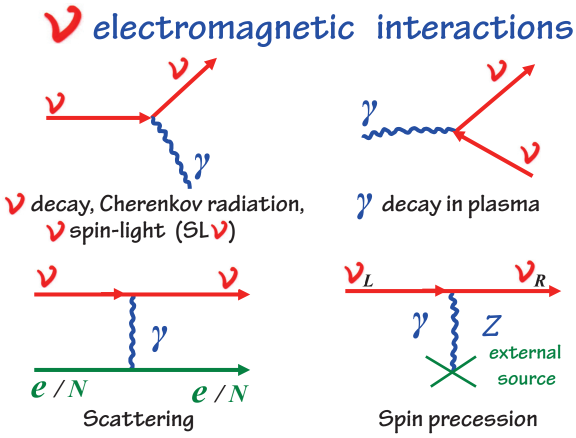

Figure 2: Basic neutrino electromagnetic processes

4. Neutrino radiative decay

A heavier neutrino mass state ν_i in case the neutrino have nonzero electric charges (millicharges) or the magnetic and electric (the diagonal and transition) dipole moments can decay into a lighter state ν_f, $m_i > m_f$, with emission of a photon. For the first time this kind of processes were discussed in [20] and in the concrete applications to neutrinos can be found in [21] and [22]. For more recent papers and detailed discussions of neutrino radiative decay $\nu_i \to \nu_f + \gamma$ see [2]. In case one neglects the effect of nonzero neutrino millicharges ($q_\nu = 0$) the neutrino electromagnetic vertex function reduces to

$$\Lambda_\mu^{if} = -i\sigma_{\mu\nu}q^\nu \left(\mu_{if} + i\epsilon_{if}\gamma_5\right), \tag{16}$$

and the corresponding neutrino interaction Lagrangian reads

$$L = \frac{1}{2}\bar{\psi}_i \sigma_{\mu\nu}(\mu_{ij} + i\gamma_5\epsilon_{ij})\psi_j F^{\mu\nu} + h.c. \tag{17}$$

The decay rate in the rest frame of the decaying neutrino ν_i is given by

$$\Gamma_{\nu_i \to \nu_f + \gamma} = \frac{1}{8\pi}\left(\frac{m_i^2 - m_f^2}{m_i}\right)^3 \left(|\mu_{fi}|^2 + |\epsilon_{fi}|^2\right). \tag{18}$$

Note that due to $m_i \neq m_f$ only the transition magnetic and electric dipole moments contribute. Therefore this expression (18) is equally valid for both Dirac and Majorana neutrinos. In the simplest extensions of the Standard Model the Dirac and Majorana neutrino decay rates are given by [23]

$$\Gamma_{\nu_D} \simeq \frac{\alpha}{2}\left(\frac{3G_F}{32\pi^2}\right)^2 \left(\frac{m_i^2 - m_f^2}{m_i}\right)^3 \left(m_i^2 + m_f^2\right)|U|^2, \quad U = \sum_{l=e,\mu,\tau} U_{li}^* U_{lf}\frac{m_l^2}{m_W^2}, \tag{19}$$

and

$$\Gamma_{\nu_M} \simeq \alpha\left(\frac{3G}{32\pi^2}\right)^2 \left(\frac{m_i^2 - m_f^2}{m_i}\right)^3 \left\{(m_i + m_f)^2 |U|^2 \; - 4m_i m_f \; (Re\ U)^2\right\}. \tag{20}$$

Thus, for both Dirac and Majorana neutrinos we have

$$\Gamma_{\nu_M} = \frac{\mu_{eff}^2}{8\pi}\left(\frac{m_i^2 - m_f^2}{m_i}\right)^3 \approx 5\left(\frac{\mu_{eff}}{\mu_B}\right)^2 \left(\frac{m_i^2 - m_f^2}{m_i}\right)^3 \left(\frac{m_i}{1\ eV}\right)^3 s^{-1}, \tag{21}$$

$$\mu_{eff}^2 = |\mu_{fi}|^2 + |\epsilon_{fi}|^2. \tag{22}$$

The corresponding life time of neutrinos in respect to the radiative decay is indeed huge

$$\tau_{\nu_i \to \nu_j + \gamma} \approx 0.19\left(\frac{m_i^2}{m_i^2 - m_j^2}\right)^3 \left(\frac{eV}{m_i}\right)\left(\frac{\mu_B}{\mu_{eff}^2}\right) s. \tag{23}$$

This is because the neutrino transition moments are suppressed by the Glashow–Iliopoulos-Maiani cancellation mechanism. Note that for the diagonal moments there is no GIM cancellation.

The neutrino radiative decay has been constrained from the absence of decay photons in studies of the solar, supernova and reactor (anti)neutrino fluxes, as well as from the absence of the spectral distortions of the cosmic microwave background radiation. However, the corresponding upper bounds on the effective neutrino magnetic moments [24] are in general less stringent than the astrophysical bounds from the plasmon decay into neutrino-antineutrino pair discussed in the next section.

5. Plasmon decay to neutrino-antineutrino pair

For constraining neutrino electromagnetic properties, and obtaining upper bounds on neutrino magnetic moments in particular, the most interesting process is the plasmon decay into a neutrino-antineutrino pair [25]. This plasmon process becomes kinematically allowed in media where the

PoS(CORFU2021)057

photon behaves as a particle with an effective mass ω_γ. In the case of nonrelativistic plasma the dispersion relation for a photon (plasmon) is $\omega_\gamma^2 + \vec{k}_\gamma^2 = \omega_\gamma^2$, where $\omega_\gamma = 4\pi N_e/m_e$ is the plasmon frequency. The plasmon decay rate is given by

$$\Gamma_{\gamma^*\to\nu\bar{\nu}} = \frac{\mu_{eff}^2}{24\pi} Z \frac{(\omega_\gamma^2 - k_\gamma^2)^2}{\omega_\gamma}, \tag{24}$$

where Z is a factor which depends on the polarization of the plasmon. A plasmon decay into a neutrino-antineutrino pair transfers the energy ω_γ to neutrinos that can freely escape from a star and thus can fasten the star cooling. The corresponding energy-loss rate per unit volume is

$$Q_{\gamma^*\to\nu\bar{\nu}} = \frac{g}{(2\pi)^3} \int \Gamma_{\gamma^*\to\nu\bar{\nu}} f_{k_\gamma} \omega_\gamma d^3 k_\gamma, \tag{25}$$

where f_{k_γ} is the photon Bose-Einstein distribution function and $g = 2$ is the number of polarization states. The magnetic moment plasmon decay enhances the Standard Model photo-neutrino cooling by photon polarization tensor, and more fast star cooling slightly reduces the core temperature. These nonstandard energy losses can delay the helium ignition in low-mass red giants. This, in turn, can be related to astronomically observable luminosity of stars before and after the helium flash. And in order not to delay the helium ignition in an unacceptable way (a significant brightness increase is constraint by observations) the upper bound on the effective neutrino magnetic was obtained in [26]

$$\mu_{eff} = \left(\sum_{fi} |\mu_{fi}|^2 + |\epsilon_{fi}|^2 \right)^{\frac{1}{2}} \leq 3 \times 10^{-12} \mu_B. \tag{26}$$

Recently new analysis [27–29] of the observed properties of globular cluster stars provides new upper bounds on the effective neutrino magnetic moment $\mu_{eff} \leq (1.2\text{–}2.6) \times 10^{-12} \mu_B$ that is valid for both cases of Dirac and Majorana neutrinos.

It is interesting to compare these astrophysical bounds on the effective neutrino magnetic moments with constraints obtained in investigations of the elastic scattering of a flavour neutrino $\nu_l + e^- \to \nu_l + e^-, \ \ l = e, \mu, \tau$ (or an antineutrino $\bar{\nu}_l$) in the laboratory experiments (see Section 2). For a detailed discussion of this issue see [2] and [16].

The plasmon decay considered in the vicinity of red giants can also be used to constrain neutrino millicharges q_ν [24]. The plasmon decay to neutrino-antineutrino pair due to the neutrino millicharge q_ν is described by the Lagrangian

$$L = -iq_\nu \bar{\psi}_\nu \gamma_\mu \psi_\nu A^\mu. \tag{27}$$

In order to avoid the delay of helium ignition in low-mass red giants the millicharge should be constraint at the level $q_\nu \leq 2 \times 10^{-14} e_0$, and from the absence of the anomalous energy-dependent dispersion of the SN1987A neutrino signal it should be $q_\nu \leq 3 \times 10^{-17} e_0$ (e_0 is the value of an electron charge).

The most stringent astrophysical constraint on neutrino millicharges $q_\nu \leq 1.3 \times 10^{-19} e_0$ was obtained [30] in consideration of the impact of the *neutrino star turning* mechanism (νST) that can shift a magnetised pulsar rotation frequency. Note the most sever upper bound on the neutrino millicharges $q_\nu \sim 10^{-21} e_0$ arrives from neutrality of the hydrogen atom.

6. Neutrino spin conversion

One of the most important for astrophysics consequences of neutrino nonzero effective magnetic moments (see [2], [24], and [14]) is the neutrino helicity change $\nu_L \to \nu_R$ with the appearance of nearly sterile right-handed neutrinos ν_R . In general, this phenomena, the helicity change $\nu_L \to \nu_R$, can proceed in two different electromagnetic mechanisms: 1) the helicity change in the neutrino magnetic moment scattering on electrons (or protons and neutrons), 2) the neutrino spin and spin-flavour precession in an external magnetic field. As it was investigated for the firs time in [31] and then studied in [32, 33], there is also nonelectromagnetic mechanism of the neutrino helicity change $\nu_L \to \nu_R$: the neutrino spin and spin-flavour precession in the transversally moving matter currents or in the transversally polarized matter at rest. The existence of the proposed phenomenon [31] have been confirmed and applied in studies of astrophysical neutrino fluxes in [34–37].

For completeness note that the important astrophysical consequence of nonzero neutrino millicharges is the neutrino deviation from the rectilinear trajectory.

The detection of the SN 1987A neutrinos provides the energy-loss limits on the effective neutrino magnetic moments related to the observed duration of the neutrino signal (see [2] and [24]). In the magnetic scattering on electrons due to the change of helicity $\nu_L \to \nu_R$ the proto-neutron star formed in the core-collapse SN can cool faster since ν_R are sterile and are not trapped in a core like ν_L are trapped for a few seconds. The escaping ν_R will cool the core very efficient and fast ($\sim 1\ s$). However, the observed $5 - 10\ s$ pulse duration in Kamioka II and IMB experiments is in agreement with the standard model ν_L trapping and cooling of the star. From these considerations it was concluded that for the Dirac neutrinos the effective magnetic moment $\mu_D \geq 10^{-12}\mu_B$ is inconsistent with the SN1987A observed cooling time.

There is another approach to constrain the neutrino magnetic moment from the data on SN1987A neutrinos related to the observed neutrino energies, [2] and [24]. The right-handed neutrinos ν_R, that appear due to the helicity change in the magnetic scattering in the inner SN core, have larger energies than ν_L emitted from the neutrino sphere. Then in the magnetic moment precession process $\nu_R \to \nu_L$ the higher-energy ν_L would arrive to the detector as a signal of SN1897A. And from the absence of the anomalous high-energy neutrinos again the bound on the level $\mu_D \leq 10^{-12}\mu_B$ can be settled.

7. Future prospects

A new phase of the GEMMA project for measuring the neutrino magnetic moment is now underway at the Kalinin Power Plant in Russia. The discussed [38] next GEMMA-3 experiment, called νGEN, is aimed at the detection of coherent Neutrino–Ge Nucleus elastic scattering. It is also expected that this experiment will further increase sensitivity to the neutrino magnetic moment and will reach the level of $\mu_{\nu_e} \sim (5–9) \times 10^{-12}\mu_B$. To reach the claimed limit on the neutrino magnetic moment the νGEN experiment setup reasonably improves characteristics in respect to those of the previous editions of the GEMMA project. The most important are the following: 1) a factor of 2 increase in the total neutrino flux at the detector because of much closer location of the detector to the reactor core, 2) a factor of 3.7 increase in the total mass of the detector, 3) the energy threshold would be improved from 2.8 keV to 200 eV. Furthermore, the νGEN experimental setup is located in the new room at the Kalinin Power Plat with much better (by an order of magnitude) gamma-background conditions and on a moveable platform. The later gives an opportunity to vary online the neutrino flux and thus suppress systematic errors.

Here we recall that the first observation of the coherent elastic neutrino-nucleus scattering at the COHERENT experiment at the Spallation Neutron Source [39] can be also used for constraining neutrino electromagnetic properties. However, as it was shown in [40] and then confirmed in recent studies (see, for instance, [41]), the bounds for the flavour neutrino magnetic moments are of the order $\mu_e, \mu_\mu \sim 10^{-8}\mu_B$.

In the recent studies [42] it is shown that the puzzling results of the XENON1T collaboration [43] at few keV electronic recoils could be due to the scattering of solar neutrinos endowed with finite Majorana transition magnetic moments of the strengths lie within the limits set by the Borexino experiment with solar neutrinos [18]. The comprehensive analysis of the existing and new extended mechanisms for enhancing neutrino transition magnetic moments to the level appropriate for the interpretation of the XENON1T data and leaving neutrino masses within acceptable values is provided in [44].

In a recent paper [45] we have proposed an experimental setup to observe coherent elastic neutrino-atom scattering using electron antineutrinos from tritium decay and a liquid helium target. In this scattering process with the whole atom, that has not been observed so far, the electrons tend to screen the weak charge of the nucleus as seen by the electron antineutrino probe. Finally, we study the sensitivity of this apparatus to a possible electron neutrino magnetic moment and we find that it is possible to set an upper limit of about $\mu_\nu \leq 7 \times 10^{-13}\mu_B$, that is about two orders of magnitude smaller than the current experimental limits from GEMMA and Borexino. A corresponding experiment involving the use of an intense $1kg$ tritium antineutrino source is currently being prepared in the framework of the research program of the National Center for Physics and Mathematics in Sarov (Russia).

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