Selected topics in neutrino astrophysics are reviewed. These include the production of low energy neutrino flux from cores of collapsing stars and the expected high energy neutrino flux from some other astrophysical sites such as the galactic plane as well as the center of some distant galaxies. The expected changes in these neutrino fluxes because of neutrino oscillations during their propagation to us are described. Observational signatures for these neutrino fluxes with and without neutrino oscillations are discussed.

I. GENERAL INTRODUCTION

In the standard model of the electro weak interactions, the lepton masses and the values of other parameters such as weak mixing angle, couplings, etc. are arbitrary and are therefore determined by experiments. These parameters are independent of each other and can not be determined uniquely, while the neutrino is taken to be massless because of maximal parity violation. The masslessness of the neutrino however does not follow from any other theoretical ground unlike the local gauge invariance for the photon [1, 2, 3].

In order to search for physics beyond the standard model of particle physics, light neutrino masses are incorporated in the extensions of the standard model. Then, on quite general grounds one expects that neutrinos will also possess non-zero magnetic moments. Experimentally, one places finite non-vanishing upper bounds on measurable neutrino masses and magnetic moments. Presently, there is some indirect experimental evidence for the masslessness of the neutrino [4].

Massive neutrinos quite likely mix. The mixing of quarks is an established fact and because of quark-lepton symmetry, it is natural to assume that leptons exhibit mixing as well. An additional argument, in this sense, is provided by grand unification models, in which quarks and leptons are described in a unified manner.

Mixing means that $\nu_e$, $\nu_\mu$, and $\nu_\tau$, i.e., the states created in weak interactions, are different from the states $\nu_1$, $\nu_2$ and $\nu_3$ that have definite masses. The neutrinos $\nu_e$, $\nu_\mu$ and $\nu_\tau$ are orthogonal combinations of $\nu_1$, $\nu_2$ and $\nu_3$ with different phases between them. In addition, the sterile neutrinos may mix with these neutrinos.

Neutrino mixing has, as its consequence, neutrino oscillations, i.e., the process of periodic (complete or partial) conversion of neutrinos of one type into another, for instance, $\nu_e \rightarrow \nu_\mu \rightarrow \nu_e \rightarrow \ldots$. The components $\nu_i$ of a mixed neutrino have different masses and hence different phase velocities. It follows that the phase difference caused by the mass difference between the $\nu_i$ vary monotonically during the propagation. This phase change manifests itself as neutrino oscillations.

* Based on lectures given at The Sixth Constantine High Energy Physics School on Strong and Weak Interactions Phenomenology, 6 – 12 April, 2002, Constantine, Algeria.
If neutrino oscillations do occur in vacuum, matter can enhance their depth (probability amplitude) up to a maximal value. That is, a monotonic change of density may lead to resonant conversions between various neutrino flavors. This follows from the fact that when neutrinos propagate through a monotonically changing density medium, $\nu_e$ and $\nu_\mu$ ($\nu_\tau$) feel different potentials, because $\nu_e$ scatters off electrons via both neutral and charged currents, whereas $\nu_\mu$ ($\nu_\tau$) scatters off electrons only via the neutral current. This induces a coherent effect in which maximal conversion of $\nu_e$ into $\nu_\mu$ take place (even for a rather small intrinsic mixing angle in the vacuum), when the phase difference arising from the potential difference between the two neutrinos cancel the phase caused by the mass difference in the vacuum.

During nearly past half a century, the empirical search for neutrinos has spanned roughly four orders of magnitude in neutrino energy $E$, from $\sim$ MeV up to $\sim 10^6$ MeV. The lower energy edge corresponds to the Solar neutrinos, whereas the upper energy edge corresponds to the Atmospheric neutrinos. A detailed early description of the Solar neutrino search can be found in [7], whereas for recent status, see [8]. The aspects of neutrino production in Atmosphere of earth related to neutrino oscillation studies are recently reviewed in [9]. The intermediate energy range corresponds to terrestrial neutrinos such as from reactors (and accelerators) and the supernova neutrinos. Thus, obviously either going in energy range below these values or above are the available frontiers. For a general introduction of the possibility of having neutrinos with energy $> 10^6$ MeV, see [10], the upper energy edge for these high energy neutrinos is limited only by the concerned experiments. More detailed general discussions in the context of high energy neutrinos can be found in [11, 12]. Despite the availability of the somewhat detailed discussion of progress cited in the last reference, the field of high energy neutrino astrophysics is still passing through its initial stage of development.

The above mentioned empirical search has already given us quite useful insight into neutrino intrinsic properties such as mass and mixing. Massive neutrinos and their associated properties such as Dirac or Majorana character of their mass, their mixings and magnetic moments can have important consequences in astrophysics. In these two lectures, I elaborate some of the selected consequences and the constraints implied by these consequences on neutrino properties as well as the insight that one may gain about the nature of the astrophysical or and cosmological sites and the interactions that produce these neutrinos. The explanation of observed Solar $\nu_e$ deficit relative to its production value in the core of the Sun, via $\nu_e \rightarrow \nu_\mu, \nu_\tau$ conversion is an interesting example in this context [13].

The general plan of the lectures is as follows. In the first lecture, I elaborate these aspects for the low energy neutrino flux ($E \sim$ MeV) emitted during the gravitational collapse of stars which may be accompanied by supernovae phenomenon. In the second lecture, I elaborate these for the expected high energy neutrino flux ($E \geq 10^6$ MeV) from some representative examples of remaining cosmos around us such as our galactic plane. This includes the one arising from the interaction of ultra high energy cosmic ray flux ($E \geq 10^{12}$ MeV) with the matter and radiation inside the sources (such as center of our and other galaxies) of as well as during propagation of ultra high energy cosmic ray flux to us.
II. LECTURE 1: NEUTRINOS FROM COLLAPSING STARS

A. Introduction

Almost \(\sim 99\%\) of the binding energy of a few Solar mass collapsing core of a star is released in the form of neutrinos typically in an order of magnitude energy interval between 5 and 50 MeV. In contrast, in the Sun, the energy released from the core in the form of neutrinos is of the order of \(\sim 1\%\). Therefore, neutrinos from collapsing cores are expected to bring relatively more useful information about the mechanism of the collapse as well as information about the intrinsic properties of neutrinos itself, though still difficulties remain in obtaining a self consistent and a clear understanding of the exact mechanism of supernova phenomena. The rather optimistic frequency of occurrence of 1987 A like supernova is roughly once per thirty years. The last one occurred in 1987 implying that it is by now already half way. On the other hand, the estimates of cumulative neutrino flux from all the supernovae that occurred in the past are also rather close to the present relevant detection threshold \([14]\). For a more detailed introduction, see \([15, 16, 17, 18]\).

B. Neutrino production

Type II supernovae are the most interesting from the point of view of neutrino flux studies. Present theories of type II supernovae assume that they are catastrophic endpoints in the evolution of massive stars with \((8 - 10) \leq M/M_\circ \leq 70\), where \(M_\circ \sim 2 \cdot 10^{33}\) g is the Solar mass. Such stars form unstable iron cores (or if the mass is closer to the lower limit, O, Ne or Mg cores) supported mainly by degenerate electron gas pressure. Partial photo dissociation of nuclei (for instance, \(\gamma + ^{56}Fe \rightarrow 13\alpha + 4n\)) and electron capture by free protons \((e^- + p \rightarrow n + \nu_e)\) as well as by nuclei \([e^- + (Z, A) \rightarrow (Z - 1, A) + \nu_e]\), known as neutronization, causes the core to undergo a dynamic collapse when its mass exceeds the Chandrasekhar mass \(M_{\text{Ch}}\) (which is function of \(Y_e\)), namely, when \(M > M_{\text{Ch}} \equiv 5.83Y_e^2M_\circ\). Here \(Y_e = n_e/n\) is the number of electrons per baryon. It is estimated that most of the energy during this collapse is carried by the neutrinos. During the collapse, the neutrinos are emitted not only by neutronization but also by the thermal emission. The significant thermal emission interactions include: pair annihilation \((e^+ + e^- \rightarrow \nu + \bar{\nu})\), plasmon decay (plasma excitation \(\rightarrow \nu + \bar{\nu}\)), photo annihilation \((\gamma + e^- \rightarrow e^- + \nu + \bar{\nu})\) and neutrino bremsstrahlung \([e^- + (Z, A) \rightarrow (Z, A) + e^- + \nu + \bar{\nu}]\). In all these interactions, the electron neutrino pair can be produced by \(Z\) or by the \(W^\pm\) exchange, whereas the muon and tau neutrino pairs can be produced only by \(Z\) exchange. However, each neutrino emission interaction has an inverse interaction corresponding to absorption. Both absorption and inelastic interactions impede the free escape of neutrinos from a collapsing core. The relevant important interactions include: free interactions \((\nu + n \rightarrow \nu + n)\), interactions of heavy nuclei with \(A > 1 \ [\nu + (Z, A) \rightarrow \nu + (Z, A)]\), nucleon absorption \((\nu_e + n \rightarrow p + e^-)\) and electron neutrino interaction \((e^- + \nu \rightarrow e^- + \nu)\). Similar interactions occur for anti neutrinos. The cross section for these interactions defines the depth of various neutrino spheres. These interactions tend to
thermalize the various neutrinos and thus contribute to neutrino opacity.

Thus, the core of a collapsing star is the source of all neutrino and anti neutrino species. The fluxes change with time and are different for different species. During the initial stages of collapse, $\nu_e$ flux from the neutronization dominates. On the $\nu$ opacity stage, comparable (but not equal) fluxes of all flavors are emitted. Due to the difference in inelastic interactions of $\nu_e$ and $\nu_\mu$ ($\nu_\tau$), the $\nu_\mu$ ($\nu_\tau$) neutrino spheres lie deeper (and consequently at higher temperatures) than the $\nu_e$ neutrino sphere. This results in the higher energies of $\nu_\mu$ ($\nu_\tau$) with respect to the $\nu_e$ energies. Typically $E_{\nu_\mu} = (2 - 3) \cdot E_{\nu_e}$ but in the smaller flux of these neutrinos: $F^0_{\nu_\mu} = \frac{1}{2} \cdot F^0_{\nu_e}$. The energies as well as fluxes of non electron type neutrinos ($\nu_\mu, \nu_\tau, \bar{\nu}_\mu, \bar{\nu}_\tau$) are considered to be roughly equal. The spectra of various neutrinos being (approximately) Fermi-Dirac spectra with different temperatures and with high energy cut.

Summarizing, the neutrino radiation of a type II supernova consists of two components: a $\sim 10$ ms electron neutrino burst from the neutronization of dense matter which is followed by $\sim 10 - 20$ s thermal radiation of $\nu - \bar{\nu}$ pairs of all types; $\nu_e, \bar{\nu}_e, \nu_\mu, \bar{\nu}_\mu, \nu_\tau, \bar{\nu}_\tau$ typically with the above characteristics. The thermal $\nu$'s are emitted via black body radiation from the surfaces of the corresponding neutrino spheres. Their fluxes and spectra are therefore determined by the neutrino spheres, which in turn are essentially determined by the $\nu$ interaction cross section as mentioned earlier.

The relevant feature of the above characteristics for present elaboration is that the neutrinos being produced in the central part of collapsing the star and crossing the matter layers with decreasing density (from $\rho \sim 10^{14}$ g/cm$^3$ to approximately zero) may undergo resonant conversions. There are in general two aspects of this phenomena. The resonant oscillations may change the properties of $\nu$ burst which is important for the burst detection. Moreover, the change of the $\nu$ fluxes may influence the evolution of star: the dynamics of the collapse itself and of the expelling envelope. I intend to elaborate the former aspect in this lecture. The basic characteristic features of the neutrino fluxes, i.e., the flux spectra of the various neutrino species, etc., are the main testable ingredients of not only a stellar collapse theory but also of a possible occurrence of matter enhanced neutrino oscillations. These can be studied by existing and future detectors in some detail for neutrino bursts from nearby supernovae ($\leq 10$ kpc, where 1 pc $\sim 3 \cdot 10^{18}$ cm).

C. Oscillations during propagation: Effects of neutrino mixing

Because the matter density inside a collapsing star is quite high so that a necessary condition for resonant neutrino conversions, namely the level crossing or resonance is satisfied (see later). I, therefore, will elaborate only the effect of matter enhanced neutrino oscillations on the neutrino flux inside the matter and magnetic fields of collapsing stars. There are no matter enhanced spin flavor oscillation effects between a nearby supernova and the earth for downward going neutrinos.

A large range of neutrino mixing parameters determining the (two, three and four) flavor oscillations are presumably
get measured in terrestrial neutrino oscillation experiments with quite a good accuracy in near future perhaps before the next nearby supernova occurrence. This statement is further supported when one includes the already existing information on neutrino mixing parameters from Solar and Atmospheric neutrino flux measurements. The pure neutrino flavor oscillation effects may thus possibly be disentangled from for instance pure neutrino spin flavor oscillation effects for supernova neutrinos. In the possible presence of relatively strong magnetic fields in collapsing stars, the role of neutrino spin flavor oscillations become relevant. Therefore, I mainly elaborate here the effects of pure neutrino spin flavor oscillations of the type $\nu_e \leftrightarrow \bar{\nu}_\mu$, $\bar{\nu}_e \leftrightarrow \nu_\mu$ in collapsing stars, assuming the smallness of neutrino flavor mixing for illustration only. The neutrino spin flavor oscillations can take place for Majorana neutrinos for vanishing vacuum mixing also as the Majorana type neutrino magnetic moment can mix both the helicity and flavor of the two neutrino states. A recent detailed discussion on pure flavor oscillation effects for neutrinos from collapsing stars is given in [19]. At the end of this subsection, I will briefly compare the observational signatures of pure flavor and pure spin flavor oscillations (in absence of flavor mixing) for completeness. The general case of neutrino spin flavor oscillations in the presence of non vanishing vacuum mixing is briefly discussed in [20].

Neutrino spin flavor oscillations are commonly studied by numerically solving a Schrödinger like system of equations with an effective Hamiltonian. Using the notation of [21], I briefly outline the main steps to obtain it in a particular neutrino basis. I start from a most general lagrangian density, $\mathcal{L}$, describing the neutrino propagation in the presence of varying matter density and electromagnetic fields:

$$\mathcal{L} = i\bar{\psi}_L \frac{\partial}{\partial t} \psi_L + i\bar{\psi}_R \frac{\partial}{\partial t} \psi_R - \bar{\psi}_L h_L \psi_L - \bar{\psi}_R h_R \psi_R - \bar{\psi}_L h_{LR} \psi_R - \bar{\psi}_R h_{RL} \psi_L,$$

where $h_L, h_R, h_{LR}, h_{RL}$ are matrices with lepton flavor and Dirac indices. These matrices are in general space time dependent. In the Dirac case, $\psi_L$ and $\psi_R$ are independent fields and thus $\mathcal{L}$ conserves total lepton number. The equation of motions that follow from $\mathcal{L}$ can be written as

$$\frac{\partial \psi}{\partial t} = H_D \psi.$$  

Here, $\psi = \psi_L + \psi_R$ and $H_D$ is given by

$$H_D = -i\alpha \cdot \nabla + \beta(h_L P_L + h_R P_R + h_{LR} P_R + h_{RL} P_L).$$

The left and right-handedness is defined by

$$P_L \psi_L = \psi_L, \quad P_R \psi_R = \psi_R, \quad \text{with} \quad P_{L,R} = \frac{1}{2}(1 \mp \gamma_5),$$

and $\alpha_j \equiv \gamma^0 \gamma^j$ ($j = 1, 2, 3$) and $\beta \equiv \gamma^0$.

A equation similar to Eq. (2) can be obtained for Majorana neutrinos. The total Hamiltonian comprising both is

$$H_T = -i\alpha \cdot \nabla + V_L P_L + V_R P_R + V_S \beta P_L + V_S^\dagger \beta P_R + \frac{1}{2}\beta (\omega P_L + \omega^\dagger P_R)\sigma_{ab} F^{ab},$$

(5)
where $V$’s are in general space time dependent potentials in matrix form in flavor basis. These describe the effects of neutrino interactions with the background particles and can be obtained using finite temperature and density field theory approach [22]. The last term in above equation describes the effective neutrino electromagnetic interactions in usual notation. The flavor matrix $\omega$ contains both electric and magnetic dipole moments. The $V_L$ and $V_R$ are hermitian. The $H_T$ evolves a system of field equations that contain both negative and positive energy solutions. From implications point of view, it proved convenient to eliminate the negative energy solutions. This can be obtained by applying the following (unitary) Foldy-Wouthuysen transformation to $H_T$ and $\psi$:

$$U_{FW} = \frac{1}{\sqrt{2}} \left( \begin{array}{c} 1 \\ \sigma_3 \\ 1 \end{array} \right) \cdot \left( \begin{array}{c} \sigma_3 \\ 1 \\ 1 \end{array} \right).$$

(6)

The resulting one dimensional ($x \equiv x^3$) Schrödinger like equations of motion including the effective Hamiltonian can be written as

$$i \left( \begin{array}{c} \dot{\nu}_{eL} \\ \dot{\bar{\nu}}_{eR} \end{array} \right) = \left( \begin{array}{cc} 0 & \mu B(r) \\ \mu B(r) & 0 \end{array} \right) \left( \begin{array}{c} \nu_{eL} \\ \bar{\nu}_{eR} \end{array} \right),$$

(7)

where only magnetic dipole moment, $\mu$ connecting the same neutrino flavor in vacuum (no neutrino interaction effects, namely $V_L = V_R = V_S = 0$) is considered. In Eq. (7), the $\cdot$ denotes differentiation w.r.t distance as I use $\hbar = c = 1$ for relativistic neutrinos. Note that the strength of the magnetic field is assumed to be varying along the neutrino trajectory.

Neutrino spin precession probability is obtained by solving above system of equations, it is defined as $P(\nu_{eL} \rightarrow \bar{\nu}_{eR}; r) = |\langle \bar{\nu}_{eR}(r) | \nu_{eL}(0) \rangle|^2$ or

$$P(\nu_{eL} \rightarrow \bar{\nu}_{eR}; r) = 1 \cdot \sin^2 \left( \mu \int_0^r B(r')dr' \right),$$

(8)

where, I have assumed that $\nu_{eL}(r = 0) = 1$ and $\bar{\nu}_{eR}(r = 0) = 0$. Note that $P$ has maximal depth in vacuum (the pre factor 1) and is independent of $E$. For $P \neq 0$, one needs $\mu = B(r) \neq 0$, simultaneously. The neutrino spin precession length, for a constant $B$, is defined as

$$l_B = \pi/2\mu B,$$

(9)

namely when the argument of $\sin^2$ in Eq. (8) is $\pi/2$, so that if $r = l_B$, then $P = 1$. As an elementary example, let me ask a question in the context of Solar neutrinos: what $B_\odot$ is required to get $P = 1/2$ for $\mu \sim 10^{-11}\mu_B$ (assuming a constant $B_\odot$ in $0.7 \leq r/R_\odot \leq 1$, for simplicity)? Here $\mu_B \equiv e/2m_e$ is Bohr magneton and $R_\odot \sim 7 \cdot 10^{10}$ cm is Solar radius. The result is $B_\odot \sim 10$ kGauss. This is referred to as the Voloshin, Vysotsky and Okun (VVO) solution for the long standing Solar neutrino problem [23]. For a review of neutrino spin flavor oscillation solution to Solar neutrino problem, see [24], whereas for a recent discussion, see [25].

Let me recall here that in the standard model of particle physics with $m_\nu \neq 0$, $\mu \propto m_\nu$ and is therefore $O(10^{-19})\mu_B$ for $m_\nu \sim O(1)\text{eV}$. However, the $\mu$ can be as high as $10^{-12}\mu_B$ in some extensions of standard model [2]. It is primarily
because $\mu$ dependence on $m_\nu$ can be changed for instance, to $m_\alpha$ (where $\alpha = e, \mu, \tau$). In some extensions of standard model, it can be achieved by introducing new symmetries in the relevant standard model lagrangian density (and then breaking these). The same can also be achieved either by enlarging the matter particle sector or higgs or/and gauge boson sector of standard model. The present upper bound based on measurement of $e$ spectrum distortions in $\bar{\nu}_e e \to \bar{\nu}_e e$ in reactor experiments is $\mu < 1.9 \cdot 10^{-10} \mu_B$ [26]. The interpretation of stellar cooling rate (of He burning stars) via plasmon decay into $\nu \bar{\nu}$ imply a more stringent upper bound $\mu < (1 - 3) \cdot 10^{-12} \mu_B$ [27].

The neutrino spin flavor precession probability can be obtained by solving the following system of equations:

$$i \left( \frac{\dot{\nu}_{eL}}{\nu_{eL}} \right) = \left( \begin{array}{ccc} 0 & \mu B(r) & \Delta m^2 / 2E \\ \mu B(r) & 0 & \Delta m^2 / 2E \\ \Delta m^2 / 2E & \Delta m^2 / 2E & 0 \end{array} \right) \left( \begin{array}{c} \nu_{eL} \\ \nu_{\mu L} \\ \nu_{\tau L} \end{array} \right),$$

which for a constant magnetic field is

$$P(\nu_{eL} \to \bar{\nu}_{\mu R}; r) = \left( \frac{(2\mu B)^2}{(\Delta m^2)^2 + (2\mu B)^2} \right) \cdot \sin^2 \left( \sqrt{\frac{(\Delta m^2)^2 + (2\mu B)^2}{2E}} \right).$$

The amplitude of $P$ is now suppressed unless $\Delta m^2/2E \ll 2\mu B$, where $\Delta m^2 = m_2^2 - m_1^2$ is the mass splitting. This $P$ connects the neutrinos of different flavor and helicity in contrast to the one given by Eq. (8).

A convenient form of the neutrino evolution equation that takes into account not only the effect of neutrino interactions with matter particles in the presence of external magnetic field $B$ but also mass splitting, is

$$i \left( \frac{\dot{\nu}_1'}{\nu_1'} \right) = \left( \begin{array}{ccc} -(M_2^2 - M_1^2)/4E & -i\theta_B & 0 \\ i\theta_B & (M_2^2 - M_1^2)/4E & 0 \\ 0 & 0 & (M_1^2 - M_2^2)/4E \end{array} \right) \left( \begin{array}{c} \nu_1' \\ \nu_2' \\ \nu_3' \end{array} \right).$$

This is a form of the neutrino evolution equations that can be used for studying numerically the propagation of mixed neutrinos [28]. Here $\nu_1' = \text{Exp}[-i(M_1^2 + M_2^2)r/4E]\nu_1$ and $\nu_2' = \text{Exp}[-i(M_1^2 + M_2^2)r/4E]\nu_2$ with

$$M^2_{2,1} = \frac{1}{2} \left[ (m_1^2 + m_2^2 + V_{SF} E) \pm \sqrt{(V_{SF}E - \Delta m^2)^2 + (4E\mu B)^2} \right],$$

where

$$\tan 2\theta_B = \frac{2\mu B}{V_{SF} - \Delta m^2 / 2E}.$$  \hspace{1cm} (14)

Here, $V_{SF} = \sqrt{2}G_F n(2Y_e - 1)$ is the interaction potential for spin flavor conversions, $n$ being the nucleon number density. The level crossing or resonance condition here imply $\theta_B = \pi/4$ or when

$$V_{SF} = \frac{\Delta m^2}{2E}.$$ \hspace{1cm} (15)

Namely, in the resonance, the effective mixing angle attains its maximal value. In the resonance layer, the maximal change in the spin flavor composition of the mixed neutrino state occurs. The following two changes give the description for pure neutrino flavor oscillation in appropriate basis: $\mu B \to \Delta m^2 / 2E \sin 2\theta$ and $V_{SF} \to V_F$ with $V_F = \sqrt{2}G_F nY_e$.

In Eq. (12), if $|\theta_B| \ll |(M_2^2 - M_1^2)/4E|$ or if $2(2\mu B)^2/\pi|V_{SF}| \gg 1$, then the off diagonal terms containing $\dot{\theta}_B$ can be ignored and $\nu_1'$ and $\nu_2'$ become eigenstates of the effective Hamiltonian. In this case, it can be shown that

$$P(\nu_{eL} \to \bar{\nu}_{\mu R}; r) = \frac{1}{2} \left( 1 - \cos 2\theta_B \cos 2\theta_B \right).$$ \hspace{1cm} (16)
where $\theta_i^B$ is initial (at production) and $\theta_f^B$ is the final (at detection) mixing angles. This situation is called the adiabatic approximation. Note that in the adiabatic approximation, the neutrino spin flavor conversion probability depends on initial and final mixing angles only. If the adiabaticity is broken at the level crossing, that is, if $\kappa_{SF} \leq 1$, where

$$\kappa_{SF} \equiv \frac{2(2\mu B)^2}{\pi |V_{SF}|},$$

(17)

then one needs to solve the system of equations from the beginning. Note that the adiabaticity parameter $\kappa_{SF}$ does not depend on $E$ explicitly. The violation of adiabaticity can be parameterized by Landau St"uckelberg Zener probability, $P_{LSZ}$ as [29]

$$P_{LSZ} = \exp \left(-\frac{\pi^2}{4} \cdot \kappa_{SF}\right).$$

(18)

A general expression for neutrino spin flavor conversion probability including the effects of violation of adiabaticity in parameterized form, is

$$P(\nu_eL \rightarrow \bar{\nu}_\mu R; r) = \frac{1}{2} - \left(\frac{1}{2} - P_{LSZ}\right) \cos 2\theta_i^B \cos 2\theta_f^B.$$  

(19)

In summary, the two necessary conditions to obtain a resonant character in neutrino oscillations are the occurrence of level crossing and fulfillment of the adiabaticity condition at the level crossing. The adiabatic approximation imply $P_{LSZ} \rightarrow 0$ in the above Eq.

The details of application of above description for neutrinos from collapsing stars is given in [30]. The reader is referred to these articles for further details. It was pointed out there that, the neutrino spin flavor conversions can occur for $\mu \leq 10^{-13} \mu_B$ in a reasonable strength of magnetic field in the isotopically neutral region of a collapsing star for $10^{-1} \leq \Delta m^2/eV^2 \leq 10^{-8}$. This is because of the peculiar behavior of the effective matter potential, $V_{SF}$ in the isotopically neutral region that a relatively small magnetic field strength is required to get an appreciable spin flavor conversion as compared to that in the Sun at the same distance from the center of the star. The above feature is possible for spin flavor conversions between active neutrinos only [31].

Briefly speaking, in the onion like structure of the progenitor of type II supernovae, below the hydrogen envelope, the layers with mainly isotopically neutral nuclei ($n_n = n_p$) follow such as $^4$He, $^{12}$C, $^{16}$O, $^{28}$Si and $^{32}$S. Thus, the region between the hydrogen envelope and the core is almost isotopically neutral. The deviation from the neutrality is the small abundance of the elements with excess of neutrons ($n_n \neq n_p$) such as in $^{22}$Ne, $^{23}$Na, $^{25}$Mg and $^{56}$Fe. This region is referred to as isotopically neutral region, it extends typically for $10^{-3} \leq r/R_\odot \leq 1$.

The existence of this region follows from the fact that during the collapse, the core and inner region of the star is neutron rich ($n_n > n_p$), whereas the outer region is proton rich, essentially hydrogen envelope ($n_n < n_p$), for a typical type II supernova. Obviously, the neutron and proton densities are almost equal in between ($n_n \sim n_p$), which defines the isotopically neutral region. In terms of $Y_e$, where $Y_e \equiv n_e/(n_n + n_p)$, the supernova phenomena imply $Y_e > 0.5$.
FIG. 1: Some examples of the $\bar{\nu}_e$ flux spectrum distortions in neutrino spin flavor conversions. The original $\bar{\nu}_e$ flux spectrum is shown by the bold solid line. The adiabatic $\bar{\nu}_e$ conversion is shown by the solid line.

in the outer parts, whereas $Y_e < 0.5$ in the inner parts (as electric neutrality of the medium implies that $n_e = n_p$). Therefore, $Y_e \sim 0.5$ in between. As $V_{SF} \propto (2Y_e - 1)$, this implies that $V_{SF}$ passes through very small values in the isotopically neutral region. In other words, it is suppressed relative to $V_F$ up to three orders of magnitude and also changes sign in the isotopically neutral region. This is not the case for Sun for the same distance from the center of the Sun because of the entirely different physics associated with the inner parts of the Sun relative to that in a collapsing star. The presence of the isotopically neutral region depends on the nuclear composition of the star just after the core collapse. It is independent of any external $B$ present in the expanding envelope of the collapsed star. Its presence is also independent of the neutrino intrinsic properties such as $\Delta m^2$, $\mu$ (and $\theta$).

In the case of a direct mass hierarchy ($\Delta m^2 > 0$) and a small flavor mixing with $\mu = B \neq 0$, the main observational signature of a neutrino spin flavor conversion is a distortion of the $\bar{\nu}_e$ flux spectrum, and specially the appearance of a high energy tail. In general, the final $\bar{\nu}_e$ spectrum is the energy dependent combinations of the original $\bar{\nu}_e$ spectrum and the hard spectrum of the non electron neutrinos (see Fig. 1). Another important signature of the spin flavor conversion can be obtained from a comparison of the spectra of different neutrino species. In particular, the $\bar{\nu}_e$ and $\nu_\mu$ spectra can be completely permuted.

On the other hand, in case of pure flavor adiabatic conversions, the above mentioned features are essentially absent in case of Large Mixing Angle (LMA) solution for Solar neutrino problem [19]. Thus, the various features of future supernova neutrino data will possibly help to identify the role of magnetic field in the propagation of mixed system of neutrinos.

The combination of the spin-flip effects with other (flavor) conversions may result in rather peculiar final spectra. For instance, $\nu_e$ may have the spectrum of the original $\bar{\nu}_e$, whereas $\bar{\nu}_e$ may have the original $\nu_\mu$ spectrum. The electron neutrino and anti neutrino spectra can be the same and coincide with the hard spectrum of the original muon neutrinos, etc.
Presently, the only nearby supernova from which neutrinos have been seen is SN 1987 A. Using the fact that no $\bar{\nu}_e$ with average energy greater than $\sim 10$ MeV is seen, an upper bound on the strength of the magnetic field profile for SN 1987A with a fixed $\mu$ value ($\mu \sim 10^{-12} \mu_B$) was obtained under the assumption that $P(\nu_\mu \rightarrow \bar{\nu}_e)$ does not depend on $E$ \[31\]:

$$F_{\bar{\nu}_e} = P(\bar{\nu}_e \rightarrow \bar{\nu}_e) F_{\bar{\nu}_e}^0 + P(\nu_\mu \rightarrow \bar{\nu}_e) \cdot \frac{1}{4} F_{\bar{\nu}_e}^0.$$ \[20\]

The difference among non electron neutrino spectra is ignored here. The upper bound is independent of any magnetic field profile inside the supernova as it is obtained by using Eq. (17) with $\kappa_{SF} = 1$. Thus, the obtained upper bound depends on the profile of $V_{SF}$ only. This is an example of constraining a relevant astrophysical quantity using neutrino observations from the core of a collapsing star.

**D. Prospects for future observations**

In a future nearby supernova occurring, the neutrino signal is expected to be largely dominated by the $\bar{\nu}_e$ flux. This was the case for SN 1987 A also. A reason being that $\sigma(\bar{\nu}_e p \rightarrow ne^+)$ is (at least) an order of magnitude higher than the $\nu_e$ interaction cross section for $5 \leq E/\text{MeV} \leq 50$ in the detector. The $\sigma(\bar{\nu}_e p \rightarrow ne^+)$ is of the order of $10^{-41}$ cm$^2$ for $E \sim 25$ MeV. A future galactic supernova will give several thousand neutrino events in a Super Kamiokande like detector. For completeness, in the following paragraph, I briefly summarize other principal characteristics and limitations of present and future detectors of neutrino bursts from supernovae.

The light water Cherenkov detectors such as Super Kamiokande and Sudbury Neutrino Observatory (SNO) have energy, time and angle resolution for $\nu_e$, $\bar{\nu}_e$ and $\nu_\mu$ ($\nu_\tau$). The neutral current reactions in these detectors have only time resolution. The ice Cherenkov detector such as Antarctic Muon and Neutrino Detector array (AMANDA) can also search for neutrinos from gravitational collapse. The heavy water Cherenkov detector (SNO) can detect $\bar{\nu}_\mu$ and $\bar{\nu}_\tau$ via the reaction $\bar{\nu}_i + d \rightarrow n + p + \bar{\nu}_i$ ($E^\text{th} = 2.22$ MeV) as well with time resolution. The scintillation detectors such as Baksan and Borexino are or will be sensitive to $\nu_e$, $\bar{\nu}_e$, $\nu_\mu$ and $\nu_\tau$ without angle resolution. The drift chamber detector, Imaging of Cosmic and Rare Underground Signals (ICARUS) is or will be sensitive to $\nu_\mu$($\nu_\tau$) with energy, time and angle resolution [although less sensitive to $\nu_\mu$($\nu_\tau$)]. More detailed discussion on prospects for future observations of supernova neutrinos can be found in [32].

In conclusion, all these detectors will in future, collectively provide the temporal, energetic, angular and flavor information for any stellar collapse in our as well as in a nearby galaxy. This information will in turn enable us to constrain the relevant astrophysical quantities such as role of supernova magnetic field strength in mixed neutrino propagation as elaborated in this lecture.
III. LECTURE 2: HIGH ENERGY NEUTRINOS FROM COSMOS

A. Introduction

The neutrinos with $E > 10^6$ MeV are expected to mainly arise from the interaction of ultra high energy cosmic rays considered to be protons ($p$) here with the matter ($p$) and/or radiation ($\gamma$) present in cosmos. Examples of the astrophysical sites where these interaction can occur include the galactic plane, other sites within our galaxy as well as distant sites such as centers of nearby active galaxies (AGNs) and cites for gamma ray bursts (GRBs).

The plan of this lecture is to briefly review the present motivations and status of phenomenological (and experimental) study of these high energy neutrinos. This include a simple classification of presently envisaged main sources, with a description of the main interactions responsible for expected high energy neutrino production. In view of recent growing evidence of neutrino flavor oscillations, I will elaborate the relative changes expected in the high energy neutrino flux because of these neutrino oscillations. I will also describe the basic crucial factors that determine the (limited) near future prospects for observations of these high energy neutrinos. Though, so far there is no observation of neutrinos with energy greater than few thousand MeV, whose origin can not be associated with the Atmosphere of earth, nevertheless, somewhat optimistically speaking, given the current status of high energy neutrino detector developments and the absolute levels of predicted high energy neutrino fluxes, it is expected that possibly the first evidence of high energy neutrinos may come within this decade.

A main motivation of high energy neutrino search is the quest of the microscopic understanding of the nature and origin of observed ultra high energy cosmic rays, namely the presently open questions such as whether they are protons, photons, neutrinos, heavy nuclei such as iron nuclei or some particles suggested beyond the standard model of particle physics, and where and how they are produced or accelerated. A positive observation of high energy neutrinos can raise the possibility of simultaneous explanation of observed high energy photons ($E_{\gamma} \simeq 10^6$ MeV) and ultra high energy cosmic rays as a result of hadron acceleration and interaction in the presently expanding universe.

The neutrinos with energy $> 10^6$ MeV can act as probes of the ultra high energy phenomena observed in the Universe. Unlike photons and charged particles such as protons and heavy nuclei, which can be absorbed or deflected by dust, other intervening matter or magnetic fields, neutrinos can more easily reach the earth because of their weak interactions with matter particles. It is therefore hoped that such neutrinos can provide information about the astrophysical (or/and cosmological) sources that will be complementary to inferences based on visual observations. A better understanding of the interactions involved in neutrino production and a more accurate estimate of resulting neutrino fluxes could entail important consequences. Among these are insights into intrinsic properties of neutrinos such as mass and mixing [33], and the possible role of gravity on neutrino propagation in astrophysical environments [34]. However, it all depends on the existence of a sizable high energy neutrino flux. Assuming an existence of a sizable high energy neutrino flux, several of the other neutrino intrinsic properties as well as the useful information about the source producing these neutrinos can be obtained, at least in principle. These include testing neutrino decay
hypothesis [35], constraints on neutrino magnetic moment [36], quantum gravity effects on neutrino propagation [37],
tests of possible violation of equivalence principle by neutrinos [38], as well as information on different properties of
relic neutrinos [39]. Also possibly enhancement in neutrino nucleon interaction cross section because of various new
physics effects may be constrained [40]. An early attempt to constrain the neutrino nucleon interaction cross section
is discussed in [41].

The relevant average physical picture in AGNs is as follows. Some galaxies have quite bright centers. The photon
luminosity of these galaxies typically reach \((10^{44} - 10^{48})\ \text{erg s}^{-1}\). These galaxies are typically several Mpc away from
us. In general, AGNs refer to these bright and compact central regions, which may extend up to several pc in the
center. These central compact regions have the remarkable property of being much more luminous than the rest of
the entire galaxy. It is hypothesized that the existence of a super massive black hole with mass, \(M_{\text{BH}} \sim (10^6 - 10^{10})\ M_\odot\),
may explain the observed brightness as this super massive black hole captures the matter around it through accretion.
This super massive black hole is presently hypothesized to be formed by the collapse of a cluster of stars. Some AGNs
give off a jet of matter that stream out from the central compact region in a transverse plane and produce hot spots
when the jet strikes the surrounding matter at its other ends. During and after accretion, the (Fermi) accelerated
protons may collide with other protons and/or with the ambient photons in the vicinity of an AGN or/and in the
associated jets/hot spots to produce unstable hadrons. These unstable hadrons decay mainly into neutral and charged
pions. The neutral pions further decay dominantly into photons and thus may explain a large fraction of the observed
brightness, whereas the charged pions mainly decay into neutrinos. AGNs, therefore, have been targeted as one
likely source of high energy neutrinos. Currently, the photohadronically \((p\gamma)\) produced flux of high energy neutrinos
originating from AGNs dominate over the flux from other sources above the relevant Atmospheric background typically
for \(E \geq 10^9\ \text{MeV}\) [42, 43]. For further reading on astrophysical super massive black holes, see [44].

Recently, fireballs are suggested as a possible production scenario for gamma ray bursts as well as high energy
neutrinos at the site [45]. Though, the origin of these gamma ray burst fireballs is not yet understood, the observations
suggest that generically a very compact source of linear scale \(\sim 10^7\ \text{cm}\) through internal or/and external shock
propagation produces these gamma ray bursts (as well as burst of high energy neutrinos) mainly in \(p\gamma\) interactions.
Typically, this compact source is hypothesized to be formed possibly due to merging of binary neutron stars or due
to collapse of a super massive star. Thus, fireballs have also been suggested as a probable scenario for the observed
gamma ray bursts, and they too are expected to emit neutrinos with energies in excess of hundreds of thousands of
MeV. For a recent review, see [46].

A nearby and more certain source of high energy neutrinos is our galactic plane. The incoming ultra high energy
cosmic ray protons interact with the ionized hydrogen clouds there and can produce high energy neutrinos in \(pp\)
interactions. Present estimates indicate that the diffuse galactic plane muon neutrino flux can dominate over the
Atmospheric one for \(E > 10^8\ \text{MeV}\).
TABLE I: Comparison of the cross sections for the three high energy neutrino production interactions discussed in the text at √s ∼ 1.2 · 10^3 MeV.

| Interaction          | σ(mb)  |
|----------------------|--------|
| pγ → Nπ±             | ≤ 5 · 10^{-4} |
| pp → Nπ±             | ∼ 3 · 10^3   |
| γγ → µ⁺µ⁻           | < 10^{-3}   |

B. Expected neutrino production

A presently favorable astrophysical scenario for high energy neutrino production is that the observed ultra high energy cosmic rays beyond GZK cutoff (see later) are dominantly protons and that the observed high energy photon flux can be associated with these. On the other hand, an unfavorable scenario is that the ultra energy cosmic rays are dominantly other than protons and that the observed high energy photon flux has purely electromagnetic origin. In the latter case, there will still be neutrino flux but at a rather suppressed level (such as in γγ interactions) as compared to the former case. The latter possibility is recently discussed in some detail in [47].

The main interactions responsible for the production of these high energy neutrinos include the pγ and pp interactions (see Table I). For the behavior of these cross sections as a function of center-of-mass energy √s in the range of interest, see [26]. There is formation of Δ resonance in pγ interactions, at √s ∼ m_Δ ∼ 1.2 · 10^3 MeV, that mainly decay into electron and muon neutrinos. Two behaviors of the pγ cross section, near √s ∼ 1.2 · 10^3 MeV make it an important channel for high energy neutrino production, the relatively large width of the Δ resonance, Γ_Δ/m_Δ ∼ 10^{-2}, and the almost constant behavior of the cross section for √s > m_Δ + Γ_Δ. Under the assumption of all other similar conditions, it is the interaction cross section that determines the absolute level of high energy neutrino production.

For illustrative purpose, Fig. 2 displays a simple classification flow chart for presently envisaged sources of high energy neutrinos. It includes the possibility of high energy neutrino production from cosmic relics, referred to as X [48]. Briefly, these relics are considered to be formed in the early epochs of the universe such as during inflation epoch. The large amount of energy trapped in these relics may be released in the form of grand unification scale gauge bosons which in turn decay/annihilate into standard model particles including neutrinos. These relics need not be far away from us. In fact, some of the models suggest that they may be a part of our galactic dark matter halo implying at a distance of ≤ 10 kpc. If these X’s can be the dominant sources of observed ultra high energy cosmic rays then this in turn severely constrain their number density n_X, life time τ_X, mass M_X, and thus determine the resulting high energy neutrino flux spectrum shape and absolute level. This possibility is referred to as the cosmological scenario for expected high energy neutrino production. Currently, the ultra high energy cosmic rays with energy E_{UHECR}^{max} up to ∼ 3 · 10^{14} MeV are observed [49].

Depending on the details of the astrophysical or cosmological model for high energy neutrino production scenario, either the observed photon flux or proton flux or both are used to determine the absolute level of the expected
Typically, an acceleration mechanism for cosmic ray protons is required with a power law flux spectrum, $F_p(E) \propto E^{-\zeta}$ with $\zeta \sim 2$. In general, astrophysical source(s) should be at a distance of $\leq 50$ Mpc.

Here, no acceleration mechanism is required for cosmic ray protons. If the cosmo relic X’s can be a dominant source for ultra high energy cosmic rays (UHECR), then it requires that $M_X c^2 \geq E_{\text{max, UHECR}}$.

FIG. 2: A simple classification flow chart for presently envisaged main sources of high energy neutrinos. Only non tau neutrino production is illustrated.

neutrinos flux. From the cosmos, presently high energy photons and ultra high energy cosmic rays (considered to be protons here) are observed in the relevant context. Their observed level of flux determines the absolute flux level of neutrinos as high energy neutrinos are secondary in nature in the sense that they are not matter particles and are not a significant fraction of the matter density associated with a specific known astrophysical or and cosmological source.
On the other hand, neutrinos are stable and neutral and therefore for this precise reason will carry useful information about the source. Supposing protons can escape the extra galactic astrophysical sources and can be a dominant fraction of the observed ultra high energy cosmic ray flux, the resulting high energy (muon) neutrino flux mainly in \( p\gamma \) and \( pp \) interactions either arising from inside the source or during propagation has to be less than this. It can be typically \( \leq 10^{-5}\) MeV(cm\(^2\)·s·sr\(^{-1}\)) for \( 10^8 < E/\text{MeV} < 10^{15} \). This bound further tightens by a factor of 1/2 once the neutrino flavor oscillation effects are taken into account (see later).

Consider now briefly the \( p\gamma \rightarrow \Delta \rightarrow p\pi \) \((N=p)\) interactions occurring during the propagation of ultra high energy cosmic rays either inside an astrophysical source or between the source and the earth in the presence of a dense photon background. This is to serve as an illustrative example for having an order of magnitude idea of the expected \( E \). The threshold energy for protons interacting at an angle \( \phi \) to form \( \Delta \) resonance, is

\[
E_{th}^p = \frac{(m_p + m_\pi)^2 - m_p^2}{2E_\gamma (1 - \cos \phi)},
\]

which in case of head on interactions further simplifies to

\[
E_{th}^p \simeq \frac{m_p m_\pi}{2E_\gamma}.
\]

For \( E_p < E_{th}^p \), the interaction \( p\gamma \rightarrow pe^+e^- \) dominates the energy loss for protons. If \( E_\gamma = E_{CMB}^\gamma \sim 2.7 \) K then \( E_{th}^p \sim 10^{14} \) MeV. The \( p\gamma \) interaction length can be defined as

\[
\lambda \sim 1/n_\gamma \sigma_{p\gamma\rightarrow p\pi}.
\]

For instance, if \( n_\gamma = n_{CMB}\gamma \sim 410 \) cm\(^{-3}\) for \( E_\gamma = E_{CMB}^\gamma \) then \( \lambda < 6 \) Mpc, where \( \sigma_{p\gamma\rightarrow p\pi} \) is given in Table I.

The propagation of ultra high energy proton flux, \( F_p \) can be studied in the presence of photon background in distance \( r \), by solving the following equation

\[
\frac{dF_p}{dr} = -\frac{1}{\lambda} F_p.
\]

The negative sign indicates the decrease in the ultra high energy proton flux because of interaction described by \( \lambda \). This results in an exponential cut off in \( p \) flux spectrum. In case of ubiquitous cmb photon background, it is commonly referred to as Greisen Zatsepin Kuzmin (GZK) cut off [51]. It occurs at \( E_{th}^p \sim 10^{14} \) MeV, according to Eq. (22). The resulting GZK (muon) neutrino flux spectrum peaks at \( \sim 10^{12} \) MeV by sharing roughly \((1/4)\cdot(1/5)\) of the \( E_{th}^p \).

The matter density in interstellar medium as well as in several of the astrophysical sites such as the galactic plane, the AGNs and the GRBs, is rather small (relative to that in Atmosphere of earth). Therefore, a rather simple formula can be used to estimate high energy neutrino flux spectrum in \( p\gamma \) and/or \( pp \) in a specific individual astrophysical site

\[
F_\nu^0(E) = \int_E^{E_{max}} dE \int F_p(E) g(E) \frac{dn_{p(\gamma,p)\rightarrow \nu Y}}{dE}.
\]

Here \( F_p(E) \) parameterizes the high energy proton flux. The function \( g(E) \equiv r/\lambda \) gives the number of \( p(\gamma,p) \) interactions within the distance \( r \). The \( dn/dE \equiv \sigma^{-1}d\sigma/dE \) is the neutrino energy distribution in above interactions. The
implicit assumption here is that the unstable hadrons and leptons produced in above interactions decay before they interact owing to the fact that the matter density in the distance \( r \) is assumed to be rather small. Also, the effects of possible red shift evolution and magnetic field of the astrophysical sources are neglected for simplicity.

There is yet another possible class of astrophysical sources of high energy neutrinos that are essentially neither constrained by observed high energy photon nor by ultra high energy cosmic ray flux. It is so because in this class of sources, the matter density is considered to be too large so that neither of the above leave the source. These sources are therefore commonly referred to as hidden sources or neutrinos only sources. The high energy neutrino production occurs in same \( pp \) (or \( p\gamma \)) interactions here also. These can only be constrained by the high energy neutrino flux (non) observations [52].

The above discussion is restricted to non tau neutrino production only. In the \( \pi^\pm \rightarrow \mu^\pm \rightarrow \nu \) decay situation, the relative ratio of resulting electron and muon neutrino flux is 1 : 2 respectively. The astrophysical tau neutrino flux is produced in decays of \( D^+_S \). For \( \sqrt{s} \sim m_\Delta \), it is known that \( \sigma[p(\gamma,p) \rightarrow D^+_S Y]/\sigma[p(\gamma,p) \rightarrow \pi^+ Y] \leq O(10^{-3} - 10^{-4}) \). The high energy tau neutrino flux is thus rather suppressed at the production sites and can therefore be taken as approximately zero, resulting in 1 : 2 : 0 [53]. For a recent review on astrophysical tau neutrinos, see [54], whereas for cosmological tau neutrinos, see, for instance [55].

C. Oscillations during propagation: Effects of neutrino mixing

There are at least two aspects of neutrino propagation effects that need somewhat careful considerations in study of neutrino mixing effects for high energy neutrinos. These are: the neutrino interactions with the background particles inside the (astrophysical) source of neutrinos as well as between the source and the earth. The present knowledge of matter density, \( \rho \) inside the known sources as well as between these sources and the earth imply that it is rather quite small (as compared to that in Sun). As a result, the level crossing condition, \( G_F \rho/m_N \sim \Delta m^2/2E \), for matter enhanced neutrino flavor oscillations is not satisfied [see also, for instance, Eq. (15)]. Level crossing is a necessary condition for occurrence of matter enhanced neutrino flavor oscillations. Therefore, there are essentially no matter effects on pure vacuum flavor oscillations. Note that this is in contrast to the situation in supernovae. Furthermore, the neutrino nucleon and neutrino electron inelastic interaction effects are also small enough to effect the mixed neutrino propagation even at ultra high energy in a significantly observable manner. This is also because of rather small matter density. Therefore, I elaborate only effects of neutrino flavor mixing in vacuum (with no matter interactions)\(^1\).

Note from the previous subsection that the high energy neutrinos are produced in the following relative ratios

\[
F^{0}_{\nu_e} : F^{0}_{\nu_\mu} : F^{0}_{\nu_\tau} = 1 : 2 : 0.
\]  

---

\(^1\) If i) \( 0.1 \leq \sin^2 2\theta \leq 0.95 \), ii) \( E \geq 10^{12} \) MeV, iii) the red shift \( z \geq 3 \) at production, and iv) \( \xi \geq 1 \), where \( \xi \equiv (n_{\nu} - n_{\bar{\nu}})/n_{\gamma} \), then a deviation from pure vacuum flavor oscillations can be of the order of few percent, when high energy neutrinos scatter over the very low energy relic neutrinos during their propagation to us in the interstellar medium [56].
It is assumed here that the high energy neutrinos and anti neutrinos originate in equal proportion from a source and are counted in the symbol $\nu$ together. Also as the absolute level of high energy neutrino flux is presently unknown, I therefore elaborate the neutrino mixing effects on relative ratios only, in the context of three flavors. Four flavor mixing effects are considered in [33, 57].

To obtain a general expression for flavor oscillation formula, I start with the connection $U$ between the flavor $|\nu_\alpha\rangle$ and mass $|\nu_i\rangle$ eigen states of neutrinos, namely

$$|\nu_\alpha\rangle = \sum_{i=1}^{3} U_{\alpha i} |\nu_i\rangle,$$

where $\alpha = e, \mu$ or $\tau$. In the context of three neutrinos, $U$ is called Maki Nakagawa Sakita (MNS) mixing matrix [58]. It can be obtained by performing the following operations to coincide with the one given in [26]:

$$U \equiv R_{23}(\theta_{23}) \cdot \text{diag}(e^{-i\delta_{13}/2}, 1, e^{i\delta_{13}/2}) \cdot R_{13}(\theta_{13}) \cdot \text{diag}(e^{i\delta_{13}/2}, 1, e^{-i\delta_{13}/2}) \cdot R_{12}(\theta_{12}),$$

where $\theta$’s are neutrino mixing angles and $\delta_{13}$ is CP violation phase. Explicitly, it reads

$$U = \begin{pmatrix}
    c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta_{13}} \\
    -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta_{13}} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta_{13}} & s_{23}c_{13} \\
    s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta_{13}} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta_{13}} & c_{23}c_{13}
\end{pmatrix}. \quad (29)$$

Here $c_{ij} = \cos \theta_{ij}$ and $s_{ij} = \sin \theta_{ij}$ (with $j = 1, 2, 3$) and that $UU^\dagger = 1$. Using it, one obtains the following well known formula for flavor oscillation probability from $\alpha$ to $\beta$ ($\beta = e, \mu$ or $\tau$) for a neutrino source at a fixed distance $L$:

$$P(\nu_\alpha \rightarrow \nu_\beta; L) \equiv P_{\alpha\beta} = \delta_{\alpha\beta} - \sum_{j \neq k} U^\ast_{\alpha j} U_{\beta j} U^\ast_{\alpha k} U_{\beta k} (1 - e^{-i\Delta m^2_{jk} L/2E}). \quad (30)$$

In the far distance approximation, namely, in the limit $L \rightarrow \infty$, one obtains

$$P(\nu_\alpha \rightarrow \nu_\beta; L \rightarrow \infty) \simeq \delta_{\alpha\beta} - \sum_{j \neq k} U^\ast_{\alpha j} U_{\beta j} U^\ast_{\alpha k} U_{\beta k},$$

$$\simeq \sum_{j=1}^{3} |U_{\alpha j}|^2 |U_{\alpha j}|^2. \quad (31)$$

Because of the assumed averaging over the rapidly oscillating phase ($l_{\text{osc}} \ll L$, where $l_{\text{osc}} \equiv 2E/\Delta m^2_{jk}$), the last two expressions are independent of $E$ and $\Delta m^2$. Under this assumption, the oscillation probability can be written as a symmetric matrix $P$ and $P$ can be written as a product of a matrix $A$:

$$P = \begin{pmatrix}
    P_{ee} & P_{e\mu} & P_{e\tau} \\
    P_{e\mu} & P_{\mu\mu} & P_{\mu\tau} \\
    P_{e\tau} & P_{\mu\tau} & P_{\tau\tau}
\end{pmatrix} \equiv AA^T, \quad (32)$$

with

$$A = \begin{pmatrix}
    |U_{e1}|^2 & |U_{e2}|^2 & |U_{e3}|^2 \\
    |U_{\mu1}|^2 & |U_{\mu2}|^2 & |U_{\mu3}|^2 \\
    |U_{\tau1}|^2 & |U_{\tau2}|^2 & |U_{\tau3}|^2
\end{pmatrix}. \quad (33)$$
A simple form for $P$ matrix can be obtained in case of vanishing $\delta_{13}$ and $\theta_{13}$ with bi maximal mixing \cite{59}:

$$P = \begin{pmatrix}
\frac{5}{8} & \frac{3}{16} & \frac{3}{16} \\
\frac{3}{16} & \frac{13}{32} & \frac{13}{32} \\
\frac{3}{16} & \frac{13}{32} & \frac{13}{32}
\end{pmatrix}. \tag{34}$$

This $P$ matrix satisfies the following unitarity conditions:

$$1 - P_{ee} = P_{e\mu} + P_{e\tau}, \quad 1 - P_{\mu\mu} = P_{e\mu} + P_{\mu\tau}, \quad 1 - P_{\tau\tau} = P_{e\tau} + P_{\mu\tau}, \tag{35}$$

namely, the disappearance of a certain neutrino flavor is equal to the appearance of this flavor into other (active) neutrino flavors. High energy neutrino flux arriving at the earth can be estimated using

$$F_{\nu\alpha} = \sum_{\beta} P_{\alpha\beta} F_{\nu\beta}, \tag{36}$$

where $P_{\alpha\beta}$ is given by Eq. (31). Note that in case of initial relative flux ratios as 1 : 2 : 0 \cite{20}, one always get

$$F_{\nu_e} : F_{\nu_\mu} : F_{\nu_\tau} = 1 : 1 : 1, \tag{37}$$

under the assumption of averaging irrespective of any specific flavor oscillation solution for Solar neutrino problem \cite{60}. A considerable enhancement in $F_{\nu_\tau}$ relative to $F_{\nu_\beta}$ because of neutrino oscillations is evident. A some what detailed numerical study that takes into account the effects of non vanishing $\delta_{13}$ and $\theta_{13}$ indicates that the deviation from these final relative ratios is not more than few percent (namely, $|\epsilon| \leq 0.1$ in $1 \pm |\epsilon|$) \cite{33}. There could, in principle, be several intrinsic neutrino properties that may lead to deviations from 1 : 1 : 1 final relative ratios other than $|\epsilon|$ as well as an energy dependence, such as neutrino spin flavor conversions \cite{61}. Astrophysical/cosmological reasons at the source can also contribute to these deviations.

In the above simplified discussion, the expression for $P$ neither depends on $\Delta m^2$ nor on $E$. However, in some situations, this need not be the case. In that case, one need to use complete expression for $P$ given by Eq. (30) and have to average over the red shift distribution of astrophysical sources, $f(z)$. This gives the effect of evolution of the sources with respect to $z$. This effect can be calculated using the $P$ given in Eq. (30) with $E \rightarrow (1 + z)E$ in following formula

$$P_{\alpha\beta}(E) = \frac{\int_0^{z_{\text{max}}} P_{\alpha\beta}(E, z)f(z)dz}{\int_0^{z_{\text{max}}} f(z)dz}. \tag{38}$$

The $f(z)$ can be found in \cite{56}.

**D. Prospects for possible future observations**

The current status of the dedicated high energy neutrino detectors is given in \cite{62}. Briefly, the detectors based on Cherenkov radiation measurement, in ice or water are the Antarctic Muon and Neutrino Detector Array (AMANDA)
and its proposed extension, the Ice Cube, the lake Baikal detector and the Astronomy with a Neutrino Telescope and Abyss environmental RESearch (ANTARES) detector array. The hybrid detectors based on particle and radiation measurement such as Pierre Auger Observatory can also detect high energy neutrinos \[63\]. Detectors based on alternative detection techniques such as radio wave detection are also in operation, such as Radio Ice Cherenkov Experiment (RICE). This detector is based on Askaryan effect. This effect is briefly defined as follows: In an electromagnetic shower generated in deep inelastic neutrino nucleon interaction, the electrons and photons in the shower generate an excess of \(\sim 10 - 20\%\) electrons in the shower because of the electron and photon interactions with the medium in which the shower develops. This in turn generate coherent radio wave pulse (in addition to other type of radiation), if the wavelength of this radio emission is greater than the size of the shower. The search for alternative high energy neutrino detection medium other than air, water and/or ice, such as rock salt has also been attempted for radio wave emission \[64\]. It might also be possible to detect the acoustic pulses generated by deep inelastic neutrino nucleon interactions near or inside the detector. An attempt in this direction is through Sea Acoustic Detection of Cosmic Objects (SADCO) detector array. Other proposals include space based detectors such as Orbiting Wide Angle Light collector (OWL/Air Watch) and Extreme Universe Space Observatory (EUSO).

Among all these, the tightest upper bounds on high energy neutrino flux are reported by AMANDA and Baikal detectors. From AMANDA, this is typically \(\leq 10^{-3}\) MeV \((cm^2 \cdot s \cdot sr)^{-1}\) in the energy range \(10^6 < E/\text{MeV} < 10^9\). The effective area for AMANDA detector is \(\sim 10^{-2}\) km\(^2\) for a \(10^7\) MeV muon neutrino. The next generation high energy neutrino detectors are considered to have an effective area of \(\sim 1\) km\(^2\).

The high energy neutrino observation can be achieved in the following two main interactions: the deep inelastic neutrino nucleon and neutrino electron interactions. The deep inelastic neutrino nucleon interaction can proceed via \(Z\) or \(W^\pm\) exchange. The former is called neutral current (NC) interaction, whereas the later is called charged current (CC) interaction. The CC interactions (\(\nu_\alpha N \rightarrow \alpha Y\)) are most relevant for prospective high energy neutrino observations. The showers, the charged particles and the associated radiation emission such as Cherenkov radiation from these interactions are the measurable quantities. The CC deep inelastic neutrino nucleon cross section \(\sigma^{CC}(E)\) over nucleons with mass \(m_N\), can be written as a function of the incoming neutrino energy \(E\) as:

\[
\sigma^{CC}(E) = \frac{2 G_F^2 m_N E}{\pi} \int dx \int dy \left( \frac{M_W^2}{Q^2 + M_W^2} \right)^2 x \cdot y \cdot \left\{ \left(1 - \frac{m_N^2 x^2 y^2}{Q^2} \right) \left[ f_d(x, Q^2) + f_s(x, Q^2) + f_u(x, Q^2) \right] + \left(1 - y^2 - \frac{m_N^2 x^2 y^2}{Q^2} \right) \left[ f_d(x, Q^2) + f_s(x, Q^2) + f_l(x, Q^2) \right] \right\} . \tag{39}
\]

The integration limit for \(x\) and \(y\) can be taken between 0 and 1. This expression can be straightforwardly obtained using \(s = 2m_N E\) in \[24\]. Here \(f_i(x, Q^2)\) are the parton distribution functions. In above Eq., \(y \equiv (E - E')/E\) is the inelasticity in the neutrino nucleon interactions. It gives the fraction of \(E\) lost in a single neutrino nucleon interaction. The \(x \equiv Q^2 / 2m_N (E - E')\) is the fraction of the nucleon’s momentum carried by the struck quark. The lepton mass
20

\[
\begin{align*}
\sigma(\bar{\nu}_e e \rightarrow W^- \rightarrow \text{hadrons}) &= \frac{\Gamma_W (\text{hadrons})}{\Gamma_W (e^+ \nu)} \frac{G_F^2 s}{3\pi} \left[ \frac{M_W^4}{(s - M_W^2)^2 + \Gamma_W^2 M_W^2} \right], \\
\sigma(\bar{\nu}_e e \rightarrow W^- \rightarrow \gamma) &= \frac{\sqrt{2} \alpha G_F}{3u^2(u-1)} \left[ 3(u^2 + 1) \ln \left( \frac{(u-1)M_W^2}{m_e^2} \right) - (5u^2 - 4u + 5) \right],
\end{align*}
\]

where \(\Gamma_W\)'s can be found in [26]. The above resonant interaction select the anti electron neutrino flavor as well as the energy. This interaction can, in principle, be used to calibrate the high energy neutrino energy provided it can possibly be discriminated from neutrino nucleon interaction in a detector. The \(\sigma(\bar{\nu}_e e \rightarrow W^- \rightarrow \text{hadrons})\) has a slight enhancement because of hard photon emission in the final state for \(\sqrt{s} \geq \Gamma_W\). It is given by [65]

\[
\sigma(\bar{\nu}_e e \rightarrow W^- \gamma) = \frac{\sqrt{2} \alpha G_F}{3u^2(u-1)} \left[ 3(u^2 + 1) \ln \left( \frac{(u-1)M_W^2}{m_e^2} \right) - (5u^2 - 4u + 5) \right],
\]

where \(u = s/M_W^2\). In Fig. 3, the three cross sections are plotted for illustration. The \(\sigma^{CC}(\bar{\nu}_e N)\) is calculated using CTEQ(5M) parton distribution functions generated by Coordinated Theoretical and Experimental Project on QCD Phenomenology and Tests of the Standard Model [66].

The high energy neutrino flux arrives at an earth based detector in three general directions in equal proportion. The downward going neutrinos do not cross any significant earth cord before reaching the detector. The horizontal and upward going neutrinos cross the earth with increasing cord length respectively.

The event rate for downward going high energy neutrinos in CC deep inelastic interactions is given by [67]

\[
\text{Rate} = A \int_{E_{\nu,\alpha}^{\text{min}}}^{E_{\nu,\alpha}^{\text{max}}} dE_{\nu,\alpha} P_{\nu,\alpha \rightarrow \alpha}(E_{\nu,\alpha}, E_{\nu,\alpha}^{\text{min}}) F_{\nu,\alpha},
\]

here \(A\) is the area of the high energy neutrino detector. The \(F_{\nu,\alpha}\) can be obtained using Eq. (36). In the above equation

\[
P_{\nu,\alpha \rightarrow \alpha}(E_{\nu,\alpha}, E_{\nu,\alpha}^{\text{min}}) = N_A \int_0^{1 - E_{\nu,\alpha}^{\text{min}}/E_{\nu,\alpha}} dy R_{\alpha}(E_{\nu,\alpha}, E_{\nu,\alpha}^{\text{min}}) \frac{d\sigma^{CC}_{\nu,\alpha N}(E_{\nu,\alpha}, y)}{dy}.
\]
FIG. 4: Left panel: Expected downward going $e$–like, $\mu$–like and $\tau$–like event rate produced by AGN neutrinos as a function of minimum energy of the corresponding charged lepton in a large km$^3$ volume ice or water neutrino detector. Three flavor neutrino mixing is assumed. Right panel: Approximate representative event topologies for the three neutrino flavors in a km$^3$ volume water or ice neutrino detector for the order of magnitude energy interval shown in left panel.

The $d\sigma/dy$ can be obtained using Eq. (39). The $N_A$ is the Avogadro’s constant. Various $R$’s are given in Table II. Note that for $\tau$ lepton, it is the decay length that is considered as its range with $E_{\tau}^{\text{min}} \sim 2 \cdot 10^9$ MeV and $E_{\nu_\tau}^{\text{max}} \sim 2 \cdot 10^{10}$ MeV as the value of $D$ is chosen as $10^5 \text{ cm}$ for illustration here [68]. Also note that $R_e \equiv R_e(E)$ only.

For upward going high energy neutrinos, a shadow factor $S$ is included in the integral given by Eq. (43). The shadow factor $S$ takes into account the effects of absorption by earth [67]. The absorption of upward going high energy neutrinos by earth is neutrino flavor dependent. For $E \geq 10^9$ MeV, the upward going tau neutrinos may reach the surface of the earth in a relatively small number by lowering their energy so that $E < 10^9$ MeV [69], whereas the upward going electron and muon neutrinos are almost completely absorbed by the earth. For further details, see [67]. For downward going high energy neutrinos, the $S$ is taken as unity. Detailed estimates of the high energy neutrino event rates are done mainly numerically. These estimates are model dependent. The event rates of downward high energy neutrinos typically vary between $\mathcal{O}(10^1)$ and $\mathcal{O}(10^2)$ in units of (yr sr)$^{-1}$ for the proposed km$^3$ volume ice or water neutrino detector. The left panel of Fig. 4 displays the three downward going event rates along with examples of event topologies for AGN neutrinos [70]. In this AGN model, the $pp$ interactions inside the

| Lepton Flavor | $R(\text{cmwe})$ |
|---------------|------------------|
| $e$           | $\frac{40}{(1 - (y(E))}\frac{E}{\text{cmwe}}}$ |
| $\mu$         | $\frac{1}{a+bE_\mu}f_{\mu}$, $a = 2 \cdot 10^{-3}$, $b = 3.9 \cdot 10^{-6}$ |
| $\tau$        | $D = \frac{E(1-y)_\tau}{D}$, $D = 10^5 \text{ cm}$ |

TABLE II: The three charged lepton ranges discussed in the text.
core of the AGN are considered to play an important role. The $e^-$-like event rate are obtained by re-scaling the $\mu^-$-like event rate. The indicated order of magnitude energy interval is relevant for proposed $\text{km}^3$ volume high energy neutrino detectors in the context of possible neutrino flavor identification.

The downward going high energy neutrinos of different flavors interact with the medium (free nucleons) of the detector, deep inelastically mainly through CC interactions. The three flavors on the average give rise to different event topologies based on these interactions and the behavior of the associated charged lepton. For instance, for $E \geq 10^9$ MeV, in proposed $\text{km}^3$ volume ice or water neutrino detectors, typically the downward going high energy electron neutrinos produce a single shower, the muon neutrinos produce muon like tracks passing through the detector (along with a single shower), whereas the tau neutrinos produce two hadronic showers connected by muon like track and is such that the amplitude of the second shower is essentially a factor of two larger as compared to the first. Here, amplitude refers to maximum number of charged particles per unit length (see right panel of Fig. 4) \cite{68, 71}. For a recent discussion on prospects for observations of near horizontal (tau) neutrinos, see \cite{72}.

IV. SUMMARY AND CONCLUSIONS

Detailed study of neutrino fluxes from different astrophysical sites such as the cores of collapsing stars, the galactic plane, as well as other more far away anticipated astrophysical and cosmological sites will provide valuable information about the neutrino intrinsic properties and the site itself.

In collapsing stars, such studies can constrain the role of supernova magnetic field in mixed neutrino propagation. In particular, if a neutrino spin flavor conversion occurs in the isotopically neutral region in a collapsing star, the resulting changes in neutrino flux spectrum are sensitive to rather small values of neutrino magnetic moment, $\mu$, such as $\mu \leq 10^{-13} \mu_B$. In this context, in the first lecture, after a brief introduction of the main neutrino production mechanisms during the supernova stage, the description for neutrino spin flavor conversions is elaborated under the assumption of small vacuum mixing. This includes a discussion that incorporates the possibility of violation of adiabaticity in neutrino spin flavor conversions. Results of its implications for supernova neutrinos are summarized.

For other more distant and energetic sources, the prospective high energy neutrino observation will provide clues for a solution of the long standing problem of origin of observed ultra high energy cosmic rays. The (non) observation of high energy neutrinos will also help to better model the underlying physics of the far away astrophysical and cosmological sites. In this context, in the second lecture, the present motivations for their searches are presented, with a description of their possible connection with ultra high energy cosmic rays. Main high energy neutrino production mechanisms are summarized via a simple classification flow chart. Three flavor neutrino oscillation description is reviewed and its implications for the three relative ratios of high energy neutrinos are given. Furthermore, the essentials of prospective observations of high energy neutrinos are briefly described including the possible relevant observational signatures.

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The main steps consist in deriving the $P_{LSZ}$ are to start from neutrino evolution equation of the type given in Eq. (10) for $\nu_e$ and $\bar{\nu}_e$ and then using appropriate boundary conditions obtain a second order differential equation for $\bar{\nu}_e$ (see a similar consideration for pure flavor case in [1]). Assuming that the matter density and magnetic field vary linearly in the $\nu$ for $\nu$ $\nu$ with $\eta \neq 0$ and $\dot{\eta} \neq 0$ with $\dot{\eta} \neq 0$. Remaining variables are: $l = \sqrt{\eta r} \exp(-i \frac{\eta r}{2})$ and $n = \frac{l^2}{r^2}$, where $\epsilon_{12} = 2\mu B$. Choice of the solution with correct asymptotic
behavior of this differential equation gives $\bar{\nu}_\mu$ which in turn can be shown to lead to exponential form of the probability that $\bar{\nu}_\mu$ stays as $\bar{\nu}_\mu$ after crossing the resonance region, given by Eq. (18).

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