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OSCILLATIONS OF THREE GENERATIONS OF NEUTRINOS

Abstract. In this paper, we investigate the evolution of the neutrino flux propagating through dense matter and an intensive magnetic field. As an example, the magnetic field of the coupled sunspots being the sources of solar flares is considered. We assume that neutrinos possess dipole magnetic and anapole moments while the magnetic field is twisting and nonpotential, and its strength may be $\geq 10^5$ Gs. The problem is investigated within three neutrino generations. Possible resonance conversions inside the neutrino flux are studied.

Keywords: neutrino oscillations, lepton flavor violation, left-right symmetric model, heavy and light neutrinos, mixing in the neutrino sector

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ОСЦИЛЛЯЦИИ ТРЕХ ПОКОЛЕНИЙ НЕЙТРИНО

Аннотация. Исследуется эволюция потока нейтрино, проходящего через конденсированное вещество и интенсивное магнитное поле. В качестве примера интенсивного магнитного поля рассматривается магнитное поле спаренных солнечных пятен, которые являются источниками солнечных вспышек. Предполагается, что нейтрино обладает дипольным магнитным и анапольным моментами, в то время как магнитное поле является скрученным, носит непотенциальный характер и его напряженность может быть $\geq 10^5$ Гс. Проблема исследуется в рамках трех нейтринных поколений. Изучаются возможные резонансные переходы в нейтринном потоке.

Ключевые слова: нейтринные осцилляции, нарушение лептонного числа, лево-право симметричная модель, тяжелые и легкие нейтрино, смешивание в нейтринном секторе

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Introduction. Neutrinos have long remained in the focus of intensive investigations, both theoretical and experimental. In the standard model (SM) neutrinos are massless particles. However, the observation of neutrino oscillation shows the necessity of neutrino masses, which implies that the SM can be modified. At the same time, the fact that neutrinos have non-zero masses paves the way for neutrino electromagnetic interactions. Neutrino electromagnetic interactions are induced by multipole moments (MM's) which are caused by radiative corrections. Due to the smallness of neutrino MM's, these interactions become essential only in the case of intensive fields. The effects of the neutrino electromagnetic properties are searched both in astrophysics and in laboratory measurements. In particular, one might expect manifestations of these effects in astrophysical sources of neutrinos such as the Sun, magnetars, neutron stars, supernovae, etc., where strong magnetic fields are known to exsist. The interaction of neutrino MM's with their magnetic fields can give rise to the phenomenon of neutrino resonant conversions and, hence, influence the neutrino fluxes from these sources.

One of the most important current tasks in neutrino physics is to understand the full implications of the solar, atmospheric, and reactor neutrino data made available in the last years. While much of the recent effort has been concentrated on establishing the values of neutrino masses and mixings, there has also been a growing understanding that the same data may incorporate a wealth of information on such neutrino properties as their interactions both with matter and with the electromagnetic field.

In this paper, aside from neutrino MM's, we will be interested in the dipole magnetic and anapole moments. For the first time the influence of the dipole magnetic moment (DMM) on the neutrino behavior in the external magnetic field was studied in [1]. A lot of papers have since appeared in which the problems of the solar neutrinos are considered with inclusion of the neutrino DMM (see [2]). Note that the investigation of the effects related to neutrino DMM's is a way to define the neutrino nature (Dirac or Majorana). There are fundamental differences between these kinds of neutrinos. Dirac neutrinos can have both flavor conserving DMM's and transition DMM's. In so doing, active (left-handed) neutrinos propagating in an external electromagnetic field could be transformed into sterile ones being right-handed neutrinos. As far as the Majorana neutrino, it has only a transition DMM for which $\mu_{ll'} = -\mu_{l'l}$. Moreover, right-handed Majorana neutrinos are no longer sterile, and their interactions are identical with those of right- handed Dirac antineutrinos.

The anapole moment (AM) was firstly introduced in elementary particle physics in [3], while the behavior of neutrinos endowed with the AM in a non-potential magnetic field was considered in [4] upon discussing the correlation between the neutrino flux and solar flare events.

It is worth noting that most of the papers studying the behavior of the neutrino beam in the magnetic field are limited by a two-flavor approximation. Within three-neutrino generations this problem was investigated in [5] for Dirac neutrinos. The goal of the present work is to consider the behavior of Majorana neutrinos in an intensive magnetic field within three neutrino generations. As an example, we shall consider the Sun's magnetic fields. In so doing, of special interest are the magnetic fields of the solar sunspots which will be the source of the solar flare (SF). It is widely believed that the magnetic field is the main energy source of the SF [6, 7]. During the years of the active Sun, a magnetic flux of 10^{24} Gauss·cm^2 [8] erupts from the solar interior and accumulates within the sunspots giving rise to the stored magnetic field. The SF formation starts from pairing big sunspots of opposite polarity (coupled sunspots – CS's). Then the process of magnetic energy storage of the CS's begins. The duration of this initial SF stage varies from several to dozens of hours. In that case, the magnetic field value for the CS's could be increased from ~10^4 Gs up to ~10^5 Gs and upwards. It is clear that when the electron neutrinos beam passing through the magnetic field of the CS's changes its composition and we can detect this changing, then the problem of the SF's prediction will be resolved.

This paper is organized as follows. We start by finding the evolution equation of the neutrino beam which travels through the region of the CS's in the preflare period. In section III we establish all possible resonance conversions of the neutrino system under study and derive the expression for the survival probability of the electron neutrino. Finally, in section IV, some conclusions are drawn.

1. Neutrino evolution equation. In order to take into account the interaction effects of the neutrinos with the electromagnetic fields, the neutrino system under study must contain both left-handed and right-handed neutrinos. Since in the Majorana case right-handed neutrinos interact as right-handed Dirac antineutrinos, then the ν_{lR} is often denoted by $\overline{\nu}_{lR}$, or only $\overline{\nu}_l$, and named an electron antineutrino. In what follows, for right-handed neutrinos we shall draw on the designation $\overline{\nu}_{lR}$.

Note that the neutrino magnetic moment predicted by the SM is proportional to the neutrino mass [9]

$$\mu_\nu = \frac{3eG_F m_\nu}{8\sqrt{2}pi^2} = 10^{-19}\mu_B\left(\frac{m_\nu}{eV}\right), \tag{1}$$

and, as a consequence, cannot lead to any observable effects in real fields. Therefore, if one uses the values of the neutrino magnetic moments which are close to the upper experimental bounds, then one should employ the SM extension containing right-handed charged currents and/or charged Higgs bosons. As an example of such a model, the left-right symmetric model (LRM) based on the $SU(2)_R \times SU(2)_L \times U(1)_{B-L}$ gauge group can exist [10–12]. The Higgs sector structure of the LRM specifies the neutrino nature. When the Higgs sector of the LRM contains a bidoublet $\Phi(1/2,1/2,0)$ and two triplets $\Delta_L(1,0,2)$, $\Delta_R(0,1,2)$ [13] (in brackets the values of S_L^W, S_R^W, and B–L are given,

S_L^W (S_R^W) being the weak left (right) isospin while B and L being the baryon and lepton numbers), then the neutrino has a Majorana nature. If the neutrino is a Dirac particle, the Higgs sector must hold a bidoublet $\Phi\ (1/2,1/2,0)$ and two doublets $\chi_L(1/2,0,1)$, $\chi_R(0,1/2,1)$ [14].

In the LRM the Lagrangian describing neutrino interaction with $W^{\pm}, Z$ gauge bosons has the form

$$\mathcal{L}_g = \frac{g}{2\sqrt{2}}\Big[\overline{\nu}_l(x)\gamma^{\mu}(1-\gamma_5)l(x)W_{\mu}^{*}(x)+\overline{l}(x)\gamma^{\mu}(1-\gamma_5)\nu_l(x)W_{\mu}(x)\Big]+$$
$$+\frac{g}{4\cos\theta_W}\Big\{\overline{\nu}_l(x)\gamma^{\mu}(1-\gamma_5)\nu_l(x)+\overline{l}(x)\gamma^{\mu}\Big[4\sin^2\theta_W-1+\gamma_5\Big]l(x)\Big\}Z_{\mu}(x), \tag{2}$$

and it must be added to the Lagrangians which are responsible for neutrino interactions with additional gauge bosons $W'^{\pm}$, Z', and singly charged Higgs bosons $h^{(\pm)}$, $\tilde{\delta}^{\pm}$ [15]. Inasmuch as the masses of $W'^{\pm}$, $Z'^{\pm}$, and $h^{(\pm)}$ lie at the TeV scale [16], then one may neglect their contributions in the neutrino Lagrangian (2). On the other hand, the $\tilde{\delta}^{\pm}$ boson does not interact with the quarks, and as a result, the more firm data for deriving the bounds on the $m_{\tilde{\delta}}$ follow from the electroweak processes. For example, results from LEP experiments (ALEPH, DELPHI, L3, and OPAL) gave the bound $m_{H^+} > 80$ GeV [16]. In the LRM, the interaction between the neutrino and the $\tilde{\delta}^{\pm}$ boson is described by the Lagrangian

$$\mathcal{L}_{\tilde{\delta}} = \frac{f_{ll'}}{\sqrt{2}}\overline{l}^{\,c}(x)(1-\gamma^5)\nu_{l'}(x)\tilde{\delta}^{+}(x), \tag{3}$$

where $f_{ll'}$ is a triplet Yukawa coupling constant (TYCC) and $l,l' = e, \mu, \tau$. Then this interaction may lead to changes of the matter potential on the value

$$V_{ll'}^{\tilde{\delta}} = -\frac{f_{el}f_{el'}}{m_{\tilde{\delta}}}n_e, \tag{4}$$

(n_e is electron density) that could be larger by several tens of percent than SM [17]. Further, for the sake of simplicity, we shall assume that only diagonal TYCCs are different from zero.

The existence of electromagnetic multipole moments caused by the radiative corrections could be taken into account in terms of the effective Lagrangian

$$\mathcal{L}_{\text{eff}} = \frac{i}{2}\overline{\nu}_l(x)\Big[\mu_{ll'}\sigma^{\beta\lambda}+a_{ll'}(\partial^{\beta}\gamma^{\lambda}-\partial^{\lambda}\gamma^{\beta})\Big](1-\gamma_5)\nu_{l'}(x)F_{\lambda\beta}(x)+\text{conj.}=$$
$$=\frac{i}{2}\overline{\nu}_a(x)\Big[\mu_{ab}\sigma^{\beta\lambda}+a_{ab}(\partial^{\beta}\gamma^{\lambda}-\partial^{\lambda}\gamma^{\beta})\Big](1-\gamma_5)\nu_b(x)F_{\lambda\beta}(x)+\text{conj.}, \tag{5}$$

where the indexes a and b refer to the mass eigenstate basis ($a, b = 1,2,3$), μ_{ab} (a_{ab}) are the dipole magnetic (anapole) moments of the mass eigenstates, and $F_{\lambda\mu} = \partial_{\lambda}A_{\mu} - \partial_{\mu}A_{\lambda}$.

Let us assume that the magnetic fields in which the neutrino beam travels have the nonpotential character

$$(\text{rot}\,\mathbf{B})_z = 4\pi j_z, \tag{6}$$

and exhibit the geometrical phase $\Phi(z)$

$$B_x \pm iB_y = B_{\perp}e^{\pm i\Phi(z)}, \tag{7}$$

which is defined by the following relation

$$\Phi(z) = \frac{\alpha\pi}{L_{mf}}z. \tag{8}$$

So, we assume that the magnetic fields are in existence over a distance L_{mf} and twist by the angle $\alpha\pi$. We shall also assert that the nondiagonal neutrino anapole moments are equal to zero while the diagonal ones are equal in magnitude $a_{\nu_e\nu_e} = a_{\nu_\mu\nu_\mu} = a_{\nu_\tau\nu_\tau} = a_{\nu\nu}$.

Then in the flavor basis the evolution equation can be represented as follows

$$i\frac{d}{dz}\begin{pmatrix} \nu_{eL} \\ \nu_{\mu L} \\ \nu_{\tau L} \\ \overline{\nu}_{eR} \\ \overline{\nu}_{\mu R} \\ \overline{\nu}_{\tau R} \end{pmatrix} = \mathcal{H}^M \begin{pmatrix} \nu_{eL} \\ \nu_{\mu L} \\ \nu_{\tau L} \\ \overline{\nu}_{eR} \\ \overline{\nu}_{\mu R} \\ \overline{\nu}_{\tau R} \end{pmatrix}, \tag{9}$$

where

$$\mathcal{H}^M = \mathcal{U}\begin{pmatrix} E_1 & 0 & 0 & 0 & 0 & 0 \\ 0 & E_2 & 0 & 0 & 0 & 0 \\ 0 & 0 & E_3 & 0 & 0 & 0 \\ 0 & 0 & 0 & E_1 & 0 & 0 \\ 0 & 0 & 0 & 0 & E_2 & 0 \\ 0 & 0 & 0 & 0 & 0 & E_3 \end{pmatrix}\mathcal{U}^{-1} +$$

$$+\begin{pmatrix} V'_{eL}+A_L & 0 & 0 & 0 & \mu_{e\mu}B_\perp & -\mu_{e\tau}B_\perp \\ 0 & V_{\mu L}+A_L & 0 & -\mu_{e\mu}B_\perp & 0 & \mu_{\mu\tau}B_\perp \\ 0 & 0 & V_{\tau L}+A_L & \mu_{e\tau}B_\perp & -\mu_{\mu\tau}B_\perp & 0 \\ 0 & -\mu_{e\mu} & \mu_{e\tau} & -V'_{eL} - A_R & 0 & 0 \\ \mu_{e\mu}B_\perp & 0 & -\mu_{\mu\tau}B_\perp & 0 & -V_{\mu L} - A_R & 0 \\ -\mu_{e\tau}B_\perp & 0 & 0 & 0 & 0 & -V_{\tau L} - A_R \end{pmatrix}, \tag{10}$$

$$\mathcal{U} = \begin{pmatrix} D & 0 \\ 0 & D \end{pmatrix},$$

$$D = \exp(i\lambda_7\psi)\exp(i\lambda_5\phi)\exp(i\lambda_2\omega) = \begin{pmatrix} c_\omega c_\phi & s_\omega c_\phi & s_\phi \\ -s_\omega c_\psi - c_\omega s_\psi s_\phi & c_\omega c_\psi - s_\omega s_\psi s_\phi & s_\psi c_\phi \\ s_\omega s_\psi - c_\omega c_\psi s_\phi & -c_\omega s_\psi - s_\omega c_\psi s_\phi & c_\psi c_\phi \end{pmatrix},$$

$\psi = \theta_{23}$, $\phi = \theta_{13}$, $\omega = \theta_{12}$, $s_\psi = \sin\psi$, $c_\psi = \cos\psi$, and so on, the λ's are Gell-Mann matrices related to the spin-one matrices of the $SO(3)$ group, V'_{eL} ($V_{\mu L}$, $V_{\tau L}$) is a matter potential describing the interaction of the ν_{eL} ($\nu_{\mu l}$, $\nu_{\tau L}$) neutrinos with a dense matter,

$$V'_{eL} = V_{eL} + V_{ee}^{\tilde{\delta}}, \quad V_{eL} = \sqrt{2}G_F(n_e - n_n/2), \quad V_{\mu L} = V_{\tau L} = -\sqrt{2}G_F n_n/2,$$

$$A_L = 4\pi a_{\nu_L\nu_L} j_z - \dot{\Phi}/2, \quad A_R = 4\pi a_{\nu_R\nu_R} j_z - \dot{\Phi}/2,$$

and n_n is neutron density.

Of course, to get the survival probabilities of definite flavor neutrinos, we can turn to the Hamiltonian (10) and immediately find all possible resonant transitions in a flavor basis. However, even though we work with the three component neutrino wave function $\Psi^T = (\nu_{eL}, \nu_{\mu L}, \nu_{\tau L})$ the physical implications will be far from transparent [18]. For this reason we shall rotate the flavor basis to make the physical meanings more obvious. In so doing, when the mixing angles ψ and ϕ tend to zero, our results must convert into those obtained within the two-flavor approximation. This could be arranged by the following transformation

$$\begin{pmatrix} \nu_1'^M \\ \nu_2'^M \\ \nu_3'^M \\ \overline{\nu}_1'^M \\ \overline{\nu}_2'^M \\ \overline{\nu}_3'^M \end{pmatrix} = \mathcal{U}' \begin{pmatrix} \nu_{eL} \\ \nu_{\mu L} \\ \nu_{\tau L} \\ \overline{\nu}_{eR} \\ \overline{\nu}_{\mu R} \\ \overline{\nu}_{\tau R} \end{pmatrix}, \qquad (11)$$

where

$$\mathcal{U}' = \begin{pmatrix} D' & 0 \\ 0 & D' \end{pmatrix}, \quad D' = \exp(-i\lambda_5\phi)\exp(-i\lambda_7\psi) = \begin{pmatrix} c_\phi & 0 & s_\phi \\ -s_\phi s_\psi & c_\psi & c_\phi s_\psi \\ -s_\phi c_\psi & -s_\phi & c_\phi c_\psi \end{pmatrix}.$$

In this basis the Hamiltonian H'^M has the form

$$\mathcal{H}'^M = \begin{pmatrix} B_\upsilon + \Lambda & M \\ -M & B_\upsilon + \overline{\Lambda} \end{pmatrix}, \qquad (12)$$

where

$$B_\upsilon = \begin{pmatrix} -\delta^{12} c_{2\omega} & \delta^{12} s_{2\omega} & 0 \\ \delta^{12} s_{2\omega} & \delta^{12} c_{2\omega} & 0 \\ 0 & 0 & \delta^{31} + \delta^{32} \end{pmatrix}, \quad \Lambda = \begin{pmatrix} V_{eL}^{\text{eff}} c_\phi^2 & 0 & V_{eL}^{\text{eff}} s_{2\phi}/2 \\ 0 & 0 & 0 \\ V_{eL}^{\text{eff}} s_{2\phi}/2 & 0 & V_{eL}^{\text{eff}} s_\phi^2 \end{pmatrix},$$

$$\overline{\Lambda} = \begin{pmatrix} -V'_{eL} c_\phi^2 - V_{\mu L}\left(1 + s_\Phi^2\right) - A_{\nu\nu}^{(\Sigma)} & 0 & -V_{eL}^{\text{eff}} s_{2\phi}/2 \\ 0 & -2V_{\mu L} - A_{\nu\nu}^{(\Sigma)} & 0 \\ -V_{eL}^{\text{eff}} s_{2\phi}/2 & 0 & -V'_{eL} s_\phi^2 - V_{\mu L}\left(1 + c_\Phi^2\right) - A_{\nu\nu}^{(\Sigma)} \end{pmatrix},$$

$$A_{\nu\nu}^{(\Sigma)} = 4\pi\left(a_{\nu_L\nu_L} + a_{\nu_R\nu_R}\right) j_z - \dot{\Phi},$$

$$M = \begin{pmatrix} B_\perp & 0 & 0 \\ 0 & B_\perp & 0 \\ 0 & 0 & B_\perp \end{pmatrix} \times$$

$$\times \begin{pmatrix} 0 & \mu_{e\mu} c_\psi c_\phi + \mu_{e\tau} s_\psi c_\phi + \mu_{\mu\tau} s_\phi & \mu_{e\mu} s_\psi + \mu_{e\tau} c_\psi \\ -\mu_{e\mu} c_\psi c_\phi - \mu_{e\tau} s_\psi c_\phi + \mu_{\mu\tau} s_\phi & 0 & -\mu_{e\mu} c_\psi s_\phi - \mu_{e\tau} s_\psi s_\phi + \mu_{\mu\tau} c_\phi \\ -\mu_{e\mu} s_\psi + \mu_{e\tau} c_\psi & \mu_{e\mu} c_\psi s_\phi + \mu_{e\tau} s_\psi s_\phi - \mu_{\mu\tau} c_\phi & 0 \end{pmatrix},$$

$$\delta^{ik} = \frac{m_i^2 - m_k^2}{4E}, \qquad V_{eL}^{\text{eff}} = \sqrt{2} G_F n_e + V_{ee}^{\tilde{\delta}}.$$

From Eq. (12) it follows that for this basis choice the resonance conditions will not contain the angle ψ while the ψ-dependence will be transported to the resonance widths and the oscillation lengths.

For Dirac neutrinos in the basis $\Psi'^T=(\nu_1'^D,\nu_2'^D,\nu_3'^D,\overline{\nu}_1'^D,\overline{\nu}_2'^D,\overline{\nu}_3'^D)$ where

$$\begin{pmatrix}\nu_1'^D\\ \nu_2'^D\\ \nu_3'^D\\ \overline{\nu}_1'^D\\ \overline{\nu}_2'^D\\ \overline{\nu}_3'^D\end{pmatrix}=\mathcal{U}'\begin{pmatrix}\nu_{eL}\\ \nu_{\mu L}\\ \nu_{\tau L}\\ \nu_{eR}\\ \nu_{\mu R}\\ \nu_{\tau R}\end{pmatrix}, \tag{13}$$

the Hamiltonian H'^D is given by the expression

$$\mathcal{H}'^D=\begin{pmatrix}B_\upsilon+\Lambda & M^D\\ M^D & B_\upsilon+\Lambda^D\end{pmatrix}, \tag{14}$$

in which

$$M^D=D'\begin{pmatrix}\mu_{ee}B_\perp & \mu_{e\mu}B_\perp & \mu_{e\tau}B_\perp\\ \mu_{e\mu}B_\perp & \mu_{\mu\mu}B_\perp & \mu_{\mu\tau}B_\perp\\ \mu_{e\tau}B_\perp & \mu_{\mu\tau}B_\perp & \mu_{\tau\tau}B_\perp\end{pmatrix}D'^{-1}, \tag{15}$$

$$\Lambda^D=\begin{pmatrix}\dot{\Phi}/2-A_L-V_{\mu L} & 0 & 0\\ 0 & \dot{\Phi}/2-A_L-V_{\mu L} & 0\\ 0 & 0 & \dot{\Phi}/2-A_L-V_{\mu L}\end{pmatrix}. \tag{16}$$

Since in the Dirac neutrino case the masses of the singly charged Higgs bosons lie at the TeV scale, then in the expression for H'^D we may neglect their contributions. It should be also stressed that here the ν_{eR}, $\nu_{\mu R}$ and $\nu_{\tau R}$ are sterile states.

2. Resonance transitions in the neutrino system. Our next task is to identify possible resonance conversions of the neutrino beam which travels in the region of the coupled sunspots being the source of solar flares. Remember that for the resonance conversion the following requirements are necessary: (i) the resonance condition must be fulfilled; (ii) the resonance width must be nonzero; (iii) the neutrino beam must pass a distance comparable with the oscillation length. Then to verify the fulfillments of these requirements we should make numerical estimations. With this aim the experimental data of the values of the multipole moments should be used.

Considering Majorana three-neutrino mixing and using the solar and KamLAND data in the work [19] the following result was obtained

$$\sqrt{|\mu_{12}|^2+|\mu_{23}|^2+|\mu_{31}|^2}<4.0\cdot10^{-10}\mu_B,\quad (90\ \%\quad \text{CL}). \tag{17}$$

On the other hand, the addition of the Rovno [20], TEXONO [21] and MUNU [22] constraints leads to the inequality

$$\sqrt{|\mu_{12}|^2+|\mu_{23}|^2+|\mu_{31}|^2}<1.8\cdot10^{-10}\mu_B,\quad (90\ \%\quad \text{CL}). \tag{18}$$

The experimental bounds on Dirac neutrinos could be found in [23, 24]. The upper bounds on μ_{ν_e} lie in the range from $2.9\cdot10^{-11}\mu_B$ to $1.1\cdot10^{-9}\mu_B$, while those on μ_{ν_μ} are from $6.8\cdot10^{-10}\mu_B$ to $8.5\cdot10^{-10}\mu_B$. For the τ-neutrino, the bounds on μ_{ν_τ} are less restrictive (see, for example, [24]), and the current upper limit is $3.9\cdot10^{-7}\mu_B$.

As for the anapole moment is concerned, we remind that its phenomenology is similar to that of neutrino charge radii (NCR) $< r^2(\nu_i) >$. In that case, the limits on the NCR apply also to the neutrino anapole moments multiplied by 6. The measurement of the elastic neutrino-electron scattering in the TEXONO experiment leads to the following bounds on the electron NCR [25]

$$-2.1\cdot 10^{-32}\,\text{cm}^2 \le < r^2_{\nu_e} > \le 3.3\cdot 10^{-32}\,\text{cm}^2. \tag{19}$$

So, the existing experimental data allow us to use the following values for neutrino MM's

$$\mu_{\nu_l \nu_{l'}} = 10^{-10}\mu_B, \qquad |a_{\nu_l \nu_{l'}}| = 3\cdot 10^{-40}\,\text{esu}\cdot\text{cm}^2. \tag{20}$$

Now we could proceed to the examination of the resonance conversions of the neutrino system under study. In so doing we shall infer that the resonance localization places are arranged rather far from one another what enables us to treat them as independent ones. Let us begin with the ${\nu'}_1^M \leftrightarrow {\nu'}_2^M$ transitions. In this case the resonance condition, the resonance width and the oscillation length are given by the expressions

$$-2\delta^{12}c_{2\omega} + V_{eL}^{\text{eff}} c_\phi^2 = 0, \tag{21}$$

$$\Gamma\left(\nu_1'^M \leftrightarrow \nu_2'^M\right) \simeq \frac{\sqrt{2}\delta^{12}s_{2\omega}}{G_F}, \tag{22}$$

$$L_{\nu_1'^M \nu_2'^M} = \frac{2\pi}{\sqrt{\left(2\delta^{12}c_{2\omega} - V_{eL}^{\text{eff}} c_\phi^2\right)^2 + \left(\delta^{12}s_{2\omega}\right)^2}}. \tag{23}$$

Comparing the expressions (21)–(23) with those for the ${\nu'}_1^D \leftrightarrow {\nu'}_2^D$ resonance we make sure that they coincide with each other when $V_{ee}^{\tilde{\delta}} = 0$. So, in the Majorana case this resonance occurs at a lower density. From Eq. (23) it follows that the oscillation length achieves its maximum value at the resonance and the following relation takes place

$$\Gamma\left(\nu_1'^M \leftrightarrow \nu_2'^M\right) = \frac{2\sqrt{2}\pi}{G_F (L_{\nu_1'^M \nu_2'^M})_{\max}}. \tag{24}$$

It should be stressed that the connection between the maximum value of the oscillation length and the resonance width holds for any resonance conversion. It is clear that the resonance under question belongs to the kind of matter-induced resonances. When $\phi = 0$ then the Eqs. (21)–(23) convert to the corresponding expressions for $\nu_{eL} \leftrightarrow \nu_{\mu L}$ resonance conversion found in two flavor approximations (Micheev – Smirnov – Wolfenstein — MSW resonance [26, 27]). As the analysis shows, in the Sun's condition this resonance will be fulfilled before the convective zone and, as a result, has no bearing on the solar flares.

Further we shall consider the $\nu_1'^M \leftrightarrow \nu_2'^M$ and $\overline{\nu}_1'^D \leftrightarrow \overline{\nu}_2'^D$ resonance conditions. Then we have

$$-2\delta^{12}c_{2\omega} + V'_{eL}c_\Phi^2 + V_{\mu L}(1 + s_\Phi^2) + A_{\nu\nu}^{(\Sigma)} = 0, \tag{25}$$

$$\left(L_{\nu_1'^M \overline{\nu}_2'^M}\right)_{\max} \simeq \frac{2\pi}{\mu_{12}^M B_\perp}, \tag{26}$$

where

$$\mu_{12}^M = \mu_{e\mu}c_\psi c_\phi + \mu_{e\tau}s_\psi c_\phi + \mu_{\mu\tau}s_\phi,$$

while for the latter the corresponding expressions will look like

$$-2\delta^{12}c_{2\omega} + V_{eL}c_{\Phi}^2 + V_{\mu L}s_{\Phi}^2 + A_{\nu_L \nu_L} - \dot{\Phi}/2 = 0, \tag{27}$$

$$\left(L_{\nu_1'^D \overline{\nu}_2'^D}\right)_{\max} \simeq \frac{2\pi}{\mu_{12}^D B_\perp}, \tag{28}$$

where

$$\mu_{12}^D = (\mu_{\tau\tau} - \mu_{\mu\mu})c_\psi s_\psi s_\phi + (\mu_{e\mu}c_\psi - \mu_{e\tau}s_\psi)c_\phi + \mu_{\mu\tau}(s_\psi^2 - c_\psi^2)s_\phi. \tag{29}$$

From comparing the foregoing expressions one may make the conclusion that the conditions of observing the $\nu_1'^M \leftrightarrow \overline{\nu}_2'^M$ and $\nu_1'^D \leftrightarrow \overline{\nu}_2'^D$ resonances are little different from each other. Then, considering this resonance in the region of the CS's we may use the results of the work [5] and argue that the $\nu_1'^M \leftrightarrow \overline{\nu}_2'^M$ resonance may also occur only at the cost of the magnetic field and, as a result, this resonance falls into the kind of magnetic-induced resonances. The value of $(\delta^{12})_{\min} \simeq 10^{-12}$ eV entering into the resonance condition (25) could be compensated by the twisting frequency $\dot{\Phi}$ only. For example, at $B_\perp = 10^5$ Gs the value of $\dot{\Phi}$ being equal to $-10\pi / L_{mf}$ ensures the existence of the $\nu_1'^M \leftrightarrow \overline{\nu}_2'^M$ resonance.

Setting $\psi = \phi = 0$ in Eqs. (25), (26), we get the resonance condition and the oscillation length found in the two flavor approximations (FA) for the $\nu_{eL} \leftrightarrow \overline{\nu}_{\mu R}$ resonance. Therefore, we may argue that the $\nu_{eL} \leftrightarrow \overline{\nu}_{\mu R}$ resonance is an analog of the $\nu_1^M \leftrightarrow \overline{\nu}_2'^M$ resonance in the two FA. The similar connection exists for Dirac neutrinos as well.

We now turn to the $\nu_1'^M \leftrightarrow \overline{\nu}_1'^M$ and $\nu_1^D \leftrightarrow \overline{\nu}_1'^D$ resonances. In the Majorana neutrino case the resonance width is equal to zero and, as a result, the $\nu_1'^M \leftrightarrow \overline{\nu}_1'^M$ resonance is not observed.

For the Dirac neutrino the resonance condition and the maximum values of the oscillation length are defined by the following expressions

$$V_{eL}c_\phi^2 + V_{\mu L}s_\phi^2 + A_{\nu_L \nu_L} - \dot{\Phi}/2 = 0, \tag{30}$$

$$\left(L_{\nu_1'^D \overline{\nu}_1'^D}\right)_{\max} \simeq \frac{2\pi}{\mu_{11}^D B\perp}, \tag{31}$$

where

$$\mu_{11}^D = \mu_{ee}c_\phi^2 + (\mu_{\mu\mu}s_\psi^2 + \mu_{\tau\tau}c_\psi^2)s_\phi^2 - 2(\mu_{e\mu}s_\psi + \mu_{e\tau}c_\psi)s_\phi c_\phi - 2\mu_{\mu\tau}s_\phi^2 c_\psi s_\psi. \tag{32}$$

For the solar neutrinos traveling the CS's magnetic field the situation with

$$j_z = 0, \text{ but } V_{eL}c_\phi^2 + V_\mu s_\phi^2 \simeq \dot{\Phi}, \tag{33}$$

is not realistic, inasmuch as in this case the twisting magnetic field must exist over the distance being much longer than the solar radius. Therefore, the $\nu_1^D \leftrightarrow \overline{\nu}_1^D$ resonance may be in existence only under fulfillment of the following requirements

$$\dot{\Phi} \ll V_{eL}c_\phi^2 + V_{\mu L}s_\phi^2, \text{ but } V_{eL}c_\phi^2 + V_{\mu L}s_\phi^2 \simeq -4\pi a_{\nu_L \nu_L} j_z. \tag{34}$$

For example, for this resonance to occur in the Sun's corona the value of j_z must be as large as $5.5 \cdot 10^{-3}$ A/cm^2.

We also see that setting $\psi = \phi = 0$ in Eqs. (30)–(32) we may deduce all the expressions for the $\nu_{eL} \leftrightarrow \nu_{eR}$ resonance conversion obtained in the two FA [28]. That admits us to consider the $\nu_{eL} \rightarrow \nu_{eR}$ resonance as the analog of the $\nu_1^D \leftrightarrow \overline{\nu}_1^D$ resonance in two FA.

Further we consider the $\nu'^M_1 \leftrightarrow \nu'^M_3$ and $\nu'^D_1 \leftrightarrow \nu'^D_3$ resonances. In the Hamiltonians H'^M and H'^D the quantity $\delta\Sigma = \delta^{31} + \delta^{32}$ is present. Since it offers the dominant term then the ν'^M_3 and ν'^D_3 states are decoupled from the remaining ones. As a result the $\nu'^M_1 \leftrightarrow \nu'^M_3$ and $\nu'^D_1 \leftrightarrow \nu'^D_3$ oscillations controlled by the Σ term could be simply averaged out in the final survival probability for neutrinos of any flavor. For similar reasons the $\nu_1'^M \leftrightarrow \overline{\nu}_3'^M$ and $\nu_1'^D \leftrightarrow \overline{\nu}_3'^D$ resonances are also not observed.

Now we proceed to the investigation of the resonance conversions of $\nu_2'^M$ and $\nu_2'^D$ neutrinos. It is oblivious that the $\nu_2'^M \leftrightarrow \nu_3'^M$ and $\nu_2'^D \leftrightarrow \nu_3'^D$ resonances transitions are forbidden. The same is also true for the $\nu_2'^M \leftrightarrow \overline{\nu}_3'^M$ and $\nu_2'^D \leftrightarrow \overline{\nu}_3'^D$. So, we should consider only the $\nu_2'^M \leftrightarrow \overline{\nu}_1'^M$ and $\nu_2'^D \leftrightarrow \overline{\nu}_1'^D$ transitions. In the former case the resonance condition and the maximal oscillation length are defined by

$$2\delta^{12}c_{2\omega} + V_{eL}c_\Phi^2 + V_{\mu L}\left(1 + s_\Phi^2\right) + A_{\nu\nu}^{(\Sigma)} = 0, \quad (35)$$

$$\left(L_{\nu_2'^M \overline{\nu}_1'^M}\right)_{\max} \simeq \frac{2\pi}{\left(\mu_{e\mu}c_\psi c_\phi + \mu_{e\tau}s_\psi c_\phi + \mu_{\mu\tau}s_\phi\right)B_\perp}. \quad (36)$$

In the latter case we deal with the following expressions

$$2\delta^{12}c_{2\omega} + V_{\mu L} + 4\pi a_{\nu_L\nu_L} j_z - \dot{\Phi} = 0, \quad (37)$$

$$\left(L_{\nu_2'^D \overline{\nu}_1'^D}\right)_{\max} \simeq \frac{2\pi}{\mu_{12}^D B_\perp}. \quad (38)$$

Since in the region of the CS's the matter potential is much more less than δ^{12} then just as for the Majorana neutrino, so for the Dirac neutrino this resonance may be realized only at the cost of the magnetic field. For example, assuming the Dirac neutrino nature and setting

$$j_z = 0, \qquad B_\perp = 10^6\,\text{Gs},$$

we get $L_{mf} \simeq (L_{\nu_2'^D \overline{\nu}_1'^D})_{\max} \simeq 10^8$ cm at the twist frequency $3\pi / L_{mf}$.

It is clear that the $\nu_{\mu L} \leftrightarrow \overline{\nu}_{eR}$ resonance represents the analog of the $\nu_2'^M \leftrightarrow \overline{\nu}_1'^M$ and $\nu_2'^D \leftrightarrow \overline{\nu}_1'^D$ resonances in the two FA.

The $\nu_2'^M \leftrightarrow \overline{\nu}_2'^M$ and $\nu_2'^D \leftrightarrow \overline{\nu}_2'^D$ resonances are the next ones. By virtue of the fact that the resonance width of the $\nu_2'^M \leftrightarrow \overline{\nu}_2'^M$ resonance is equal to zero, it appears to be forbidden.

For the Dirac neutrino the resonance condition and the maximal oscillation length are given by the expressions

$$V_{\mu L} + 4\pi a_{\nu_L\nu_L} j_z - \dot{\Phi} = 0, \quad (39)$$

$$\left(L_{\nu_2^D \overline{\nu}_2^D}\right)_{\max} \simeq \frac{2\pi}{\mu_{22}^D B_\perp}, \quad (40)$$

where

$$\mu_{22}^D = \mu_{\mu\mu}c_\psi^2 + \mu_{\tau\tau}s_\psi^2 - 2\mu_{\mu\tau}c_\psi s_\psi. \quad (41)$$

This resonance could be observed in the solar corona. For example, assuming that the twist frequency $\dot{\Phi}$ of the magnetic field of the CS's is much less than the corona matter potential ($V_{\mu L} \simeq 10^{-30}/6$ eV), we see that the resonance condition (39) will be fulfilled provided $j_z \simeq 10^{-3}$ A/cm^2. It is clear that the $\nu_{\mu l} \leftrightarrow \mu_{\mu R}$ resonance is the analog of the $\nu_2'^D \leftrightarrow \overline{\nu}_1'^D$ resonance in the two FA.

Now let us discuss the $\nu_3'^D \leftrightarrow \overline{\nu}_3'^D$ resonance (because of the zero width the $\nu_3'^M \leftrightarrow \overline{\nu}_3'^M$ resonance is not observed). The corresponding formulas are the following

$$V_{eL}s_\phi^2 + V_{\mu L}c_\phi^2 + 4\pi a_{\nu_L\nu_L} j_z - \dot{\Phi} = 0, \quad (42)$$

$$\left(L_{\nu_3'^D \overline{\nu}_3'^D}\right)_{\max} \simeq \frac{2\pi}{\mu_{33}^D B_\perp}, \quad (43)$$

where

$$\mu_{33}^D = \mu_{ee}s_\phi^2 + (\mu_{\mu\mu}s_\psi^2 + \mu_{\tau\tau}c_\psi^2)c_\phi^2 + 2(\mu_{e\mu}s_\psi + \mu_{e\tau}c_\psi)c_\phi s_\phi + 2\mu_{\mu\tau}c_\phi^2 c_\psi s_\psi. \tag{44}$$

Again, just as in the cases of the $\nu_1'^D \leftrightarrow \overline{\nu}_1'^D$ and $\nu_2'^D \leftrightarrow \overline{\nu}_2'^D$ resonances the $\nu_3'^D \leftrightarrow \overline{\nu}_3'^D$ resonance could be realized in the region of the CS's.

Let us show that the formulas of the three neutrino generations convert into those for the two FA. Owing to our assumption that the resonance regions are well separated, we could interpret the resonances independently from each other. Next, for our purpose, we need an expression for the probability of oscillatory transitions between two neutrino states.

As such, we could take the expression related to the simplest case, namely, when n_e, n_n, $\dot{\Phi}$ and j_z are constant. It has the form

$$P_{\nu_\alpha \leftrightarrow \nu_\beta} = \sin^2\theta_{\text{eff}} \sin^2\left(\frac{\pi z}{L_{\nu_\alpha \nu_\beta}}\right), \tag{45}$$

where $\sin^2\theta_{\text{eff}}$ and $L_{\nu_\alpha\nu_\beta}$ are expressed through the elements of the effective Hamiltonian matrix H as

$$\sin^2\theta_{\text{eff}} = \frac{4\mathcal{H}_{\alpha\beta}^2}{4\mathcal{H}_{\alpha\beta}^2 + (\mathcal{H}_{\beta\beta}^2 - \mathcal{H}_{\alpha\alpha}^2)^2}, \quad L_{\nu_\alpha\nu_\beta} = \frac{2\pi}{\sqrt{4\mathcal{H}_{\alpha\beta}^2 + (\mathcal{H}_{\beta\beta}^2 - \mathcal{H}_{\alpha\alpha}^2)^2}}. \tag{46}$$

Using explicit forms of H'^M and H'^D we could determine all transition probabilities in the bases of H′ states both for Majorana and Dirac neutrinos. Then, taking into account the flavor content of these states, we could find the survival probabilities. For the electron neutrinos we obtain

$$P_{\nu_e\nu_e}^D = 1 - \Big\{c_\phi^2(P_{\nu_1'^D\overline{\nu}_1'^D} + P_{\nu_1'^D\nu_2'^D} + P_{\nu_1'^D\overline{\nu}_2'^D}) + s_\phi^4 s_\psi^2 P_{\nu_1'^D\nu_2'^D} + s_\phi^2 s_\psi^2 P_{\nu_1'^D\overline{\nu}_1'^D} +$$
$$+ s_\psi^2 s_\phi^2(2 - s_\phi^2 s_\psi^2)P_{\nu_1'^D\overline{\nu}_2'^D} + s_\phi^2 c_\psi^2(2 - c_\phi^2 c_\psi^2)P_{\nu_3'^D\overline{\nu}_3'^D}\Big\}, \tag{47}$$

in the Dirac neutrino case and

$$P_{\nu_e\nu_e}^M = 1 - \Big\{c_\phi^2(P_{\nu_1'^M\nu_2'^M} + P_{\nu_1'^M\overline{\nu}_2'^M}) + s_\phi^4 s_\psi^2 P_{\nu_1'^M\nu_2'^M} + s_\phi^2 s_\psi^2 P_{\nu_2'^M\overline{\nu}_1'^M}\Big\}, \tag{48}$$

for the Majorana neutrino case. When in Eqs. (47), (48) we put ϕ and ψ equal to zero, then, as would be expected, these expressions convert into the electron neutrino survival probability found in two FA

$$P_{\nu_e\nu_e} = 1 - \left\{\lambda P_{\nu_{eL}\nu_{eR}} + P_{\nu_{eL}\nu_{\mu L}} + P_{\nu_{eL}\nu_{\mu R}}\right\}, \tag{49}$$

where $\lambda = 1$ ($\lambda = 0$) for Dirac (Majorana) neutrinos. Using Eqs. (47), (48), (49) we could find the discrepancy between the predictions obtained within the two FA and the three neutrino generations

$$\Delta P_{\nu_e\nu_e}^{M,D} = P_{\nu_e\nu_e}^{M,D} - P_{\nu_e\nu_e}^{M,D}\Big|_{\psi=\phi=0}. \tag{50}$$

Conclusions. In this work we considered the behavior of the neutrino flux in dense matter and an intensive magnetic field within three neutrino generations. We investigated both Majorana and Dirac neutrinos. It was assumed that the neutrinos possess both the dipole magnetic moment and the anapole moment. As for the magnetic field, it had a twisting nature and displayed a nonpotential character. As an example, we covered fields of the coupled sunspots (CS's) being the source of solar flares. To make the results physically more transparent we passed from the flavor basis to the new one in which the resonance conditions did not depend on the angle θ_{23} while the θ_{23}-dependence had been transported on the resonance widths and the oscillation lengths.

In the Sun's conditions the possible resonance conversions of the active neutrinos were examined. In spite of the similar behavior of the neutrino beam in the Majorana and Dirac pictures there was a principal difference between these cases. It lies in the fact that in the Dirac neutrino case all magnetic-induced resonances transfer active neutrinos into sterile ones while in the Majorana neutrino case we deal with active neutrinos only. So, if the neutrino exhibits the Majorana nature, then the solar electron neutrino flux traveling through the region of the CS's can be converted into the active right-handed neutrinos $(\overline{\nu}_{eR}, \overline{\nu}_{\mu R}, \overline{\nu}_{\tau R})$. As an example, the emergence of $\overline{\nu}_{eL}$ neutrinos can be recorded by terrestrial detectors through the inverse β-decay reaction

$$\overline{\nu}_{eL} + p \rightarrow n + e^{+}, \tag{51}$$

having a threshold E_{th} = 1.8 MeV. Note that this reaction is at the heart of antineutrino detectors used for nuclear reactor monitoring on-line. On the other hand, since in the Dirac neutrino case the magnetic-induced resonances convert ν_{eL} neutrinos into sterile ν_{lR} neutrinos, then only decreasing the number of ν_{eL} can be observed when the solar neutrino flux passes the CS's region. As regards the phenomena of depleting the solar electron neutrino flux, the observation of decreasing the β-decay rates of some elements during the SF's [29–31] may be speculated to be its experimental confirmation. It should be stressed that the existence of the connection between the intensity of the solar neutrino flux and the β-decay rates must be supported by other experiments.

We have also demonstrated that the expressions for the survival probability of electron neutrinos found in the three neutrino generations convert into the well known expressions of the two FA provided $\phi = \psi = 0$.

In summary, we emphasize that the investigation of neutrino fluxes emitted from stellar objects will enable us to obtain information not only about such neutrino properties as multipole moment values and their nature (Dirac or Majorana) but about the stellar object structure too.

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