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Neutrino and astroparticle physics: Working group report

S MOHANTY and U A YAJNIK*
Physical Research Laboratory, Navrangpura, Ahmedabad 380 009, India
*Indian Institute of Technology, Powai, Mumbai 400 005, India

Abstract. The contributions made to the Working Group activities on neutrinos and astrophysics are summarized in this article. The topics discussed were inflationary models in Raman–Sundrum scenarios, ultra high energy cosmic rays and neutrino oscillations in 4 flavour and decaying neutrino models.

Keywords. Left-right symmetry; domain walls; leptogenesis; inflation; extra dimensions.

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1. Leptogenesis in the left-right symmetric model

J Cline and U A Yajnik

It was shown in [1] and [2] that topological defects such as domain walls and cosmic strings are generic to the left-right $(L–R)$ symmetric model with two triplet and one bidoublet Higgs fields. In particular, the discrete $L–R$ symmetry of the triplet Higgs potential implies that at the scale v_R, in a given region in the Universe, either the SU(2)$_L$ or the SU(2)$_R$ could break. Immediately after this phase transition therefore we would find domains with either of the above groups broken. Such domains will be separated by walls, dubbed $L–R$ walls in [2]. Phenomenology demands that the SU(2)$_R$ bosons must eventually acquire the larger masses, and that the domain structure disappears in order to not dominate the energy density in the Universe. This will be achieved if GUT scale physics induces small corrections to the triplet Higgs potential, making it energetically favourable for the SU(2)$_L$ bosons to remain the lighter species. This will be assumed in the following.

Recently with the increasing difficulties in explaining the baryon asymmetry either in the Standard Model or its minimal supersymmetric extension, the idea of leptogensis which then produces the baryon asymmetry has gained ground. In the $L–R$ model, the disappearance of the domains with broken SU$(2)_L$ provides a preferred direction for the motion of the domain walls. This can fulfill the out-of-thermal-equilibrium requirement needed for leptogenesis.

Consider the interaction of neutrinos from the $L–R$ wall which is moving into the energetically disfavored phase. The left-handed neutrinos, ν_L, are massive in this domain, whereas they are massless in the phase behind the wall. This can be seen from the Majorana mass term $h_M \Delta_L \overline{\nu^c_L} \nu_L$, and the fact that $\langle \Delta_L \rangle$ has a kink-like profile, being zero behind the wall and $O(v_R)$ in front of it.

To get leptogenesis, one needs an asymmetry in the reflection and transmission coefficients from the wall between ν_L and its CP conjugate $(\nu^c_L)^*$. This can happen if a CP-violating condensate exists in the wall as discussed below. Then there will be a preference for transmission of, say, ν_L. The corresponding excess of antineutrinos (ν^c_L) reflected in front of the wall will quickly eqilibrate with ν_L due to helicity-flipping scatterings, whose amplitude is proportional to the large Majorana mass. However the transmitted excess of ν_L survives because it is not coupled to its CP conjugate in the region behind the wall, where $\langle \Delta_L \rangle = 0$.

A quantitative analysis of this effect can be made either in the framework of quantum mechanical reflection, valid for domain walls which are narrow compared to the particles' thermal de Broglie wavelengths, or using the classical force method [3], which is appropriate for walls with larger widths. We adopt the latter here. The classical CP-violating force of the wall on the neutrinos, whose sign is different for ν_L and ν^c_L, is

$$F = \pm \frac{1}{2E^2} \left(m^2_\nu(x) \alpha'(x) \right)', \tag{1}$$

where $m^2_\nu(x)$ is the position-dependent mass and α is the spatially varying CP-violating phase. One can then write a diffusion equation for the chemical potential μ of the ν_L as seen in the wall rest frame

$$-D_\nu \mu'' - v_w \mu' + \theta(x) \, \Gamma_{\rm hf} \, \mu = S(x). \tag{2}$$

Here D_ν is the neutrino diffusion coefficient, v_w is the wall velocity, $\Gamma_{\rm hf}$ is the rate of helicity flipping interactions taking place in front of the wall (hence the step function $\theta(x)$), and S is the source term, given by

$$S(x) = -\frac{v_w D_\nu}{\langle \vec{v}^{\,2} \rangle} \langle v_x F(x) \rangle', \tag{3}$$

where $\vec{v}$ is the neutrino velocity and the angular brackets indicate thermal averages.

The spatially varying complex neutrino mass can be found from the finite-temperature effective potential for the L–R model, thus specifying the source (1), (3). The diffusion equation (2) can be solved using standard Green's function methods [4].

1.1 *CP violation*

An attractive feature of the left-right symmetric model is dynamical generation of the CP violation [5]. After accounting for the phases that can be eliminated by global symmetries and field redefinitions, the two remaining phases can be introduced into the VEVs [5,6]

$$\langle \Delta_R \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ v_R \mathrm{e}^{i\theta} & 0 \end{pmatrix} \quad \text{and} \quad \langle \phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} k_1 \mathrm{e}^{i\alpha} & 0 \\ 0 & k_2 \end{pmatrix}.$$

Here ϕ is the bidoublet and we use $\tilde{\phi}$ for its $\mathrm{SU}(2)$ conjugate matrix. We need the effective CP violating phase to enter as a position dependent function in the wall profile. It can be shown that the terms

$$\beta_1 \mathrm{Tr}(\phi \Delta_R \phi^\dagger \Delta_L^\dagger) + \beta_2 \mathrm{Tr}(\tilde{\phi} \Delta_R \phi \dagger \Delta_L^\dagger) + \beta_3 \mathrm{Tr}(\phi \Delta_R \tilde{\phi}^\dagger \Delta_L^\dagger) + \mathrm{h.c.}$$

give rise to a potential for the phases α and θ of the form

$$\beta_1 k_1 k_2 v_L v_R \cos(\alpha - \theta) + \beta_2 k_1^2 v_L v_R \cos(2\alpha - \theta) + \beta_3 k_2^2 v_L v_R \cos\theta.$$

The bidoublet is not expected to acquire a VEV at the temperature scale of v_R. However, in the presence of the position dependent VEV of the triplets, it can turn on in the wall interior. Analysis of the above potential shows that this gives rise to position dependent VEV's for the two phases as well.

Thus our preliminary analysis gives a strong indication of feasible leptogenesis during the epoch of disappearance of the L–R walls, which are generic to the left-right symmetric model. Depending on the parameters of the model, the lepton asymmetry could be quite large. If it is distributed in a certain way between the different lepton flavors, this large lepton asymmetry can be converted by sphalerons into a small baryon asymmetry [7]. Otherwise one expects a baryon asymmetry of the same order of magnitude as the lepton asymmetry.

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2. Inflation with bulk fields in the Randall–Sundrum warped compactification?

J Cline and U A Yajnik

2.1 *Introduction*

The Randall–Sundrum proposal for solving the hierarchy problem has received much attention in the last year [1]. They considered an extra compact dimension with coordinate y and line element

$$\mathrm{d}s^2 = a^2(y)\eta_{\mu\nu}\mathrm{d}x^\mu \mathrm{d}x^\nu + b^2 \mathrm{d}y^2, \tag{4}$$

where $a(y) = \mathrm{e}^{-kb|y|}$, $y \in [-1, 1]$ and the points y and $-y$ are identified so the extra dimension is an orbifold with fixed points at $y = 0$ and $y = 1$. The four-dimensional geometry is conformal to Minkowski space. The scale k in the warp factor $a(y)$ is determined by the 5-D cosmological constant, Λ and the analog to the Planck mass, M, by $k = (-\Lambda/6M^3)^{1/2}$. Hence Λ must be negative, and the 5-D space is anti-deSitter.

At $y = 0$ there is a positive tension brane (the Planck brane) on which particle masses are naturally of order the Planck mass, M_p, while at $y = 1$ there is a negative tension brane (called the TeV brane), where particle masses are suppressed by the warp factor e^{-kb}. By adjusting the size of the extra dimension, b, so that $kb \sim 37$, the masses of particles on the TeV brane will be in the TeV range, even if all the underlying mass parameters (including Λ, M and k) are of order M_p to the appropriate power.

Although this is desirable for solving the hierarchy problem, it makes it difficult to understand the origin of inflation. The density perturbations from inflation, $\delta\rho/\rho$, are suppressed by inverse powers of M_p. To get $\delta\rho/\rho \sim 10^{-5}$, one needs a mass scale much larger than 1 TeV in the numerator. By construction, the TeV scale is the cutoff on the TeV brane, so it is hard to see where such a scale could come from unless some physics outside of the TeV brane is invoked.

2.2 *Inflation with a bulk scalar*

In this study we investigate what happens when the inflaton is a bulk scalar field. The simplest possibility is chaotic inflation with a free field [2]. We will assume that the line element (4) is modified by replacing the Minkowski metric with 4-D deSitter space, whose line element is $\mathrm{d}s^2 = -\mathrm{d}t^2 + \mathrm{e}^{2Ht}\mathrm{d}\vec{x}^2$. The action for the bulk scalar is then

$$S = \frac{1}{2}\int \mathrm{d}^4x\,\mathrm{d}y\,b\,\mathrm{e}^{3Ht}a^4(y)\left[a^{-2}(y)\dot{\phi}^2 - b^{-2}\phi'^2 - m^2\phi^2\right]. \tag{5}$$

The equation of motion for ϕ is

$$a^{-2}\left(\ddot{\phi} + 3H\dot{\phi}\right) - b^{-2}\left(\phi'' - 4kb\phi'\right) + m^2\phi = 0 \tag{6}$$

and, assuming that the size of the extra dimension is stabilized [3], the Hubble rate is given by

$$H^2 = \frac{4\pi G}{3}\int_0^1 \mathrm{d}y\,b\,a^4(y)\left(a^{-2}\dot{\phi}^2 + b^{-2}\phi'^2 + m^2\phi^2\right), \tag{7}$$

where G is the ordinary 4-D Newton's constant.

We look for a separable solution, $\phi = \phi_0(t)f(y)$. We take ϕ_0 to have dimensions of (mass)^1 so that it is the canonically normalized field in an effective 4-D description, and f has dimensions of (mass)$^{1/2}$. We will also assume the slow roll condition is fulfilled so that the terms $\ddot{\phi}$ and $\dot{\phi}^2$ can be ignored in the last two equations. The equation of motion becomes

$$\dot{\phi}_0 = \frac{\mathrm{e}^{-2kby}}{3\hat{H}}\left(-m^2 + \frac{1}{b^2 f}\left(f'' - 4kbf'\right)\right) = \text{constant} \equiv -\Omega \tag{8}$$

with

$$\hat{H}^2 \equiv \frac{H^2}{\phi_0^2} = \frac{4\pi G}{3} \int_0^1 \mathrm{d}y\, b\, \mathrm{e}^{-4kby} \left(b^{-2} f'^2 + m^2 f^2 \right). \tag{9}$$

The solution for ϕ_0 is obviously linear, $\phi_0(t) = C - \Omega t$. For chaotic inflation we want $C \gg M_p$ (necessary to fulfill the slow roll condition [2]) and $\Omega > 0$, so that ϕ_0 is rolling to the minimum of its potential.

The equation for f becomes

$$f'' - 4kbf' - b^2 \left(m^2 - 3\hat{H}\Omega \mathrm{e}^{2kby} \right) f = 0. \tag{10}$$

This is the same equation as (7) of ref. [4]. The solutions f_n (called y_n in [4]) are discrete, such that $\hat{H}\Omega$ is quantized:

$$3k^{-2}\hat{H}\Omega \mathrm{e}^{2kb} = x_{n\nu}^2. \tag{11}$$

Here $x_{n\nu}$ is the nth root of the equation $2J_\nu(x_{n\nu}) + x_{n\nu}J'_\nu(x_{n\nu}) = 0$, where the order of the Bessel function is $\nu = \sqrt{4 + m^2/k^2}$. For m/k in the range 0.5–3, the lowest mode $x_{1,\nu}$ ranges from 4 to 6. This assumes the boundary condition that $f' = 0$ at $y = 0, 1$, but other choices of boundary conditions will lead to essentially identical conclusions, as we will explain below. The f_n's are normalized so that

$$\int_0^1 \mathrm{d}y\, b\, \mathrm{e}^{-2kby} f_n(y) f_m(y) = \delta_{mn}. \tag{12}$$

With the solution for f_n we can evaluate the rescaled Hubble rate, $\hat{H}$, in eq. (9). After a partial integration and use of the equation of motion (10), the integral in (9) becomes identical to that of (12), times $3\hat{H}\Omega$. This gives $\hat{H}^2 = 4\pi G\hat{H}\Omega$, which together with eq. (11) determines Ω, the rate at which ϕ_0 is rolling to its minimum:

$$\Omega = \frac{k x_{n\nu} \mathrm{e}^{-kb}}{\sqrt{12\pi G}} \sim M_p \times 1\ \mathrm{TeV}. \tag{13}$$

We used the fact that e^{-kb} is supposed to be of order (TeV)$/M_p$.

2.3 *Density perturbations*

We can now estimate the magnitude of density perturbations in this model. Using $\delta\rho/\rho \sim H^2/|\dot{\phi}_0|$,

$$\frac{\delta\rho}{\rho} \sim (4\pi G)^2\, \Omega\, (C - \Omega t)^2 \sim \frac{\mathrm{TeV}}{M_p}. \tag{14}$$

Although C is presumed to be super-Planckian, $(C - \Omega t)$ will not be orders of magnitude larger than M_p near the end of inflation, when the perturbations with COBE-scale wavelengths were being produced; hence we take $(C - \Omega t) \sim M_p$ in the above estimate. This suppression of the density perturbations makes our model not viable.

Nevertheless, it is interesting to compare to what would happen if we tried to do chaotic inflation using a scalar field trapped on the TeV brane. The mass of the field is now constrained to be of order $m \sim$ TeV because of the suppression of masses by the warp factor. The equation of motion during the slow roll regime is

$$\dot{\phi} = -\frac{m^2\phi}{3H} = -\frac{m^2\phi}{\sqrt{4\pi G m^2\phi^2/3}}; \tag{15}$$

hence ϕ evolves linearly with time, and $\dot{\phi}$ is of order $M_p \times 1$ TeV, just as with the bulk scalar field. And the estimate for $\delta\rho/\rho$ has the same parametric form. All this, despite the fact that we started with a bulk scalar whose mass is Planck-scale in the 5-D Lagrangian.

In retrospect, this result is not surprising. Reference [4] noted that the modes of the bulk scalar behave similarly to TeV-scale particles on the brane. This can be understood by the form of the solutions [4],

$$f_n \sim \mathrm{e}^{2ky} J_\nu \left(x_{n\nu} \mathrm{e}^{kb(y-1)} \right), \tag{16}$$

which are strongly peaked near $y = 1$. There is thus little practical difference between the low-lying modes of the bulk field and a field confined to the TeV brane.

One might wonder if this conclusion depends on the choice of boundary conditions for the modes f_n. However the fact that the modes peak at the TeV brane comes from bulk energetics, not boundary conditions. Since the mass of the bulk field is effectively varying like e^{-kby}, it is energetically much more efficient for the field to be concentrated near $y = 1$.

2.4 *Conclusion*

The simplest chaotic inflation models seem to be ruled out in the Randall–Sundrum scenario, whether the inflaton is a bulk field or one restricted to the TeV brane. One could, alternatively, put the inflaton on the Planck brane (at $y = 0$), at the cost of reintroducing a hierarchy problem—why should m/M_P be $O(\delta\rho/\rho) \sim 10^{-5}$? This fine-tuning problem always occurs in inflation, but the RS setting casts it in a somewhat new light. One could invent an intermediate brane for the inflaton, which has just the right mass scale, but this seems artificial. Perhaps the RS idea, if correct, is telling us that hybrid inflation (involving more than one field) is necessary.

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3. A parametrization for vacuum mixing matrix in four generations

Sandhya Choubey, Gautam Dutta, Srubabati Goswami, D Indumathi, Debashis Majumdar, M V N Murthy

Many participants mentioned above, either together or separately, are studying the solar, atmospheric and supernova neutrinos and their signatures at the earthborne detectors. It has become clear, since the announcement of the LSND results, that one may need a fourth sterile neutrino in order to account for all the experimental data. During the discussions the participants attempted to arrive at a parametrization for the vacuum mixing matrix which would be used by all the members in analysing the data from different detectors as well as different phenomena involving neutrinos. This would then facilitate comparison of the results emanating from different groups. Basically the following constraints arising from experiments were imposed.

1. The solar and atmospheric neutrino data indicate the existence of two very different scales, 10^{-6} eV^2 and 10^{-3} eV^2, for the mass squared differences between neutrinos. However LSND data requires in addition a scale in the range of eV^2, different from both atmospheric and solar neutrino results. The favoured scenario is that in which there are two neutrino doublets. The mass squared difference in each doublet is given by the solar and atmospheric neutrino requirements. The doublets themselves are separated by a scale corresponding to the LSND data.
2. While this defines the structure of the mixing matrix, further constraint is imposed using the results from the CHOOZ reactor neutrino experiment. This result basically imposes an upper limit for the conversion of electron anti-neutrino to any other flavour. The limit on the relevant mixing parameter is $\epsilon \leq 0.07$,

We make use of the fact that the parameter ϵ is small compared to unity, and parametrize the unitary mixing matrix. The mixing matrix, which relates the flavour and mass eigenstates in the four generation scenario has six angles and the CP-violating phases which we do not consider here. As is the convention, we denote the mixing angle in the lower doublet by ω and the mixing angle in the upper doublet by ψ. The resulting complicated mixing matrix involving six mixing angles is greatly simplified by application of the CHOOZ constraint. Then the flavour eigenstates are related to the four mass eigenstates in vacuum through a unitary transformation,

$$\begin{bmatrix} \nu_e \\ \nu_s \\ \nu_\mu \\ \nu_\tau \end{bmatrix} = U^v \begin{bmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \\ \nu_4 \end{bmatrix}, \tag{17}$$

where the superscript v on the r.h.s. stands for vacuum. Within the two doublet scheme the 4×4 unitary matrix, U^v, may be written as

$$U^v = \begin{pmatrix} (1-\epsilon^2)c_\omega & (1-\epsilon^2)s_\omega & \epsilon & \epsilon \\ -(1-\epsilon^2)s_\omega - 2\epsilon^2 c_\omega & (1-\epsilon^2)c_\omega - 2\epsilon^2 s_\omega & \epsilon & \epsilon \\ \epsilon(s_\omega - c_\omega)(c_\psi + s_\psi) & -\epsilon(s_\omega + c_\omega)(c_\psi + s_\psi) & (1-\epsilon^2)c_\psi - 2\epsilon^2 s_\psi & (1-\epsilon^2)s_\psi \\ \epsilon(s_\omega - c_\omega)(c_\psi - s_\psi) & -\epsilon(s_\omega + c_\omega)(c_\psi - s_\psi) & -(1-\epsilon^2)s_\psi - 2\epsilon^2 c_\psi & (1-\epsilon^2)c_\psi \end{pmatrix}, \tag{18}$$

where $s_\omega = \sin\omega$ and $c_\omega = \cos\omega$, etc. The angles ω and ψ can take values between 0 and $\pi/2$. The mixing matrix given above is unitary up to order ϵ^2. Since ϵ is a small parameter, we do not need to go beyond this order.

4. Three flavour MSW fits to super-K data

S R Dugad, Mohan Narayan and Uma Shankar

There are several puzzles in the super-Kamiokande solar neutrino data:

1. Higher energy bins (12–14 MeV) see less suppression than others. The energy dependence of the observed neutrino flux does not follow any of the three oscillation models – vacuum oscillations, small angle MSW or large angle MSW.
2. There is less suppression seen in the neutrino flux in the early part of night than at midnight. This is at odds with the expectation that regeneration by the earth core is most effective at midnight.

This group intends to do MSW fits again using only the gallium data and the super-K data on spectrum and time of night variations.

5. ν decay solution to solar neutrino problem

S Choubey, S Goswami and D Mazumdar

This group examines models where there is a mixing between $\nu_e \to \nu_\mu$ and in addition there is a decay of the heavier neutrino $\nu_{2L} \to \nu_{1R} + \mathrm{majoron}$. The survival probability of ν_e is then given by (for $\Delta m^2 > 10^{-4}\,\mathrm{eV}^2$)

$$P(\nu_e \to \nu_e) = (1 - |U_{e2}|^2)^2 + |U_{e2}|^2 \exp\left(-\frac{Lm_2}{\tau_0 E}\right). \tag{19}$$

A fit with the super-K spectrum data gives $\chi^2_{\min} = 14.87$ for 15 d.o.f with the best fit values of the parameters $|U_{e1}|^2 = 0.63, m_2/\tau_0 = 5.62 \times 10^{-11}\mathrm{eV}^2, \chi_n = 12.6/5.15$ and is allowed at 46.08% C.L. with respect to BP-98 standard solar model.

In comparison vacuum oscillations model gives $\chi^2_{\min} = 13.07$ for 15 d.o.f and MSW gives $\chi^2_{\min} = 17.23$ for 15 d.o.f.

Detailed calculations of rates and spectral data analysis is planned.

6. Role of collective neutrino forces in large scale structures

S Mohanty, V Sahni and A M Srivastava

Numerical simlations of cold + hot dark matter models show that the density perturbations do not match observations at small scales (100 kpc) and large scales (100 Mpc). The CDM component of dark matter clusters more than what is observed at scales smaller than

100 kpc and the DM component free-streams and gives excess power at scales larger than 100 Mpc.

The proposal of this group is to study the collective neutrino force on the CDM component which can be first order in G_F

$$F_c = G_F(-\nabla n_\nu + \vec{v}_c \times \nabla \times \vec{J}_\nu). \tag{20}$$

This is different from the usual collisional force which is $f = G_F^2 n_x$ which decouples at very early times ($t \sim 1\,\mathrm{sec}$).

The Boltzmann equation of the CDM particles

$$\vec{v}_c \cdot \nabla_x f_c + \vec{F}_c \cdot \nabla_p f_c = 0. \tag{21}$$

where the force term includes the collective weak interactions in addition to gravity

$$F_c = -m_c \nabla_x \phi + G_F(-\nabla n_\nu + \vec{v}_c \times \nabla \times \vec{J}_\nu). \tag{22}$$

A detailed analysis is proposed to check if the introduction of this collective weak force can prevent excessive clustering of CDM.

7. Signature of Dirac vs Majorana nature of neutrinos from HECR observations

K R S Balaji, P Bhattacharjee, S Mohanty and S N Nayak

We examine the compatibilty of the Wyler mechanism which explains the observed UHECR above the GZK cutoff energies (10^{19} GeV) via annihilation of high energy neutrinos with cosmic background neutrinos [1] – with the neutrino mass matrix suggested by the solar and atmospheric neutrino problems. If the neutrinos are Dirac then there is a resonant spin flip by the magnetic fields of AGN's from active to sterile chirality and the flux of UHECR would be decreased by about 75%.

Resonant spin flip in AGN's

Neutrinos are produced in AGNs or radio galaxies from pions arising from collisions of protons after undergoing shock acceleration. Each π^+ ultimately produces one each of ν_μ, ν_μ^c and ν_e. The flavour oscillations between different active species will not change the number of real Z's produced by annihilation with the cosmic background neutrinos. Oscillation from an active to sterile neutrinos which is possible in the case of Dirac neutrinos will reduce the number of Z's produced, and will have an observable consequence in the decrease of UHECR produced $CB\nu$ annihilation.

The hamiltonian governing the propagation of two neutrino species $(\nu_{Le}, N_{R,\alpha})$ where $\alpha = \mu$ or τ is given by

$$i\frac{\mathrm{d}}{\mathrm{d}t}\begin{pmatrix} \nu_{Le} \\ N_{R,\alpha} \end{pmatrix} = \begin{pmatrix} V_c + V_n & \mu_{e\alpha} B \\ \mu_{e\alpha} B & \frac{\Delta m^2}{2E}\cos 2\theta \end{pmatrix} \begin{pmatrix} \nu_{Le} \\ N_{R,\alpha} \end{pmatrix}, \tag{23}$$

where $\Delta m^2 = (m_2^2 - m_1^2)$ and θ is the vacuum mixing angle between ν_e and ν_α. The matter potentials due to charged and neutral current are $V_c = 2\sqrt{2} G_F N_e$ and $V_n = -\sqrt{2} G_F N_n$ respectively (N_e and N_n are the number densities of electrons and neutrons respectively).

The mixing angle in matter and magnetic field is given by

$$\tan 2\theta_M = \frac{2\mu_{e\alpha} B}{\frac{\Delta m^2 \cos 2\theta}{2E} - (V_c + V_n)}. \tag{24}$$

The transition probability between $\nu_{Le} \to N_{R\alpha}$ is given by

$$P(\nu_{Le} \to N_{R\alpha}) = \frac{1}{2} - \left(\frac{1}{2} - \mathrm{e}^{-\pi\kappa/2}\right) \cos 2\theta \cos 2\theta_M(r_i), \tag{25}$$

where the matter mixing angle is evaluated at the point of production of ν_{Le} and where the adiabaticity parameter κ is given by

$$\kappa(r_{\rm res}) = \frac{E\mu^2 B^2}{\Delta m^2 \cos 2\theta |\mathrm{d} \ln(V_c + V_n)/\mathrm{d}r|}. \tag{26}$$

We now calculate the amount of suppression of active neutrino flux due to resonant spin-flavour precession (RSFP) in the AGN core model as a source of the high energy neutrinos.

Active galactic nuclei are the most luminous objects in the universe with luminosities in the range 10^{42}–10^{48} ergs/sec. These are supposed to be powered by matter accreting into black holes of masses 10^4–$10^{10} M_\odot$. According to Szabo and Protheroe [6], the matter density outside the core has the profile

$$\begin{aligned} \rho(r) &= \rho_0 \, (r/R_s)^{-2.5} (1 - 0.1 (r/R_s)^{0.31})^{-1} \quad \left(\frac{10^{48} \mathrm{ergs/sec}}{L_{\rm AGN}}\right) \\ \rho_0 &= 1.4 \times 10^{-15} \mathrm{g/cm}^3, \end{aligned} \tag{27}$$

$R_s = 3 \times 10^{11} (M_{\rm AGN}/10^8 M_\odot)$ being the Schwarzchild radius of the AGN. The magnetic field profile is given by

$$\begin{aligned} B(r) &= B_0 \, (r/R_s)^{-1.75} (1 - 0.1 (r/R_s)^{0.31})^{-0.5} \quad \left(\frac{10^{48} \mathrm{ergs/sec}}{L_{\rm AGN}}\right)^{1/2} \\ B_0 &= 5.5 \times 10^3 \mathrm{G}. \end{aligned} \tag{28}$$

At the neutrino energies of interest, $E = 10^{20}$ eV, the resonance condition $\frac{\Delta m^2 \cos 2\theta}{2E} - (V_c + V_n) = 0$ is attained in the AGN (at distances $r > 10 R_s$) for $\Delta m^2 \cos 2\theta = 10^{-11}$ eV$^2 (\rho/10^{-18} \text{ gm cm}^{-3})(E/10^{20} \text{ eV})$. This mass difference is in the same range as required for vacuum oscillations (between ν_e and ν_μ) in order to solve the solar neutrino neutrino problem.

The magnetic field at the core ($r = 10 R_s$) is $B = 10$ G, we have $\mu B = 1.7 \times 10^{-26}$ eV. The mixing angle (24) at the core is $\cos\theta_M = \cos(\pi - \arctan(10^{-23}/10^{-24}) = -0.8$. The probability for conversion of ν_{Le} to a sterile $N_{R,\mu}$ is

$$P(\nu_L \to N_{R\nu}) = 0.85; \quad \Delta m^2 = 5 \times 10^{-11} \mathrm{eV}^2 \text{ and } \sin^2 2\theta = 6 \times 10^{-3}. \tag{29}$$

Hence the flux of active ν_e from the AGN's at energies 10^{22} eV is reduced by (67–85)% if the neutrinos are Dirac. If on the other hand they are Majorana then the transitions will be between two active species and there will be no reduction in the number of Z's produced when they anhilate the relic neutrinos.

For the case of $(\nu^c_{R,\mu})$ neutrinos produced at the core, the Hamiltonian governing its oscillation to a sterile N_{Le} will be the one in (23) with extra minus signs on both the diagonal terms and with $(V_c = 0)$. The resonance condition is therefore the same and $\nu^c_{R,\mu}$ will convert to a sterile neutrino with a probabilty (0.6–0.8).

For the $\nu_{L,\mu}$ neutrinos produced at the AGN core there is no matter resonanace. These will precess in the magnetic field with a probability

$$P(\nu_{L\mu} \to N_R) = \sin^2(\mu B L) \to \frac{1}{2}. \tag{30}$$

So the total flux of all three species of neutrinos produced at the core reduces by about 75% if the neutrinos have a Dirac mass of about 1 eV.

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