Bounds on neutrino transition magnetic moments in random magnetic fields

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Abstract

We consider the conversions of active to sterile Majorana neutrinos ν_a and ν_s, due to neutrino transition magnetic moments in the presence of random magnetic fields (r.m.f.) generated at the electroweak phase transition. From a simple Schrödinger-type evolution equation, we derive a stringent constraint on the corresponding transition magnetic moments and display it as a function of the domain size and field geometry. For typical parameter choices one gets limits much stronger than usually derived from stellar energy loss considerations. These bounds are consistent with the hypothesis of seeding of galactic magnetic fields by primordial fields surviving past the re-combination epoch. We also obtain a bound on active-sterile neutrino transition magnetic moments from supernova energy loss arguments. For r.m.f. strengths in the range 10^7 to 10^{12} Gauss we obtain limits varying from $\mu_{as}^\nu \lesssim 10^{-13}\mu_B$ to $\mu_{as}^\nu \lesssim 10^{-18}\mu_B$, again much stronger than in the case without magnetic fields.

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1. Introduction

Recently there has been a renewed interest concerning neutrino propagation in media with random magnetic fields, both from the point of view of the early universe hot plasma as well as astrophysics [1]. It has been shown that random magnetic fields can strongly influence neutrino conversion rates and this could have important cosmological and astrophysical implications, especially in the case of conversions involving a light sterile neutrino ν_s. Indeed, note that the present hints from solar and atmospheric neutrino observations [2,3] as well as from the COBE data on cosmic background temperature anisotropies on large scales [4,5] indicate that a light sterile neutrino ν_s [6,7] might exist in nature.

So far the most stringent constraints for the neutrino mass matrix including a sterile neutrino species, ν_s, are obtained from the nucleosynthesis bound on the maximum number of extra neutrino species that can reach thermal equilibrium before nucleosynthesis and change the primordially produced helium abundance [8]. This has been widely discussed in the case of vanishing magnetic field [9] and for the case of a large random magnetic field that could arise from the electroweak phase transition and act as seed for the galactic magnetic fields [10].

The effect of active-sterile neutrino conversions in a supernova has also been discussed, both in the case where no magnetic field is present [11], as well as in the presence of random magnetic field as large as 10^{16} Gauss [12], following a suggestion made in ref. [13].

In this paper we focus on the important effect of relatively small random magnetic fields on the active sterile neutrino conversion rates when nonzero Majorana neutrino transition magnetic moments are taken into account [14,15]. We apply this to the case of ν_a to ν_s conversions in the early universe hot plasma, as well as in a supernova, showing how their effect can place limits that are substantially more stringent than those that apply in the absence of a magnetic field. First we focus on the nucleosynthesis constraint on active to sterile transition magnetic moments, and display this constraint as a function of the domain size and field geometry. For typical choices we obtain $\mu_{as}^\nu \lesssim 10^{-15}\mu_B$, which is much stronger than usually derived from stellar energy loss considerations neglecting r.m.f. effects. These bounds are consistent with the hypothesis of seeding of galactic magnetic fields by primordial fields generated at the electroweak phase transition. We also obtain a bound on active-sterile neutrino transition magnetic moments from supernova energy loss arguments. For modest r.m.f. field strengths in the range 10^7 to 10^{12} Gauss we obtain limits that vary from $\mu_{as}^\nu \lesssim 10^{-13}\mu_B$ to $\mu_{as}^\nu \lesssim 10^{-18}\mu_B$, again much stronger than in the case without magnetic fields [16].

2. Neutrino conversions

Let us consider the Schrödinger equation describing a system of two neutrinos, one active and one sterile, ν_a and ν_s in a plasma with a random magnetic field $\mathbf{B}(t)$. In general one has a four-dimensional system of equations [14]. We will use the ultra-relativistic limit and will neglect the corresponding mixing angle, $s = 0$, $c = 1$, in which case one may write [17]

$$i\frac{d}{dt}\begin{pmatrix} \nu_a \\ \nu_s \end{pmatrix} = \begin{pmatrix} V_a - \Delta + \mu_{eff}B_{\parallel}(t) & \mu B_{\perp}(t) \\ \mu B_{\perp}(t) & 0 \end{pmatrix}\begin{pmatrix} \nu_a \\ \nu_s \end{pmatrix}. \tag{2.1}$$

Here V_a (for ν_a, $a = e, \mu, \tau$) is the active neutrino vector interaction potential. For instance, for electron left-handed neutrinos propagating in the early universe hot plasma one has [18],

$$V = 3.45 \times 10^{-20}\Big(\frac{T}{MeV}\Big)^5 MeV, \tag{2.2}$$

while for the case of a supernova one has

$$V \simeq 4 \times 10^{-6}\rho_{14}(3Y_e + 4Y_{\nu_e} + 2Y_{\nu_\mu} - 1)MeV. \tag{2.3}$$

where ρ_{14} is the core density in units of 10^{14} g/cm^3 .

In the case of the muon left-handed neutrino the last abundance factor in eq. (2.3) is changed to $f(Y) = Y_e - 1 + 4Y_{\nu_\mu} + 2Y_{\nu_e}$, where we took into account ν-ν forward scattering amplitude in a supernova matter during the main neutrino burst time, about 10-20 seconds.

The term $\mu_{eff}B_{\parallel}$ is produced by the mean axial vector current of charged leptons in an external magnetic field $B_{\parallel} = \mathbf{B}.\mathbf{q}/q$. In the case of the early universe hot plasma the coefficient μ_{eff} is given as

$$\mu_{eff} = -12c_A \times 10^{-13}\mu_B\Big(\frac{T}{MeV}\Big), \tag{2.4}$$

while for the case of an ultra-relativistic degenerate electron gas of a supernova one has

$$\mu_{eff} = -8.6c_A \times 10^{-13}\mu_B(p_{F_e}/MeV) . \tag{2.5}$$

Here $c_A = \mp 0.5$ is the axial constant (upper sign for electron neutrino, lower one for $\nu_{\mu,\tau}$).

Note that even though the quantity μ_{eff} has the dimensions of a magnetic moment it is not a real magnetic moment since it is helicity conserving. Indeed, the additional energy splitting term obtained in eq. (2.1) does not lead to any helicity change. In contrast, the off-diagonal entries of the Hamiltonian eq. (2.1) involve the presence of the real neutrino transition magnetic moment $\mu = \mu_{as}^{\nu}$ as well as the transversal magnetic field component, $\mathbf{B}_{\perp}$.

Note also that our initial equation eq. (2.1) is quite general, since it applies to the case of Majorana neutrinos. It remains valid also for the description of the active to sterile neutrino conversions in the limit where these form a Dirac neutrino, in which case $\Delta = (m_2^2 - m_1^2)/2q = 0$ and μ becomes the usual magnetic moment.

Denoting $\mu B_\perp(t) = \tilde{H}_\perp(t)$ and using the auxiliary functions $R = Re(\langle \nu_e^* \nu_s \rangle)$ and $I = Im(\langle \nu_e^* \nu_s \rangle)$, we obtain from eq. (2.1) the standard system of first order differential equations,

$$\dot{P}(t) = -2\tilde{H}_\perp(t)I(t),$$

$$\dot{I}(t) = [V - \Delta + \mu_{eff}B_\|(t)]R(t) + \tilde{H}_\perp(t)(2P(t) - 1),$$

$$\dot{R}(t) = -[V - \Delta + \mu_{eff}B_\|(t)]I(t),$$

from which we derive the integro-differential equation for the neutrino conversion probability $P_{\nu_e \to \nu_s}(t) \equiv P(t)$

$$\dot{P}(t) = -2\int_0^t \tilde{H}_\perp(t)\tilde{H}_\perp(t_1)\cos\left(\int_{t_1}^t V(t_2)dt_2\right)[2P(t_1) - 1]dt_1, \qquad (2.6)$$

where $\int_{t_1}^t V(t_2)dt_2 = (V - \Delta)(t - t_1) + \mu_{eff}\int_{t_1}^t B_\|(t_2)dt_2$.

For r.m.f. domain sizes much smaller than the typical neutrino conversion length $l_{conv} \sim \Gamma^{-1}$, i.e. $L_0 \ll l_{conv}$ (see below eq. (2.11)) one can average this equation over the random magnetic field distribution. The averaged probability $\langle P(t) \rangle = \mathcal{P}(\sqcup)$ must depend only on even powers of the random field, since $\langle (B_j(t))^{2n+1} \rangle = 0$.

In what follows we neglect neutrino collisions with dense matter. One can show that averaging of the neutrino conversion probability over fast collisions does not lead to essential change of our results [3].

For uncorrelated random magnetic field domains,

$$\langle B_i(\vec{x})B_j(\vec{y}) \rangle = \frac{1}{2\lambda}\delta_{ij}\delta^{(3)}(\vec{x} - \vec{y}), \qquad (2.7)$$

we obtain the mean values

$$\langle B_\|(t) \rangle = \langle B_\perp(t) \rangle = \langle B_\|(t)B_\perp(t_1) \rangle = 0,$$

$$\langle B_\|(t)B_\|(t_1) \rangle = \langle B_\|^2 \rangle L_0 \delta(t - t_1), \qquad (2.8)$$

$$\langle B_\perp(t)B_\perp(t_1) \rangle = \langle B_\perp^2 \rangle L_0 \delta(t - t_1),$$

The correlation length λ is determined by the domain size L_0 and by the value of root mean squared field $B \equiv B_{rms} \equiv \sqrt{\langle B^2 \rangle}$ at the horizon scale $L = l_H$ [19]:

$$\frac{1}{\lambda} = \frac{3}{\pi(3 - 2p)}B_{rms}^2(l_H)L_0^3,$$

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[3] Moreover, in [19] it was shown that in an appropriate kinetic approach for the neutrino spin evolution in uncorrelated random magnetic fields the final result does not depend on the collision frequency at all

for $p \neq 3/2$, and

$$\frac{1}{\lambda} = \frac{3}{\pi} \ln \frac{l_H}{L_0} B_{rms}^2(l_H) L_0^3,$$

for $p = 3/2$.

In eq. (2.8) the mean squared fields $\langle B_\parallel^2 \rangle = \langle B^2 \rangle/3$, $\langle B_\perp^2 \rangle = 2\langle B^2 \rangle/3$ are given by the r.m.s. field $B \equiv \sqrt{\langle B^2 \rangle}$.

Averaging the master equation eq. (2.6) over a random magnetic field distribution with the use of eq. (2.8) we easily obtain the simple differential equation

$$\dot{\mathcal{P}} + \Gamma_\perp \mathcal{P} = \frac{\bar{}_\perp}{\in}, \tag{2.9}$$

which has the solution

$$\mathcal{P} = \frac{\infty}{\in}(\infty - \exp(-\bar{}_\perp \sqcup)) \approx \frac{\bar{}_\perp \sqcup}{\in}, \tag{2.10}$$

where the magnetic field damping parameter $\Gamma_\perp$ is defined as

$$\Gamma_\perp = 4\mu^2 \langle B_\perp^2 \rangle L_0 = \frac{8}{3}\mu^2 B^2 L_0. \tag{2.11}$$

Note that the result eq. (2.10) obtained from our Schrödinger equation eq. (2.1) coincides with the analogous ones obtained in [19] for the Dirac neutrino magnetic moment, using neutrino spin kinetic equation [20], or the Redfield equation for the density matrix describing a nuclear spin interacting with the lattice vibrations in solids [21].

Note also that in the simplified case of negligible neutrino mixing considered here, the assumed δ-correlated form of the random magnetic field implies that there is no dependence of the averaged neutrino conversion probability on the neutrino energy splitting due to the matter term V, to the term $\mu_{eff} B_\parallel$ or to the mixing parameter Δ.

3. Nucleosynthesis constraint

The rate for producing sterile neutrinos in the early universe hot plasma is given as the product of a typical weak interaction rate Γ_W times our averaged conversion probability $\mathcal{P}$. The constraint that follows from nucleosynthesis may be simply estimated as

$$\Gamma_s = \Gamma_W \mathcal{P} \lesssim \mathcal{H}, \tag{3.1}$$

where $\Gamma_W = 4.0 G_F^2 T^5$ is the rate for producing the standard model active neutrinos, $H = 4.46 \times 10^{-22}(T/MeV)^2 MeV$ is the Hubble parameter, and $\mathcal{P}$ is given by eq. (2.10).

Different mechanisms of the magnetic field generation in the early universe give the following main phenomenological formula

$$B = B_0 \Big(\frac{T}{T_0}\Big)^2 \Big(\frac{L_0}{l_H}\Big)^p, \tag{3.2}$$

where we put the maximal scale is chosen as the horizon scale [4] $L = l_H(t)$.

Note that if we put in eq. (3.2) $B_0 \sim T_0^2$ (as in Vachaspati's mechanism, where $B_0 \sim T_{EW}^2$ [22]) the r.m.s. field in eq. (2.10) depends only on the temperature of the universe T ($T \simeq T_{QCD}$ in eq. (3.3)) and on two parameters describing the geometry of the r.m.f. geometry: (a) the scaling parameter p, and (b) the size of random domain $L_0^{min} = a(MeV/T)$.

Using this equation we get the following constraint [5]

$$\mu \lesssim \frac{1.7 \times 10^{-21} \mu_B}{(L_0^{min})^{p+1/2} \Gamma_W^{1/2} H^p}, \tag{3.3}$$

where L_0^{min} is the minimal scale of the random magnetic field domains. Due to the copious production of left-handed neutrinos e.g. from pion decays after the quark hadron phase transition temperature $T \simeq T_{QCD} \simeq 200$ MeV this QCD temperature plays a crucial role in the estimate of the nucleosynthesis constraint. As a result all parameters in eq. (3.3) are evaluated at $T \simeq T_{QCD} \simeq 200$ MeV.

The allowed region of transition moments μ_ν we have obtained from nucleosynthesis is shown in Fig. 1 as a function of the scale parameter p and domain size L_0.

There is an independent bound from primordial nucleosynthesis on the strength of the random magnetic field. Indeed, it has been shown [23] that a too large r.m.f. strength would enhance the rates for the relevant weak processes at $T_{NS} \sim 0.1$ MeV, resulting in additional helium production. This leads to $B \lesssim 3 \times 10^{10}$ Gauss [23] which implies, at the present time [19],

$$B_{NS} \lesssim 1.8 \times 2.4^p \times 10^{-7-11p} \text{Gauss} \tag{3.4}$$

Clearly, all values of the parameter p between $p = 0$ (corresponding to uniform field) and $p = 3/2$ (corresponding to 3-dimensional elementary cells) can obey both nucleosynthesis limits on magnetic transition moments eq. (3.3) and on r.m.f. strength, eq. (3.4).

It is important to stress that, for a wide choice of p and L_0, our nucleosynthesis constraint on the active sterile transition moments is substantially stronger than the astrophysical one from supernova 1987A [16].

Until now our arguments are quite general. Now we show that our constraints are consistent with the hypothesis of a primordial origin of the observed galactic magnetic fields [22,24] which would require additional restrictions on the parameters p and L_0, resulting in a correspondingly stronger limit on the neutrino transition magnetic moments.

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[4]As explained in ref. [19], this follows from the definition of the parameter λ for the initial δ-correlator in eq. (2.7).

[5]Clearly in this approximation eq. (3.1) leads to the same constraint on the transition magnetic moment $\mu = \mu_{as}$ as obtained previously using kinetic theory for the case of Dirac neutrino diagonal magnetic moment [19].

This would require $L_0^{min} \gtrsim 10^3 (MeV/T)$ cm and $p \leq 1$. The first restriction comes from requiring that primordial fields survive against ohmical dissipation after recombination time and could seed the observed galactic magnetic field [23]. On the other hand, the limit $p \leq 1$ follows from eq. (3.4) and the lower dynamo theory astrophysical bound on the strength of the galactic magnetic field $B_{seed}(T_{now}) \gtrsim 10^{-18}$ Gauss [25]. The corresponding region is hatched in Fig. 1. This corresponds to an upper limit on the neutrino transition moment $\mu \lesssim 10^{-16} \mu_B$ (left corner of hatched region for $p = 1$). Moreover, in order to be a seed field for the Milky Way or Andromeda, one would require $B_{seed}(T_{now}) \gtrsim 10^{-13}$ Gauss [25] which corresponds to $p \lesssim 0.5$ [26], leading to a much tighter bound $\mu \lesssim 10^{-19} \mu_B$.

4. Supernova bounds

Now let us move to the case of a supernova with magnetic field. If this magnetic field is generated after collapse it could be viewed as a random superposition of many small dipoles of size $L_0 \sim 1$ km [13]. One can show that this hypothesis is consistent, even for large magnetic field strengths, with the non observation of gamma radiation from SN 1987A, if the magnetic field diffusion time [13]

$$t_{diff} \simeq 10^2 (B_0/10^9 G)^2 sec \qquad (4.1)$$

does not exceed the time necessary for the diffusion of the X-rays through the supernova mantle. This leads to an upper limit on the seed field B_0, $B_0 \lesssim 10^{12} G$. Thus, as long as this is fulfilled, the random magnetic field could well influence the SN 1987A neutrino burst without being observable at present via X-ray emission [6].

As we will show below, one can derive constraints on the active to sterile Majorana neutrino magnetic moments which are so stringent that they become relevant even for values of the r.m.s field strength as low as $B \lesssim 10^9$ Gauss which could be quite reasonable from the astrophysical point of view.

First, let us consider the strong small-scale random magnetic fields generated in a supernova as in [13]. For this case we can use the requirement that in the non-trapping regime the sterile neutrino can be emitted from anywhere inside the stellar core with a rate [12]

$$\frac{dQ}{dt} \simeq 2.8 \times 10^{55} \mathcal{P} \rho_{\infty\triangle} \mathcal{E}_{\infty''}^{\in} \frac{\mathcal{J}}{\int} \lesssim \infty\prime^{\triangle} \frac{\mathcal{J}}{\int}, \qquad (4.2)$$

where ρ_{14} is the core density in units of 10^{14} g/cm^3 and E_{100} is the neutrino energy in units of 100 MeV. In the following we will set these parameters to unity. The last

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[6] Alternatively, for the case of neutron stars the assumption of random magnetic field domains is consistent with their observed magnetic fields in the limit where the small domains merge together to form larger ones due to kinematic evolution of the random magnetic field energy transferred down the spectra $E(k) \sim k^{-5/3}$ [1] with the characteristic time $\tau(L_0) \sim L_0^{2/3}$

bound corresponds to the maximum observed integrated neutrino luminosity. For instance, for the case of SN1987A, this is $\sim 10^{46} J$.

Substituting our solution eq. (2.10) to the luminosity bound in eq. (4.2) we obtain the new astrophysical constraint on the neutrino transition magnetic moment,

$$\mu \lesssim \frac{10^{-8} \mu_B}{(B/1G)\sqrt{(L_0/1km) \times (R_{core}/10km)}}, \qquad (4.3)$$

Thus one sees that, even for a reasonable r.m.s. field $B \gtrsim 10^8$ G, this gives a more stringent bound than obtained from SN1987A [16]. It is also a more stringent bound than obtained above from nucleosynthesis in the presence of a primordial r.m.f. generated at the electroweak phase transition.

Note that the validity of eq. (4.3) requires averaging over small random r.m.f. distributions, so that the domain size L_0 must be much less than the core radius R_{core} and this, in turn, much less than the damping length $\Gamma_\perp^{-1}$. Indeed one can rewrite eq. (4.3) in such a way that both of the conditions

$$L_0 \ll R_{core} \ll \Gamma_\perp^{-1} = [(8/3)(\mu B)^2 L_0]^{-1} \qquad (4.4)$$

are verified, provided $L_0 \ll R_{core}$. The bound shown in Fig. 2 corresponds to $L_0 \sim 10$ meters in eq. (4.3).

5. Discussion and conclusions

We have derived bounds on transition neutrino magnetic moments connecting active to sterile Majorana neutrinos ν_a and ν_s for neutrinos which propagate in the presence of random magnetic fields. We treated both the constraints that follow from nucleosynthesis as well as those that follow from supernova energy loss arguments. In both cases our constraints are typically much stronger than in the case without random magnetic fields. The allowed region of parameters corresponding to each case are shown in Fig. 1 and Fig. 2. In the nucleosynthesis case our bounds are consistent with the hypothesis of seeding of galactic magnetic fields by primordial fields generated at the electroweak transition, and surviving past the re-combination epoch. In the supernova case, even for a reasonable r.m.s. field $B \gtrsim 10^8$ G, our constraint is more stringent than obtained from SN1987A. It is also a more stringent than the nucleosynthesis bound obtained in this paper in the presence of a primordial r.m.f. generated at the electroweak phase transition.

Our results were derived from a simple Schrödinger-type evolution equation ignoring neutrino interactions. Our results were derived for the general case of Majorana neutrino transition magnetic moments, but in the approximation of negligible mixing.

They are identical to the ones derived from nucleosynthesis in ref. [19] for the case of diagonal Dirac neutrino magnetic moments. This suggests that kinetic theory effects neglected here are not too important and should encourage one to perform a complete study in which the mixing between the two Majorana neutrino species is taken into account. We plan to come back to this question elsewhere.

Figure Captions

Fig.1.

Constraints on active to sterile neutrino transition magnetic moments derived from early universe nucleosynthesis in the presence of a strong random magnetic field. This limit is given as a function of the domain size and field geometry parameter p, explained in the text. The region corresponding to the seeding hypothesis of galactic magnetic fields is indicated by the hatching.

Fig.2.

Constraints on active to sterile neutrino transition magnetic moments derived from supernova cooling arguments. This limit is given as a function of the random magnetic field strength for typical values of neutrino energy, core density and size, and lepton abundances. Note however that, for the dynamo mechanism suggested in ref. [13] the seed field B_0 should be larger than $B_0 \gtrsim 2 \times 10^8 G$.

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