Neutrinos under Strong Magnetic Fields

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Abstract

In this talk we review the results on neutrino propagation under external magnetic fields. We concentrate on the effects of strong magnetic fields \( M^2_W \gg B \gg m_e^2 \) in neutral media. It is shown that the neutrino energy density get a magnetic contribution in the strong-field, one-loop approximation, which is linear in the Fermi coupling constant as in the charged medium. It is analyzed how this correction produces a significant oscillation resonance between electron-neutrinos and the other two active flavors, as well as with sterile neutrinos. The found resonant level-crossing condition is highly anisotropic. Possible cosmological applications are discussed. Effects due to primordial hypermagnetic fields on neutrinos propagating in the symmetric phase of the electroweak model are also presented. At sufficiently strong hypermagnetic fields, \( B \geq T^2 \), the neutrino energy is found to be similar to that of a massless charged particle with one-degree of freedom.

1 Introduction

A wide spectrum of strong magnetic fields can be found and predicted in nature. Magnetic fields can be as strong as \( 10^{12} \sim 10^{13} \) Gauss in the surface of typical radio pulsars [1], an even stronger \( (10^{14} \sim 10^{15} \) Gauss) in magnetars (with interior fields that may range up to \( 10^{16} \sim 10^{17} \) Gauss) [2]. Superconducting magnetic strings, if created after inflation, could generate fields of \( 10^{30} \) Gauss in their vicinity [3]. Primordial magnetic fields of \( 10^{24} \) Gauss at the electroweak (EW) scale have been proposed as the possible origin of the seed field needed to generate through galactic dynamo effect the large-scale magnetic fields observed in a number of galaxies, and galaxy clusters [4].

Strong magnetic fields can affect matter since they confine electrons perpendicular to its direction and consequently increase the atom binding energies. This effect can create matter bound states as molecular chains, magnetized three-dimensional condensed matter, etc. [5]. On the other hand, strong
enough magnetic fields have also important quantum electrodynamics effects, as the modification of the dielectric property of the medium, the polarization of photon modes [3] and the nonlinear photon splitting [4].

In this talk we present the effect of strong magnetic fields on the neutrino energy spectrum [8], [9] and discuss its consequences for neutrino oscillations. Although the neutrino, being a neutral particle, cannot directly interact with the magnetic field, we know that its propagation can be modified in the presence of an external field through quantum corrections. As shown below, in the presence of a strong magnetic field ($M_W^2 \gg B \gg m_e^2$, where $M_W$ and $m_e$ are the W-boson and electron masses respectively) the modification of the neutrino energy, due to one-loop corrections of the self-energy operator at finite temperature, gives rise to a resonant level-crossing condition in neutrino oscillations that is linear in the Fermi coupling constant. As known, in the weak-field approximation linear order modifications of the neutrino energy only appears in charged media [10] (at zero field this is the well known MSW effect [11]).

Taking into account that the particle-antiparticle asymmetry of the universe is at the level of $10^{-10} - 10^{-9}$, the cosmological medium can be considered neutral, therefore for primordial neutrino oscillations the MSW effect can be disregarded. Nevertheless, if a strong primordial magnetic field was present during the neutrino decoupling era, the results we will discuss can be significant for cosmology, and specifically for primordial nucleosynthesis.

In this talk we will also report some recent results [12] on neutrino propagation in a constant hypermagnetic field. Effects of different nature in the presence of primordial hypermagnetic fields have been also considered by several authors [13]. If a primordial magnetic field existed prior to the EW phase transition, only its $U(1)$ gauge component, i.e. the hypermagnetic field, would penetrate the EW plasma for infinitely long times. Its non-Abelian component would decay because of its infrared magnetic mass $\sim g^2 T$, which is generated non-perturbatively through the non-linear interactions of the non-Abelian fields in the thermal bath [14]. On the other hand, for Abelian and non-Abelian electric fields a Debye screening will be always generated by thermal effects producing a short-range decay for both fields [15]. Hence, the only large scale primordial field that can penetrate the EW symmetric phase is the hypermagnetic field.

By studying the finite-temperature neutrino dispersion relations in the chiral phase of the EW model in the presence of a strong hypermagnetic field, we will show that the hypermagnetic field counteracts the thermal effect responsible for the creation of effective masses for the chiral leptons [16]. As a consequence, in the strong field approximation, where leptons are basically constrained to the lower Landau level (LLL), we find that neutrinos behave as massless particles with an anisotropic propagation.

2 Neutrino Self-Energy in Strong Magnetic Field

To obtain the quantum correction to neutrino energy in a magnetized thermal bath, we should calculate the neutrino self-energy in the presence of a magnetic
field at finite temperature. As it is known, the neutrino self-energy operator is a Lorentz scalar that can be formed in the spinorial space taking the contractions with all the independent elements of the Dirac ring. From its explicit chirality it reduces to

\[ \sum (p) = R \sum (p)L, \quad \sum (p) = V_\mu \gamma^\mu \]  

(1)

where \( L, R = \frac{1}{2}(1 \pm \gamma_5) \) are the chiral-projector operators, and \( V_\mu \) is a Lorentz vector that in covariant notation can be given as a superposition of four basic vectors formed from the characteristic tensors of the corresponding problem. In the present case

\[ \sum (p) = a \hat{p}_\parallel + b \hat{p}_\perp + cp^\mu \hat{F}_{\mu\nu} \gamma^\nu + idp^\mu \hat{F}_{\mu\nu} \gamma^\nu. \]  

(2)

The presence of the magnetic field, given through the dimensionless magnetic field tensor \( \hat{F}_{\mu\nu} \) and its dual \( \hat{\bar{F}}_{\mu\nu} \), allows the covariant separation in (2) between longitudinal and transverse momentum terms that naturally appears in magnetic backgrounds

\[ \hat{p}_\parallel = p^\mu \hat{F}_{\mu\nu} \hat{\bar{F}}_{\mu\nu} \gamma^\nu, \quad \hat{p}_\perp = p^\mu \hat{F}_{\mu\nu} \hat{\bar{F}}_{\mu\nu} \gamma^\nu. \]  

(3)

The coefficients \( a, b, c, \) and \( d \) are Lorentz scalars that depend on the parameters of the theory and the approximation used. We are interested in one-loop corrections, thus for a neutral medium (i.e. in the absence of chemical potentials) the leading contribution is given by the bubble diagram with internal lines of virtual electrons and W-bosons. Since both virtual particles are electrically charged, the magnetic field interacts with both of them producing the Landau quantization of the corresponding transverse momenta. Thus, we end up with two set of Landau quantum numbers \([8]\), one for the electron, and other for the W-boson.

We assume a strong-field approximation, \( M_H^2 \gg B \gg m_e^2 \). Since in this case the gap between the electron Landau levels is larger than the electron mass square, it is consistent to use the LLL approximation for the electron, while for the W-boson it is obvious that we must sum in all W-boson Landau levels.

To justify such an approximation for cosmological applications we should recall that due to the equipartition principle, the magnetic energy can only be a small fraction of the universe energy density. This argument leads to the relation between field and temperature \( B/T^2 \sim 2 \). For such fields, the effective gap between the Landau levels (LL) is \( eB/T^2 \sim O(1) \). In this case the weak-field approximation (where the sum in all LL is important), cannot be used because field and temperature are comparable. On the other hand, because the thermal energy is of the same order of the energy gap between LL’s, it is barely enough to induce the occupation of just a few of the lower electron LL’s, since, as we are considering, the electron mass is much smaller than the magnetic field. Therefore, it is natural to expect that the LLL approximation in the electron spectrum will provide a good qualitative description of the neutrino propagation
in the presence of strong magnetic fields, although clearly a more quantitative treatment of the problem in a field $B \sim 2 T^2$ would require numerical calculations due to the lack of a leading parameter.

In the leading order in the expansion in powers of the Fermi coupling constant the scalar coefficients of Eq. (2) are obtained in our approximation as

$$a = -c \simeq \frac{g^2 e B}{8 \pi M^2_W} \left[ -\frac{T^2}{3M^2_W} \right] \exp(-p^2_/eB), \quad b = d \simeq 0$$  \hspace{1cm} (4)

Notice that by comparison the thermal contribution is smaller in a factor of $1/M^2_W$ with respect to the field-dependent vacuum contribution. Then, in the strong-field limit the thermal contribution has the same second order in the Fermi coupling constant as it is in the zero-[17], and weak-field [15] cases. On the other hand, the self-energy field-dependent vacuum contribution in (4) is of the same order in the Fermi coupling constant as the one found in a charged medium at zero [11] and weak [10] fields.

Using the zero-temperature weak-field results of Ref. [19] to identify the scalar coefficients of the general structure (2) in that approximation, we have that in the weak-field limit they are given by

$$a = b \simeq 0, \quad c = -\frac{6}{4} d \simeq \frac{6g^2 e B}{(4\pi)^2 M^2_W}$$  \hspace{1cm} (5)

We see that at weak field, the neutrino self-energy has also linear contributions in the Fermi coupling constant, but they are associated with different structures in (2) as compared with the result in the strong-field limit (4). The role of the different self-energy structural members into the neutrino energy spectrum will be clear in the next section. There, we will show that the results in the strong-field limit [11] produce a magnetic field dependence in the neutrino energy which is linear in the Fermi coupling constant, while the weak-field results [10] produce a smaller second-order contribution.

3 Neutrino Energy Spectrum and Index of Refraction

The neutrino field equation of motion is

$$[\hat{p} - \sum(p)]\Psi_L = 0$$  \hspace{1cm} (6)

The dispersion relation is obtained by solving Eq. (6), or equivalently, by finding the nontrivial solution of Eq. (6) through the equation

$$\det[\hat{p} - \sum(p)] = 0$$  \hspace{1cm} (7)

where $\sum(p)$ is given in its general covariant form by Eq. (2).

In the strong-field limit [11], the solution of Eq. (7) is [8].
\[ E_p = \pm \left| \vec{p} + \sqrt{2a} (\vec{p} \times \vec{B}) \right| = \pm |\vec{p}| [1 + 2a \sin^2 \alpha] \] (8)

where \( \alpha \) is the angle between the direction of the neutrino momentum and that of the applied magnetic field.

To obtain the neutrino index of refraction \( n \), we substitute (8) into the formula

\[ n \equiv \frac{|\vec{p}|}{E_p} \] (9)

to find

\[ n \simeq 1 - a \sin^2 \alpha \] (10)

From Eqs. (8) and (10) we can see that neutrinos moving with different directions in the magnetized space will have different dispersion relations and consequently, different indexes of refraction. It is interesting to notice that, although neutrinos are electrically neutral particles, the magnetic field, through quantum corrections, can produce anisotropic neutrino propagation, having a maximum effect for neutrinos propagating perpendicularly to the magnetic field (see Eq. (10)). The order of the energy correction is \( g^2 |\vec{eB}| M_W^2 \), and so is the order of the asymmetry.

If we consider the weak-field results \( \delta \) in the dispersion relation \( \delta \) we obtain

\[ E'_p = \pm |\vec{p}| [1 + \frac{5}{18} c^2 \sin^2 \alpha] \] (11)

Here, as the energy depends on \( c^2 \), we can see that the weak field produces a negligible second order in the Fermi coupling constant expansion. It can be corroborated that the inclusion of temperature in this approximation also produces a second order correction \( \delta \).

An asymmetric neutrino propagation of linear order in the Fermi coupling constant expansion was previously found \( \delta \) for weak magnetic fields but in a charged medium \( (\mu \neq 0) \). There, the energy correction was given by

\[ E''_p = a' \pm \left| \vec{p} - b' \vec{B} \right| \] (12)

where the coefficients \( a' \) and \( b' \) are proportional to the electron/positron number densities \( n_\pm \), and electron/positron distribution functions \( f_\pm \) respectively

\[ a' = \frac{g^2}{4M_W^2} (n_- - n_+), \quad b' = \frac{eg^2}{2M_W^2} \int \frac{d^3p}{(2\pi)^3 2E} \frac{dE}{dE} (f_- - f_+) \] (13)

If we compare our results for strong magnetic fields in a neutral medium \( \delta \) with those for weak field in a charged medium \( \delta \), we see that while in the first
case the maximum field effects occur for neutrinos propagating perpendicularly to the field direction, in the second case this propagation modes are not affected by the field, but on the contrary, the maximum effect takes place for neutrinos propagating along the field lines. Moreover, the asymmetric term in the dispersion relation found in Ref. [10] changes its sign when the neutrino reverses its motion. This property is crucial for a possible explanation of the peculiar high pulsar velocities [20]. Nevertheless, in [8] we find that in our approximation the neutrino energy-momentum relation is invariant under the change of $\alpha$ by $-\alpha$. Finally, the dispersion relations [12] have different values for neutrinos and antineutrinos respectively, since in a charged medium the CP-symmetry is violated. In our result [3], particle and antiparticle have the same energy, as it should be in a neutral medium.

4 Neutrino Oscillations and Resonance in a Strongly Magnetized Neutral Medium

Assuming that the magnetic field strength is confined within the range $m_e^2 \ll B \ll m_\mu^2$, the strong-field approximation becomes valid for electron-neutrinos, but not for muon-neutrinos and tau-neutrinos. For the last two, the magnetic field will have the weak effect in the energy corrections discussed below Eq. (11) and therefore can be neglected.

Let us consider the evolution equation in the presence of a strong magnetic field for a two-level system

$$\frac{d}{dt} \begin{pmatrix} \nu_e \\ \nu_\mu \end{pmatrix} = H_B \begin{pmatrix} \nu_e \\ \nu_\mu \end{pmatrix}$$

where the Hamiltonian $H_B$ is given by

$$H_B = p + \frac{m_e^2 + m_\mu^2}{4p} + \begin{pmatrix} \frac{-\Delta m^2}{4p} \cos 2\theta + E_B & \frac{\Delta m^2}{4p} \sin 2\theta \\ \frac{\Delta m^2}{4p} \sin 2\theta & \frac{\Delta m^2}{4p} \cos 2\theta \end{pmatrix}$$

Here, $\theta$ is the vacuum mixing angle, $\Delta m^2 = m_\mu^2 - m_e^2$ is the mass square difference of the two mass eigenstates, and the magnetic energy density contribution to the electron-neutrino is

$$E_B = \frac{g^2 e B}{2(4\pi)^2 M_W^2} |\vec{p}| \sin^2 \alpha$$

In Eq. (15) the magnetic field contribution to the muon-neutrino in the second diagonal term has been neglected, taking into account that it will be of second order in the Fermi coupling constant as corresponds to the weak-field approximation.

To find the evolution Hamiltonian corresponding to the mass eigenstates in the magnetized space we need to transform $H_B$ according to
\( \tilde{H} = U_B^\dagger H_B U_B \) \hspace{1cm} (17)

with transformation matrix

\[
U_B = \begin{pmatrix}
\cos 2\theta_B & \sin 2\theta_B \\
-\sin 2\theta_B & \cos 2\theta_B
\end{pmatrix}
\] \hspace{1cm} (18)

depending on the new mixing angle \( \theta_B \) given through the relation

\[
\tan 2\theta_B = \frac{\tan 2\theta}{2\Delta m^2 p / \cos 2\theta - E_B}
\] \hspace{1cm} (19)

If we consider that the only flavor present at the initial time was the electron-neutrino, using the evolution equation (15) we find that the appearance probability for the muon neutrino is given by

\[
P_B (\nu_e \to \nu_\mu) = \sin^2 \theta_B \sin^2 \pi x / \lambda
\] \hspace{1cm} (20)

where \( \lambda \) is the oscillation length in the magnetized space

\[
\lambda = \frac{\lambda_0}{\sin^2 2\theta + (\cos 2\theta - \lambda_0 / \lambda_e)^2}^{1/2}
\] \hspace{1cm} (21)

written in terms of the vacuum (\( \lambda_0 \)) and magnetic (\( \lambda_e \)) oscillation lengths

\[
\lambda_0 = \frac{4\pi p}{\Delta m^2}, \quad \lambda_e = \frac{2\pi}{E_B}
\] \hspace{1cm} (22)

and

\[
\sin^2 2\theta_B = \frac{\sin^2 2\theta}{(\cos 2\theta - \lambda_0 / \lambda_e)^2 + \sin^2 2\theta}
\] \hspace{1cm} (23)

is the probability amplitude defined through the mixing angle \( \theta_B \) (19).

If the resonant condition defined through the mixing angle \( \theta_B \) (19).

\[
\frac{\lambda_0}{\lambda_e} = \cos 2\theta
\] \hspace{1cm} (24)

is satisfied, then the probability amplitude (23) will get its maximum value independently of the value of the mixing angle in vacuum \( \theta \). The condition (24) is a resonant level-crossing condition, and as usual, it can be also obtained by equating the two diagonal Hamiltonian elements in (15). At a magnetic field for which the condition (24) is satisfied, a maximum transmutation between the two flavors will occur (the same effect will be obtained between electron-neutrinos and tau or sterile neutrinos). We should notice that the resonant effect in the magnetized neutral medium will be anisotropic, depending on the direction of propagation of the electron-neutrino with respect to the magnetic field (see that \( \lambda_e \) depends on \( \alpha \)).
The resonant phenomenon here is similar to that in the well known MSW effect in a charged medium. Nevertheless, we should stress that in the magnetized neutral medium the oscillation process will not differentiate between neutrinos and antineutrinos, while in the charged medium, we have that if there is resonance for the neutrino there will be no resonance for the antineutrino and vice versa.

5 Consequences for Cosmology

The early Universe, unlike the dense stellar medium, is almost charge symmetric ($\mu = 0$), since, as already mentioned, the particle-antiparticle asymmetry in the Universe is believed to be at the level of $10^{-10} - 10^{-9}$, while in stellar material it is of order one. It is known that the contribution to the neutrino energy density of pure thermal effects $^{17}$ or of quantum corrections obtained in a weakly magnetized neutral medium $^{18, 19}$, are both of second order in the Fermi coupling constant, therefore negligible small. Nevertheless, as we have shown in $^{13}$, if sufficiently strong magnetic fields were present in the early Universe, they would give rise to corrections to the energy density that are linear in the Fermi coupling constant. These corrections can produce effects as significant as those associated to the MSW mechanism.

The existence of strong magnetic fields in the early Universe seems to be a very plausible idea $^{4, 21}$, as they may be required to explain the observed galactic magnetic fields, $B \sim 2 \times 10^{-6} \, G$ on scales of the order of 100 kpc $^{1, 22}$.

The strength of the primordial magnetic field in the neutrino decoupling era can be estimated from the constraints derived from the successful nucleosynthesis prediction of primordial $^4\text{He}$ abundance $^{23}$, as well as on the neutrino mass and oscillation limits $^{24}$. These constraints, together with the energy equipartition principle, lead to the relations

$$m_e^2 \leq eB \leq m_\mu^2, \quad B/T^2 \sim 2 \quad (25)$$

Then, it is reasonable to assume that between the QCD phase transition epoch and the end of nucleosynthesis a primordial magnetic field in the range given by relations $^{25}$ could have been present $^{3}$.

If such a field existed, it could significantly modify the $\nu_e \leftrightarrow \nu_\mu, \nu_\tau$ and $\nu_e \leftrightarrow \nu_s$ resonant oscillations in the way we have shown in this paper, and consequently affect primordial nucleosynthesis $^{24}$.

Another interesting question regarding the effect of strong primordial fields is how these fields would affect neutrino propagation prior to the EW phase transition in case they were originated at earlier times in the Universe evolution. Since in this phase only a hypermagnetic fields matters, as the non-Abelian component decays due to the acquired infrared mass, it is enough to investigate the propagation of neutrinos in the presence of a background hypermagnetic field.

Notice that there are essential differences between the interactions of neutrinos with hypermagnetic and magnetic fields. We recall that the neutrino,
being a neutral particle with respect to the electromagnetic group, has instead non-zero hypercharge and thus can interact with a hypermagnetic field already at the bare level.

In the symmetric phase of the EW model, fermions are in the chirally invariant phase with separated left-handed and right-handed representations of the gauge group. In that phase, a fermion mass term in the Lagrangian density is forbidden by the symmetry of the theory. However, as it was found in Ref. [16], finite temperature corrections induce a pole in the fermion Green’s function that plays the role of an effective mass. The induced effective “mass” modifies the fermion dispersion relation in the primeval plasma opening the door for possible cosmological consequences [16, 24].

As it was shown in [12] a strong hypermagnetic field can modify the neutrino self-energy at finite temperature, so at the one-loop level it is given by

$$\Sigma_{\nu_L}(p) = \left[ A\hat{p}_\| + B\hat{p}^\mu \widehat{H}_{\mu\nu}\gamma^\nu \right] L$$

with

$$A = -B = \frac{G_{\nu_L}^2}{2\pi} \left[ \frac{a}{4\pi} + \frac{T^2}{|g' H|} \right]$$

In (27), $G_{\nu_L}^2 = \left[ \frac{g'^2}{4} + \frac{3g^2}{4} \right]$, $a$ is a positive constant of order one ($a \simeq 0.855$), and $\widehat{H}_{\mu\nu}$ is the dual of the dimensionless hypermagnetic field tensor ($\widehat{H}_{\mu\nu} = H_{\mu\nu}/H$).

The neutrino dispersion relation in the presence of a constant hypermagnetic field including radiative corrections is given by

$$\det \left[ \vec{p} \cdot \gamma + \Sigma \right] = 0,$$

where the neutrino generalized momentum in the presence of the hypermagnetic field is $\vec{p}_{\nu_L} = (p_0, 0, -\text{sgn}(g'H)\sqrt{|g'H|n, p_3})$, with the integer $n = 0, 1, 2, ...$ labelling the neutrino Landau levels.

The fact that the two structures in (26) have the same coefficient (i.e. $A = -B$), together with the dependence of $\vec{p}_{\nu_L}$ in the strong-field approximation on its longitudinal components $p_0$ and $p_3$ only (i.e. in the LLL approximation ($n = 0$)), has significant consequences for the propagation of chiral leptons in the hypermagnetized medium, as one can immediately corroborate by explicitly solving the lepton dispersion relation (28) to obtain

$$-(1 + 2A)p_3^2 + (1 + 2A)p_3^2 = 0$$

This result indicates that in a strong hypermagnetic field neutrinos behave as massless particles, so the hypermagnetic field counteracts the temperature effects on the neutrino dispersion relation. As mentioned above, according to Ref. [16] the dispersion relations of the chiral leptons are modified due to high-temperature effects in such a way that a temperature-dependent pole appears
in their one-loop Green’s functions (i.e. an effective mass). As we have shown here, if a primordial hypermagnetic field could exist in the EW plasma prior to the symmetry-breaking phase transition, with such a strength that the leptons would be mainly confined to the LLL for the existing temperatures, then the thermal effective mass found for the leptons in that phase at zero field \[16\] will be swept away. This effect can be of interest for cosmology, since it will alter the behavior of the lepton masses with temperature during the EW phase transition.

Another important outcome of the dispersion relation (29) is the large anisotropy of the neutrino propagation in strong hypermagnetic field domains. This is the consequence of the degeneracy in the energy, which does not depend on the tranverse momentum components. No degeneracy appears in the case of neutrinos propagating in magnetic fields (see Eq. 8). This different behavior can be understood from the fact that the neutrino is minimally coupled to the hypermagnetic field because of its nonzero hypercharge, but due to its electrical neutrality, it can couple to the magnetic field only through radiative corrections.

It is worth to notice that the leading contribution to the one-loop correction of the neutrino self-energy at small momentum gives rise to a dispersion relation that is formally the same as the bare dispersion relation in the strong hypermagnetic field (it is the typical dispersion relation of a massless charged particle in a magnetic field strong enough to confine the particle to its LLL).

From the above results, one can envision that if large primordial fields were present in the early universe, the neutrino propagation would be highly anisotropic, an effect that could leave a footprint in a yet to be detected relic neutrino cosmic background. If such a footprint were observed, it would provide a direct experimental proof for the existence of strong primordial fields.

6 Concluding Remarks

In this talk we have studied, within the framework of the standard model, the quantum effects on neutrinos propagating in a strong magnetic field. For fields in the range \(m_\mu^2 \gg B \gg m_e^2\), we found a field-dependent contribution to the electron-neutrino energy density that is linear in the Fermi coupling constant. For the other neutrino flavors, the strong-field approximation is not valid and consequently the magnetic corrections to the energy density are negligible small.

This separation between the magnetic contribution to the energy density of different flavors directly affects neutrino oscillations leading to a level-crossing resonance that depends on the value of the applied magnetic field. This resonant effect is similar to the well known MSW effect, but with the difference that it can take place in a neutral medium, and that it has an anisotropic character, depending on the direction of the neutrino propagation with respect to the magnetic field.

The fact that a resonant oscillation can take place in a neutral medium at the leading order in the Fermi coupling constant has significant interest for cosmology in case that a sufficiently strong primordial field were present during
the neutrino decoupling era, since it would affect primordial nucleosynthesis.

On the other hand, in the symmetric phase of the EW theory the effect of a primordial field can only be carried out by the hypermagnetic field. In the presence of this field, neutrinos behave as charged particles, and in the strong field approximation, when leptons are confined to their LLL, the leading contribution to the energy density of the neutrino thermal bath is similar to that of a massless charged particle with only one degree of freedom (the energy only depends on the longitudinal momentum). That is, the strong hypermagnetic field swept away the thermal mass found at zero field and high temperature [10]. This result can be important for the study of the EW phase transition at finite temperature and in the presence of strong primordial fields.

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References
[1] P. P. Kronberg, Rep. Prog. Phys., 57, 325 (1994); R. Beck et. al., Ann. Rev. Astron. Astrophys. 34, 153 (1996).
[2] R. C. Duncan and C. Thompson, Astrophys. J. 392, L9 (1992); B. Paczynski, Acta Astron. 42, 145 (1992); V. V. Usov, Nature 357, 472 (1992); C. Thompson and R. C. Duncan Astrophys. J. 473, 322 (1996).
[3] J. P. Ostriker, C. Thompson and E. Witten, Phys. Lett. B 180, 231 (1986).
[4] D. Grasso and H.R. Rubinstein, Phys. Rep. 348, 163 (2001).
[5] For a recent review see D. Lai, Rev. Mod. Phys. 73, 629 (2001).
[6] H. Perez-Rojas ans A. E. Shabad, Ann. of Phys. 121, 432 (1979); 138, 1 (1982); H. Perez-Rojas, Zh. Eksp. Teor. Fiz. 76, 1 (1979); A. K. Ganguly, S. Konar and P. B. Pal, Phys. Rev. D 60, 105014 (1999); J. C. D’Olivo, J. F. Nieves and S. Sahu, Phys. Rev. D 67, 025018 (2003).
[7] S. L. Adler, Ann. Phys. 67, 599 (1971); V. O. Papanyan and V. I. Ritus, Zh. Eksp. Teor. Fiz. 61, 2231 (1971) (Sov. Phys. JETP 34, 1195 (1972)); ibid 65, 1756 (1973), (Sov. Phys. JETP 38, 879 (1974)); V. N. Baier, A. I. Milstein, and R. Z. Shaisultanov, Phys. Rev. Lett. 77, 1691 (1996); J. I. Weise, M. G. Baring, and D. B. Melrose Phys. Rev D57, 5526 (1998); A. V. Kuznetsov, N. V. Mikheev, and M. V. Chistyakov, Phys. of Atom. Nuc. 62, 1638 (1999).
[8] E. Elizalde, E. J. Ferrer and V. de la Incera, Ann. of Phys. 295, 33 (2002); E. J. Ferrer, "Non-Perturbative Effects of Strong Magnetic Fields" in Quantization, Gauge Theory, and Strings, edited by A. Semikhatov, M. Vasiliev and V. Zaikin, Proc. of the International Conference dedicated to the memory of Prof. Efim Fradkin, Sci. World Pub. Co. 2001, pp. 301-306.

[9] E. Elizalde, E. J. Ferrer and V. de la Incera, Neutrino Propagation in a Strongly Magnetized Medium, in preparation.

[10] J. C. D’Olivo, J. F. Nieves and P. B. Pal, Phys. Rev. D40, 3679 (1989).

[11] L. Wolfenstein, Phys. Rev. D17, 2369 (1978); S. P. Mkheyev, A. Yu. Smirnov, Sov. J. Nucl. Phys. 42, 913 (1985).

[12] J. Cannellos, E. J. Ferrer and V. de la Incera, Phys. Lett. B542 123 (2002).

[13] M. Giovannini and M. E. Shaposhnikov, Phys. Rev. D57, 2186 (1998); M. Giovannini Phys. Rev. D61, 063004 (2000); ibid 063502; A. Ayala, J. Besprosvany, G. Pallares and G. Piccinelli, Phys. Rev. D64, 123529 (2001); A. Ayala, G. Piccinelli and G. Pallares, Phys. Rev. D66, 1103503 (2002); A. Ayala, and J. Besprosvany, Nucl. Phys. B651, 211 (2003).

[14] A.D. Linde, Rep. Prog. Phys., 42, 389 (1979); Phys. Lett B396, 178 (1980); K. Kajantie, M. Laine, K. Rummukainen, and M. Shaposhnikov, Nucl. Phys. B493, 413 (1997).

[15] E.S. Fradkin, Proceedings Lebedev Phys. Inst., Vol. 29, 7 (1965) (Eng. Transl., Consultant Bureau, New York (1967)); A. K. Rebhan, Phys. Rev. D48, R3967 (1993); Nucl. Phys. B430, 319 (1994); E. Braaten, and A. Nieto, Phys. Rev. Lett 73, 2402 (1994); R. Baier, and O.K. Kalashnikov, Phys. Lett B328, 450 (1994).

[16] H.A. Weldon, Phys. Rev. D26, 2789 (1982).

[17] D. Notzold and G. Raffelt, Nucl. Phys. B307, 924 (1988).

[18] P. Elmfors, D. Grasso, G. Raffelt, Nucl. Phys. B479, 3 (1996).

[19] G. McKeon, Phys. Rev. D24, 2744 (1981).

[20] A. Kusenko and G. Segre, Phys. Rev. Lett. 77, 4872 (1996).

[21] K. Enquist, Int. J. Mod. Phys. D7, 331 (1998).

[22] Y. Sofue, M. Fujimoto, and R. Wielebinski, Ann. Rev. Astron. Astrophys. 24, 459 (1986); R. Beck et. al., Ann. Rev. Astron. Astrophys. 34, 153 (1996).

[23] K. A. Olive, D. N. Schramm, G. Steigman, and T. Walker, Phys. Lett. B 236, 454 (1990); L. M. Krauss and P. Romanelli, Astrophys. J. 358, 47 (1990).
[24] K. Enqvist, V. Semikoz, A. Shukurov, and D. Sokoloff, Phys. Rev. D48, 4557 (1993).

[25] A. D. Dolgov, Phys. Rep. 370, 333 (2002).