astro-ph/0603372 FERMILAB–CONF–06/029–T

EXTREMELY HIGH ENERGY COSMIC NEUTRINOS AND RELIC NEUTRINOS

CHRIS QUIGG

Fermi National Accelerator Laboratory
P.O. Box 500, Batavia, Illinois 60510 USA
and
CERN, CH-1211 Geneva 23, Switzerland
E-mail: quigg@fnal.gov

ABSTRACT

I review the essentials of ultrahigh-energy neutrino interactions, show how neutral-current detection and flavor tagging can enhance the scientific potential of neutrino telescopes, and sketch new studies on neutrino encounters with dark matter relics and on gravitational lensing of neutrinos.

1. Neutrino Observatories: Expectations

An early goal of the next generation of neutrino telescopes will be to detect the flux of cosmic neutrinos that we believe will begin to show itself above the atmospheric-neutrino background at energies of a few TeV. A short summary of the science program of these instruments is to prospect for cosmic-neutrino sources, to characterize the sources, to study neutrino properties, and to be sensitive to new phenomena in particle physics. The expected sources include active galactic nuclei (AGN) at typical distances of roughly 100 Mpc.[a] If neutrinos are produced there in the decay of pions created in pp or $p\gamma$ collisions, then we anticipate—at the source—equal numbers of neutrinos and antineutrinos, with a flavor mix $2\gamma + 2\nu_\mu + 2\bar{\nu}_\mu + 1\nu_e + 1\bar{\nu}_e$, provided that all pions and their daughter muons decay. I denote this standard flux at the source by $\Phi^0_{\mathrm{std}} = \{\varphi^0_e = \frac{1}{3}, \varphi^0_\mu = \frac{2}{3}, \varphi^0_\tau = 0\}$.

We expect that a neutrino observatory with an instrumented volume of 1 km^3 will be able to survey the cosmic-neutrino flux over a broad range of energies, principally by detecting the charged-current interaction $(\nu_\mu, \bar{\nu}_\mu)N \to (\mu^-, \mu^+) + \text{anything}$. Important open questions are whether we can achieve efficient, calibrated $(\nu_e, \bar{\nu}_e)$ and $(\nu_\tau, \bar{\nu}_\tau)$ detection, and whether we can record and determine the energy of neutral-current events. One of my aims in this talk will be to illustrate how adding these capabilities will enhance the scientific potential of neutrino observatories.

The cross section for deeply inelastic scattering on an isoscalar nucleon may be

[a] 1 Mpc $\approx 3.1 \times 10^{22}$ m.

written in terms of the Bjorken scaling variables $x = Q^2/2M\nu$ and $y = \nu/E_\nu$ as

$$\frac{d^2\sigma}{dxdy} = \frac{2G_F^2 M E_\nu}{\pi}\left(\frac{M_W^2}{Q^2+M_W^2}\right)^2\left[xq(x,Q^2) + x\bar{q}(x,Q^2)(1-y)^2\right] , \tag{1}$$

where $-Q^2$ is the invariant momentum transfer between the incident neutrino and outgoing muon, $\nu = E_\nu - E_\mu$ is the energy loss in the lab (target) frame, M and M_W are the nucleon and intermediate-boson masses, and $G_F = 1.16632 \times 10^{-5}$ GeV^{-2} is the Fermi constant. The parton densities are

$$\begin{aligned} q(x,Q^2) &= \frac{u_v(x,Q^2)+d_v(x,Q^2)}{2} + \frac{u_s(x,Q^2)+d_s(x,Q^2)}{2} \\ &\quad + s_s(x,Q^2) + b_s(x,Q^2) \qquad (2) \\ \bar{q}(x,Q^2) &= \frac{u_s(x,Q^2)+d_s(x,Q^2)}{2} + c_s(x,Q^2) + t_s(x,Q^2), \end{aligned}$$

where the subscripts v and s label valence and sea contributions, and u, d, c, s, t, b denote the distributions for various quark flavors in a *proton*. The W-boson propagator, which has a negligible effect at low energies, modulates the high-energy cross section and has important consequences for the way the cross section is composed.

I was drawn to this problem by the observation[1)] that the W-boson propagator squeezes the significant contributions of the parton distributions toward smaller values of x with increasing energy. There the QCD-induced growth of the small-x parton distribution enhances the high-energy cross section. This stands in contrast to the familiar effect of QCD evolution at laboratory energies, which is to diminish the total cross section as the valence distribution is degraded at high values of Q^2. At that moment, my colleagues and I had developed for our study of supercollider physics[2)] the first all-flavor set of parton distributions appropriate for applications at small x and large Q^2, so I had in my hands everything needed for a modern calculation of the charged-current cross section at ultrahigh energies. In a sequence of works on the problem,[3,4,5,6)] we have tracked the evolving experimental understanding of parton distributions and investigated many facets of ultrahigh-energy neutrino interactions.

Let us recall some of the principal lessons. The left panel of Figure 1 shows the contributions of various parton species to the total charged-current cross section, as a function of energy. Note that the valence contribution, which dominates at laboratory energies, becomes negligible above about 10^{16} eV, whereas strange- and charm-quark contributions become significant. The cross section integrals in the right panel of Figure 1 illustrate the increasing prominence of small values of x as the neutrino energy increases. At $E_\nu = 10^5$ GeV, nearly all of the cross section comes from $x \gtrsim 10^{-3}$, but by $E_\nu = 10^9$ GeV, nearly all of the cross section lies below $x = 10^{-5}$, where we lack direct experimental information. Reno has given a comprehensive review of small-x uncertainties and the possible influence of new phenomena on the total cross section.[7)]

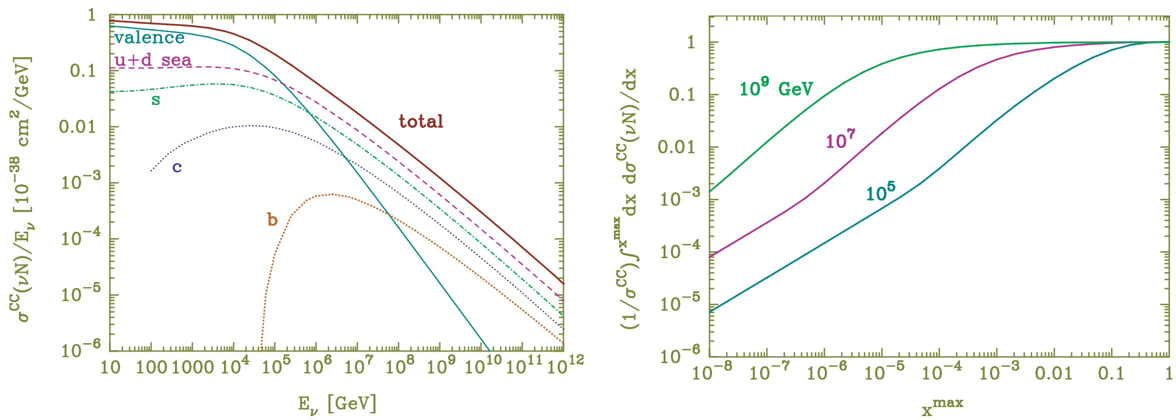

Figure 1: Left panel: Components of the νN charged-current cross section as functions of the neutrino energy for the CTEQ3 distributions. Right panel: Integral cross section $(1/\sigma)\int_0^{x^{\max}} dx\, d\sigma/dx$ for the charged-current reaction $\nu_\mu N \to \mu^- +$ anything at $E_\nu = 10^5, 10^7$, and 10^9 GeV. As the neutrino energy increases, the dominant contributions come from smaller values of x.[5)].

Figure 2 compares our first calculation, using the 1984 EHLQ structure functions, with the recent CTEQ6 parton distributions.[8)] At the highest energies plotted, the cross section is about 1.8× our original estimates, because today's parton distributions rise more steeply at small x than did those of two decades ago. HERA measurements have provided the decisive new information.[9)] At 10^{12} GeV, the QCD enhancement of the small-x parton density has increased the cross section sixty-fold over the parton-model prediction without evolution. HERA measurements of the charged-current reaction $ep \to \nu +$ anything at an equivalent lab energy near 40 TeV observe the damping due to the W-boson propagator and agree with standard-model cross sections.[10)]

The rising cross sections have important implications for neutrino telescopes. The left panel of Figure 3 shows that the Earth is opaque to neutrinos with energies above 40 TeV. This means that the strategy of looking down to distinguish charged-current interactions from the rain of cosmic-ray muons needs to be modified at high energies. On the other hand, the Universe at large is exceptionally poor in nucleons, and so the (νN) interaction length of ultrahigh-energy neutrinos in the cosmos is effectively infinite. The right panel of Figure 3 shows the interaction cross section for neutrinos on electrons in the Earth, which is generally several orders of magnitude longer than the νN interaction length. An important exception is the $\bar{\nu}_e e \to W^-$ resonance at $E_\nu \approx 6 \times 10^{15}$ eV.

2. New Physics in Neutrino-Nucleon Interactions?

New physics typically contributes equally to charged-current and neutral-current cross sections, whereas standard electroweak interactions favor the charged-current

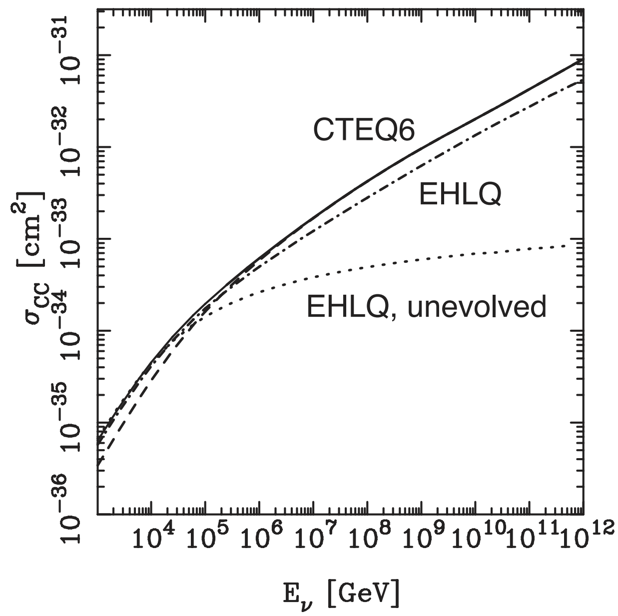

Figure 2: The solid curve shows the charged-current νN cross section calculated using the CTEQ6 parton distributions;[8] the dash-dotted line shows the situation in 1986, using Set 2 of the EHLQ parton distributions.[2] The dotted curve shows the energy dependence of the cross section without QCD evolution, i.e., with the EHLQ distributions frozen at $Q^2 = 5$ GeV^2.[7]

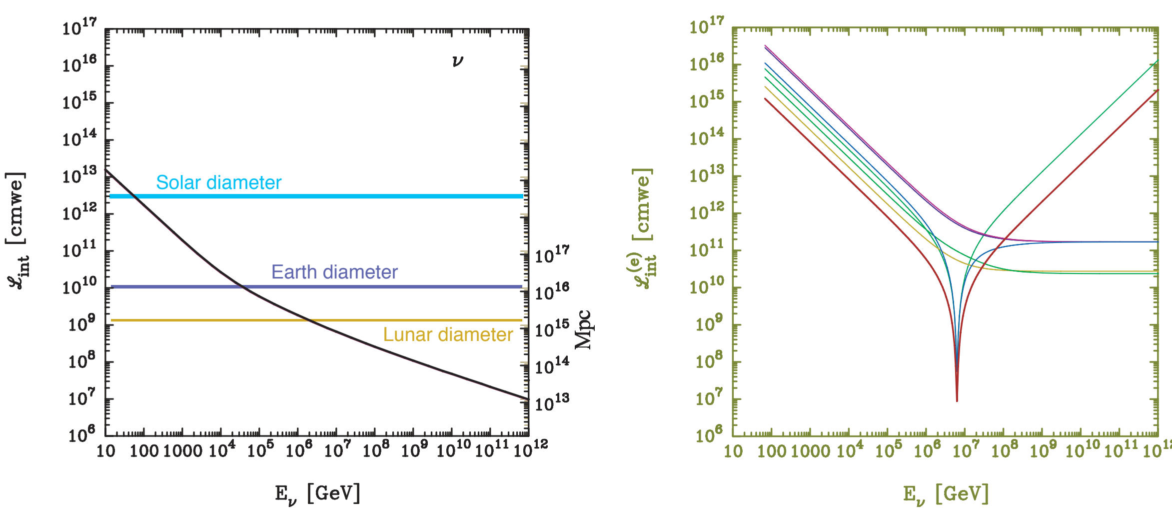

Figure 3: Left panel: Interaction length $\mathcal{L}_{\text{int}}^{\nu N} = 1/\sigma_{\nu N}(E_\nu)N_{\text{A}}$, where N_{A} is Avogadro's number, for the reactions $\nu N \to$ anything as a function of the incident neutrino energy. The left-hand scale, in cmwe, is appropriate for terrestrial applications; the right-hand scale, in Mpc for the current Universe, is appropriate for transport over cosmological distances.[6,11] Right panel: Interaction lengths for neutrino interactions on electron targets. At low energies, from smallest to largest interaction length, the processes are (i) $\bar{\nu}_e e \to$ hadrons, (ii) $\nu_\mu e \to \mu\nu_e$, (iii) $\nu_e e \to \nu_e e$, (iv) $\bar{\nu}_e e \to \bar{\nu}_\mu \mu$, (v) $\bar{\nu}_e e \to \bar{\nu}_e e$, (vi) $\nu_\mu e \to \nu_\mu e$, (vii) $\bar{\nu}_\mu e \to \bar{\nu}_\mu e$.[5]

over neutral current, by a factor of two or three. A step or bump in the neutral-current to charged-current ratio, measured as a function of energy, is thus an excellent diagnostic for the onset of new phenomena. The example of squark production through R-parity–violating interactions that we studied some time ago[12)] offers a specific illustration of this general rule.

3. Influence of Neutrino Oscillations

In the early days of planning for neutrino telescopes, people noticed that observing τ production through the double-bang signature might provide evidence for neutrino oscillations, since—to good approximation—no ν_τ are produced in conventional sources of ultrahigh-energy neutrinos. The discovery of neutrino oscillations is of course already made; the phenomenon of neutrino oscillations means that the flavor mixture at Earth, $\Phi = \{\varphi_e, \varphi_\mu, \varphi_\tau\}$, will be different from the source mixture $\Phi^0 = \{\varphi_e^0, \varphi_\mu^0, \varphi_\tau^0\}$. The essential fact is that the vacuum oscillation length is very short, in cosmic terms. For $|\Delta m^2| = 10^{-5}\ \mathrm{eV}^2$, the oscillation length

$$\mathcal{L}_{\mathrm{osc}} \;=\; 4\pi E_\nu/|\Delta m^2| \approx 2.5 \times 10^{-24}\ \mathrm{Mpc} \cdot (E_\nu/1\ \mathrm{eV}) \tag{3}$$

is a fraction of a megaparsec, even for $E_\nu = 10^{20}$ eV. Accordingly, neutrinos oscillate many times between cosmic source and terrestrial detector.

Neutrinos in the flavor basis $|\nu_\alpha\rangle$ are connected to the mass eigenstates $|\nu_i\rangle$ by a unitary mixing matrix, $|\nu_\alpha\rangle = \sum_i U_{\beta i}|\nu_i\rangle$. It is convenient to idealize $\sin\theta_{13} = 0$, $\sin 2\theta_{23} = 1$, and consider

$$U_{\mathrm{ideal}} = \begin{pmatrix} c_{12} & s_{12} & 0 \\ -s_{12}/\sqrt{2} & c_{12}/\sqrt{2} & 1/\sqrt{2} \\ s_{12}/\sqrt{2} & -c_{12}/\sqrt{2} & 1/\sqrt{2} \end{pmatrix} . \tag{4}$$

Then the transfer matrix $\mathcal{X}$ that maps the source flux Φ^0 into the flux at Earth Φ takes the form

$$\mathcal{X}_{\mathrm{ideal}} = \begin{pmatrix} 1-2x & x & x \\ x & \frac{1}{2}(1-x) & \frac{1}{2}(1-x) \\ x & \frac{1}{2}(1-x) & \frac{1}{2}(1-x) \end{pmatrix} , \tag{5}$$

where $x = \sin^2\theta_{12}\cos^2\theta_{12}$. Because the second and third rows are identical, the ν_μ and ν_τ fluxes that result from any source mixture Φ^0 are equal: $\varphi_\mu = \varphi_\tau$. Independent of the value of x, $\mathcal{X}_{\mathrm{ideal}}$ maps $\Phi^0_{\mathrm{std}} \to \{\frac{1}{3}, \frac{1}{3}, \frac{1}{3}\}$.[b]

The variation of φ_e with the ν_e source fraction φ_e^0 is shown as a sequence of small black squares (for $\varphi_e^0 = 0, 0.1, \ldots, 1$) in Figure 4 for the value $x = 0.21$, which corresponds to $\theta_{12} = 0.57$, the central value in a recent global analysis.[14)] The ν_e fraction at Earth ranges from 0.21, for $\varphi_e^0 = 0$, to 0.59, for $\varphi_e^0 = 1$.

[b] I owe this formulation to Stephen Parke.

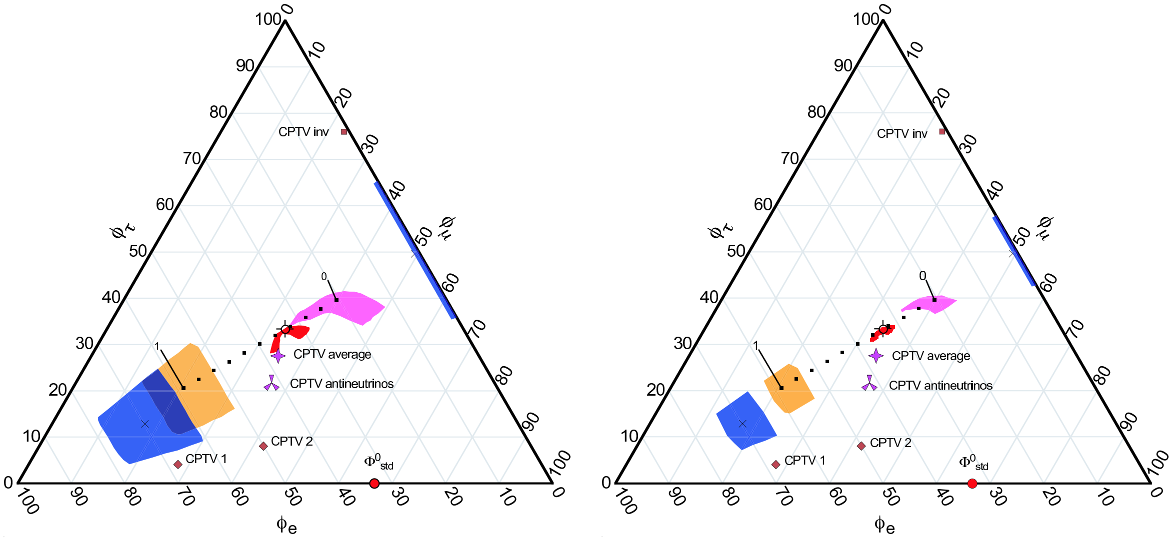

Figure 4: Ternary plots of the neutrino flux Φ at Earth, showing the implications of current (left pane) and future (right pane) knowledge of neutrino mixing. The small black squares indicate the ν_e fractions produced by the idealized transfer matrix $\mathcal{X}_{\text{ideal}}$ as φ_e^0 varies from 0 to 1 in steps of 0.1. A crossed circle marks the standard mixed spectrum at Earth, $\Phi_{\text{std}} = \{\frac{1}{3}, \frac{1}{3}, \frac{1}{3}\}$; for comparison, a red dot marks the standard source spectrum, $\Phi^0_{\text{std}} = \{\frac{1}{3}, \frac{2}{3}, 0\}$. Colored swaths delimit the fluxes at Earth produced by neutrino oscillations from the source mixtures $\Phi^0_0 = \{0, 1, 0\}$ (pink), Φ^0_{std} (red), and $\Phi^0_1 = \{1, 0, 0\}$ (orange), using 95% CL ranges for the oscillation parameters. Black crosses ($\times$) show the mixtures at Earth that follow from neutrino decay, assuming normal ($\varphi_e \approx 0.7$) and inverted ($\varphi_e \approx 0$) mass hierarchies. The blue bands show how current and future uncertainties blur the predictions for neutrino decays. The violet tripod indicates how CPT-violating oscillations shape the mix of antineutrinos that originate in a standard source mixture, and the violet cross averages that $\bar{\nu}$ mixture with the standard neutrino mixture. The brown squares denote consequences of CPT violation for antineutrino decays.[13]

The simple analysis based on $\mathcal{X}_{\text{ideal}}$ is useful for orientation, but it is important to explore the range of expectations implied by global fits to neutrino-mixing parameters. We take[13] $0.49 < \theta_{12} < 0.67$, $\frac{\pi}{4} \times 0.8 < \theta_{23} < \frac{\pi}{4} \times 1.2$, and $0 < \theta_{13} < 0.1$. With current uncertainties in the oscillation parameters, a standard source spectrum, $\Phi^0_{\text{std}} = \{\frac{1}{3}, \frac{2}{3}, 0\}$, is mapped by oscillations onto the red boomerang near $\Phi_{\text{std}} = \{\frac{1}{3}, \frac{1}{3}, \frac{1}{3}\}$ in the left pane of Figure 4. Given that $\mathcal{X}_{\text{ideal}}$ maps $\Phi^0_{\text{std}} \to \Phi_{\text{std}}$ for any value of θ_{12}, it does not come as a great surprise that the target region is of limited extent. The variation of θ_{23} away from $\frac{\pi}{4}$ breaks the identity $\varphi_\mu \equiv \varphi_\tau$ of the idealized analysis.

One goal of neutrino observatories will be to characterize cosmic sources by determining the source mix of neutrino flavors. It is therefore of interest to examine the outcome of different assumptions about the source. We show in the left pane of Figure 4 the mixtures at Earth implied by current knowledge of the oscillation parameters for source fluxes $\Phi^0_0 = \{0, 1, 0\}$ (the purple band near $\varphi_e \approx 0.2$) and $\Phi^0_1 = \{1, 0, 0\}$ (the orange band near $\varphi_e \approx 0.6$). For the Φ^0_{std} and Φ^0_1 source spectra, the uncertainty in θ_{12} is reflected mainly in the variation of φ_e, whereas the uncertainty in θ_{23} is expressed in the variation of φ_μ/φ_τ For the Φ_0^0 case, the influence of the two angles is not so orthogonal. For all the source spectra we consider, the uncertainty in θ_{13} has little effect on the flux at Earth. The extent of the three regions, and the absence of a clean separation between the regions reached from $\Phi_{\rm std}^0$ and Φ_0^0 indicates that characterizing the source flux will be challenging, in view of the current uncertainties of the oscillation parameters.

Over the next five years—roughly the time scale on which large-volume neutrino telescopes will come into operation—we can anticipate improved information on θ_{12} and θ_{23} from KamLAND and the long-baseline accelerator experiments at Soudan and Gran Sasso. We base our projections for the future on the ranges $0.54 < \theta_{12} < 0.63$ and $\frac{\pi}{4} \times 0.9 < \theta_{23} < \frac{\pi}{4} \times 1.1$, still with $0 < \theta_{13} < 0.1$. The results are shown in the right panel of Figure 4. The (purple) target region for the source flux Φ_0^0 shrinks appreciably and separates from the (red) region populated by $\Phi_{\rm std}^0$, which is now tightly confined around $\Phi_{\rm std}$. The (orange) region mapped from the source flux Φ_1^0 by oscillations shrinks by about a factor of two in the φ_e and $\varphi_\mu - \varphi_\tau$ dimensions.

4. Reconstructing the Neutrino Mixture at the Source[13)]

What can observations of the blend Φ of neutrinos arriving at Earth tell us about the source? Inferring the nature of the processes that generate cosmic neutrinos is more complicated than it would be if neutrinos did not oscillate. Because ν_μ and ν_τ are fully mixed—and thus enter identically in $\mathcal{X}_{\rm ideal}$—it is not possible fully to characterize Φ^0. We can, however, reconstruct the ν_e fraction at the source as $\varphi_e^0 = (\varphi_e - x)/(1 - 3x)$, where $x = \sin^2\theta_{12}\cos^2\theta_{12}$. The reconstructed source flux φ_e^0 is shown in Figure 5 as a function of the ν_e flux at Earth. The heavy solid line represents the best-fit value for θ_{12}; the light blue lines and thin solid lines indicate the current and future 95% CL bounds on θ_{12}.

A possible strategy for beginning to characterize a source of cosmic neutrinos might proceed by measuring the ν_e/ν_μ ratio and estimating φ_e under the plausible assumption—later to be checked—that $\varphi_\mu = \varphi_\tau$. Very large ($\varphi_e \gtrsim 0.65$) or very small ($\varphi_e \lesssim 0.15$) ν_e fluxes cannot be accommodated in the standard neutrino-oscillation picture. Observing an extreme ν_e fraction would implicate unconventional physics.

Determining the energy dependence of φ_e^0 may also be of astrophysical interest. In a thick source, the highest energy muons may interact and lose energy before they can decay. In the limit of $\varphi_e^0 = 0$, the arriving flux will be $\Phi = \{x, \frac{1}{2}(1-x), \frac{1}{2}(1-x)\} \approx \{\frac{1}{5}, \frac{2}{5}, \frac{2}{5}\}$ (cf. Figure 4). More generally, measured ν_e fractions that depart significantly from the canonical $\varphi_e = \frac{1}{3}$ would suggest nonstandard neutrino sources. An observed flux $\varphi_e = 0.5 \pm 0.1$ points to a source flux $0.47 \lesssim \varphi_e^0 \lesssim 1$, with current uncertainties, whereas $\varphi_e = 0.25 \pm 0.10$ indicates $0 \lesssim \varphi_e^0 \lesssim 0.35$.

Constraining the source flux sufficiently to test the nature of the neutrino production process will require rather precise determinations of the neutrino flux at

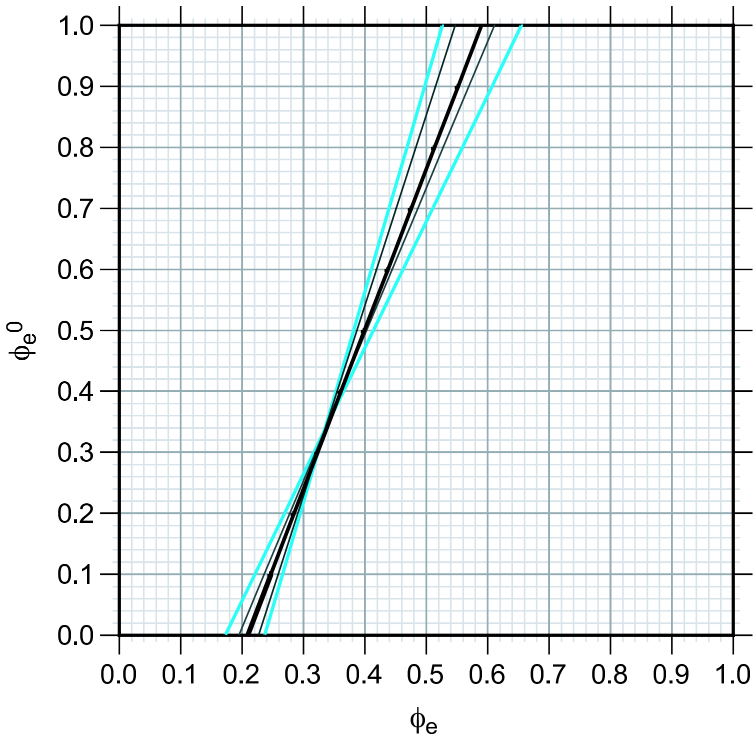

Figure 5: The source flux φ_e^0 of electron neutrinos reconstructed from the ν_e flux φ_e at Earth, using the ideal transfer matrix $\mathcal{X}_{\text{ideal}}$ of Eqn. (5). The heavy solid line refers to $\theta_{12} = 0.57$. The light blue lines refer to the current experimental constraints (at 95% CL), and the thin solid lines refer to a projection of future experimental constraints.[13)]

Earth. Suppose we want to test the idea that the source flux has the standard composition Φ^0_{std}. With today's uncertainty on θ_{12}, a 30% measurement that locates $\varphi_e = 0.33 \pm 0.10$ implies only that $0 \lesssim \varphi_e^0 \lesssim 0.68$. For a measured flux in the neighborhood of $\frac{1}{3}$, the uncertainty in the solar mixing angle is of little consequence: the constraint that arises if we assume the central value of θ_{12} is not markedly better: $0.06 \lesssim \varphi_e^0 \lesssim 0.59$. A 10% measurement of the ν_e fraction, $\varphi_e = 0.33 \pm 0.033$, would make possible a rather restrictive constraint on the nature of the source. The central value for θ_{12} leads to $0.26 \lesssim \varphi_e^0 \lesssim 0.43$, blurred to $0.22 \lesssim \varphi_e^0 \lesssim 0.45$ with current uncertainties.

5. Influence of Neutrino Decays

Beacom and Bell [15)] have shown that observations of solar neutrinos set the most stringent plausible lower bound on the reduced lifetime of a neutrino of mass m as $\tau/m \gtrsim 10^{-4}$ s/eV. This rather modest limit opens the possibility that some neutrinos do not survive the journey from astrophysical sources. Decays of unstable neutrinos over cosmic distances can lead to mixtures at Earth that are incompatible with the oscillations of stable neutrinos.[15,16,13] The candidate decays are transitions between mass eigenstates by emission of a very light particle, $\nu_i \to (\nu_j, \bar{\nu}_j) + X$. Dramatic effects occur when the decaying neutrinos disappear, either by decay to invisible products or by decay into active neutrino species so degraded in energy that they contribute negligibly to the total flux at the lower energy.

If the lifetimes of the unstable mass eigenstates are short compared with the flight time from source to Earth, all the unstable neutrinos will decay, and the (unnormalized) flavor ν_α flux at Earth will be $\widetilde{\varphi}_\alpha(E_\nu) = \sum_{i=\text{stable}} \sum_\beta \varphi^0_\beta(E_\nu)|U_{\beta i}|^2|U_{\alpha i}|^2$, with $\varphi_\alpha = \widetilde{\varphi}_\alpha / \sum_\beta \widetilde{\varphi}_\beta$. Should only the lightest neutrino survive, the flavor mix of neutrinos arriving at Earth is determined by the flavor composition of the lightest mass eigenstate, *independent of the flavor mix at the source.*

For a normal mass hierarchy $m_1 < m_2 < m_3$, the ν_α flux at Earth is $\varphi_\alpha = |U_{\alpha 1}|^2$. Consequently, the neutrino flux at Earth is $\Phi_{\text{normal}} = \{|U_{e1}|^2, |U_{\mu 1}|^2, |U_{\tau 1}|^2\} \approx \{0.70, 0.17, 0.13\}$ for our chosen central values of the mixing angles. If the mass hierarchy is inverted, $m_1 > m_2 > m_3$, the lightest (hence, stable) neutrino is ν_3, so the flavor mix at Earth is determined by $\varphi_\alpha = |U_{\alpha 3}|^2$. In this case, the neutrino flux at Earth is $\Phi_{\text{inverted}} = \{|U_{e3}|^2, |U_{\mu 3}|^2, |U_{\tau 3}|^2\} \approx \{0, 0.5, 0.5\}$. Both Φ_{normal} and Φ_{inverted}, which are indicated by crosses ($\times$) in Figure 4, are very different from the standard flux $\Phi_{\text{std}} = \{\varphi_e = \frac{1}{3}, \varphi_\mu = \frac{1}{3}, \varphi_\tau = \frac{1}{3}\}$ produced by the ideal transfer matrix from a standard source. Observing either mixture would represent a departure from conventional expectations.

The fluxes that result from neutrino decays *en route* from the sources to Earth are subject to uncertainties in the neutrino-mixing matrix. The expectations for the two decay scenarios are indicated by the blue regions in Figure 4, where we indicate the consequences of 95% CL ranges of the mixing parameters now and in the future. With current uncertainties, the normal hierarchy populates $0.61 \lesssim \varphi_e \lesssim 0.77$, and allows considerable departures from $\varphi_\mu = \varphi_\tau$. The normal-hierarchy decay region based on current knowledge overlaps the flavor mixtures that oscillations produce in a pure-ν_e source, shown in orange. (It is, however, far removed from the standard region that encompasses Φ_{std}.) With the projected tighter constraints on the mixing angles, the range in φ_e swept out by oscillation from a pure-ν_e source or decay from a normal hierarchy shrinks by about a factor of two. Neutrino decay then populates $0.65 \lesssim \varphi_e \lesssim 0.74$, and is separated from the oscillations. The degree of separation between the region populated by normal-hierarchy decay and the one populated by mixing from a pure-ν_e source depends on the value of the solar mixing angle θ_{12}. For the seemingly unlikely value $\theta_{12} = \frac{\pi}{4}$, both mechanisms yield $\Phi = \{\frac{1}{2}, \frac{1}{4}, \frac{1}{4}\}$.

The mixtures that result from the decay of the heavier members of an inverted hierarchy entail $\varphi_e \approx 0$. These mixtures are well separated from any that would result from neutrino oscillations, for any conceivable source at cosmic distances.

The energies of neutrinos that may be detected in the future from AGNs and other cosmic sources range over several orders of magnitude, whereas the distances to

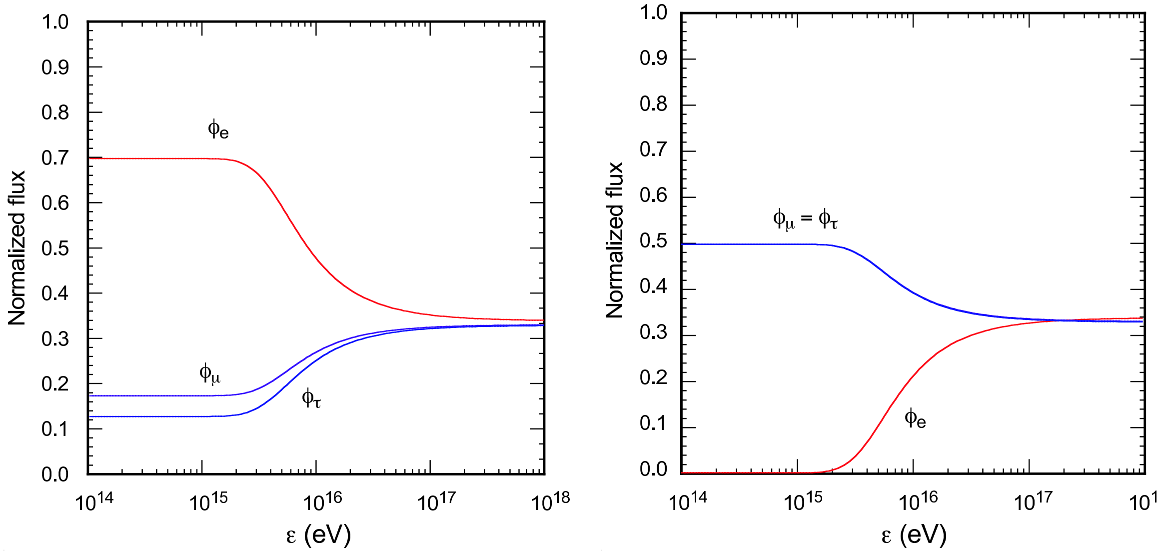

Figure 6: Energy dependence of normalized ν_e, ν_μ, and ν_τ fluxes, for the two-body decay of the two upper mass eigenstates, with the neutrino source at $L = 100$ Mpc from Earth and $\tau/m = 1$ s/eV. The left panel shows the result for a normal mass hierarchy; the right panel shows the result for an inverted mass hierarchy. With suitable rescaling of the neutrino energy $[E_\nu = \varepsilon(1\text{ s/eV})/(\tau_\nu/m_\nu) \cdot L/(100\text{ Mpc})]$, these plots apply for any combination of path length and reduced lifetime.[13)]

such sources vary over perhaps one order of magnitude. The neutrino energy sets the neutrino lifetime in the laboratory frame; more energetic neutrinos survive over longer flight paths than their lower-energy companions.[c] Under propitious circumstances of reduced lifetime, path length, and neutrino energy, it might be possible to observe the transition from more energetic survivor neutrinos to less energetic decayed neutrinos.

If decay is not complete, the (unnormalized) flavor ν_α flux arriving at Earth from a source at distance L is given by $\widetilde{\varphi}_\alpha(E_\nu) = \sum_i \sum_\beta \varphi^0_\beta(E_\nu)|U_{\beta i}|^2|U_{\alpha i}|^2 e^{-(L/E_\nu)(m_i/\tau_i)}$, with the normalized flux $\varphi_\alpha(E_\nu) = \widetilde{\varphi}_\alpha(E_\nu)/\sum_\beta \widetilde{\varphi}_\beta(E_\nu)$. An idealized case will illustrate the possibilities for observing the onset of neutrino decay and estimating the reduced lifetime. Assume a normal mass hierarchy, $m_1 < m_2 < m_3$, and let $\tau_3/m_3 = \tau_2/m_2 \equiv \tau/m$. For a given path length L, the neutrino energy at which the transition occurs from negligible decays to complete decays is determined by τ/m. The left pane of Figure 6 shows the energy evolution of the normalized neutrino fluxes arriving from a standard source; the energy scale is appropriate for the case $\tau/m = 1$ s/eV and $L = 100$ Mpc.

If we locate the transition from survivors to decays at neutrino energy $E^\star$, then we can estimate the reduced lifetime in terms of the distance to the source as

$$\tau/m \approx 100\text{ s/eV} \cdot \left(\frac{L}{1\text{ Mpc}}\right)\left(\frac{1\text{ TeV}}{E^\star}\right) . \qquad (6)$$

[c] A similar phenomenon is familiar for cosmic-ray muons.

In practice, ultrahigh-energy neutrinos are likely to arrive from a multitude of sources at different distances from Earth, so the transition region will be blurred.[d] Nevertheless, it would be rewarding to observe the decay-to-survival transition, and to use Eqn. (6) to estimate—even within one or two orders of magnitude—the reduced lifetime. If no evidence appears for a flavor mix characteristic of neutrino decay, then Eqn. (6) provides a lower bound on the neutrino lifetime. For that purpose, the advantage falls to large values of $L/E^{\star}$, and so to the lowest energies at which neutrinos from distant sources can be observed. Observing the standard flux, $\Phi_{\text{std}} = \{\frac{1}{3}, \frac{1}{3}, \frac{1}{3}\}$, which is incompatible with neutrino decay, would strengthen the current bound on τ/m by some seven orders of magnitude, for 10-TeV neutrinos from sources at 100 Mpc.

6. UHE Neutrino Annihilation on Relic Neutrinos

The neutrino gas that we believe permeates the present Universe has never been detected directly. By observing resonant annihilation of extremely-high-energy cosmic neutrinos on the background neutrinos through the reaction $\nu\bar{\nu} \to Z^0$,[17,18,19,20,21] we could hope to confirm the presence of the relic neutrinos and learn the absolute neutrino masses and the flavor composition of the neutrino mass eigenstates. I have recently made a detailed study of the prospects for cosmic-neutrino annihilation spectroscopy with Gabriela Barenboim and Olga Mena.[11] I summarize some of our main findings here.

I present the components of the neutrino-(anti)neutrino cross sections in Figure 7. The feature that matters is the Z^0-formation line that occurs near the resonant energy $E_\nu^{Z\text{res}} = M_Z^2/2m_{\nu_i}$. Existing knowledge of neutrino oscillations allows us to project the neutrino mass spectrum in terms of the unknown mass of the lightest neutrino. The expectations are shown in Figure 8 for normal and inverted mass hierarchies.

In an idealized *Gedankenexperiment,* we may consider an extremely high-energy neutrino beam traversing a very long column with the relic-neutrino properties of the current Universe. We neglect for now the expansion of the Universe and the thermal motion of the relic neutrinos. The "cosmic neutrino attenuator" is thus a column of length L with uniform neutrino density $n_{\nu 0} = 56\ \text{cm}^{-3}$ of each neutrino species, $\nu_e, \bar{\nu}_e, \nu_\mu, \bar{\nu}_\mu, \nu_\tau, \bar{\nu}_\tau$. If the column of relic neutrinos is thick enough to attenuate neutrinos appreciably through resonant absorption at the Z^0 gauge boson, the energies that display absorption dips point to the neutrino masses through the resonant-energy condition. The relative depletion of ν_e, ν_μ, ν_τ in each of the lines measures the flavor composition of the corresponding neutrino mass eigenstate.

Even if we had at our disposal an adequate neutrino beam (with energies extending beyond 10^{26} eV), the time required to traverse one interaction length for $\nu\bar{\nu} \to Z^0$ annihilation on the relic background in the current Universe (1.2×10^4 Mpc = 39 Gly) exceeds the age of the Universe, not to mention the human attention span. If we are

[d] Our assumption that $\tau_3/m_3 = \tau_2/m_2 \equiv \tau/m$ is also a special case.

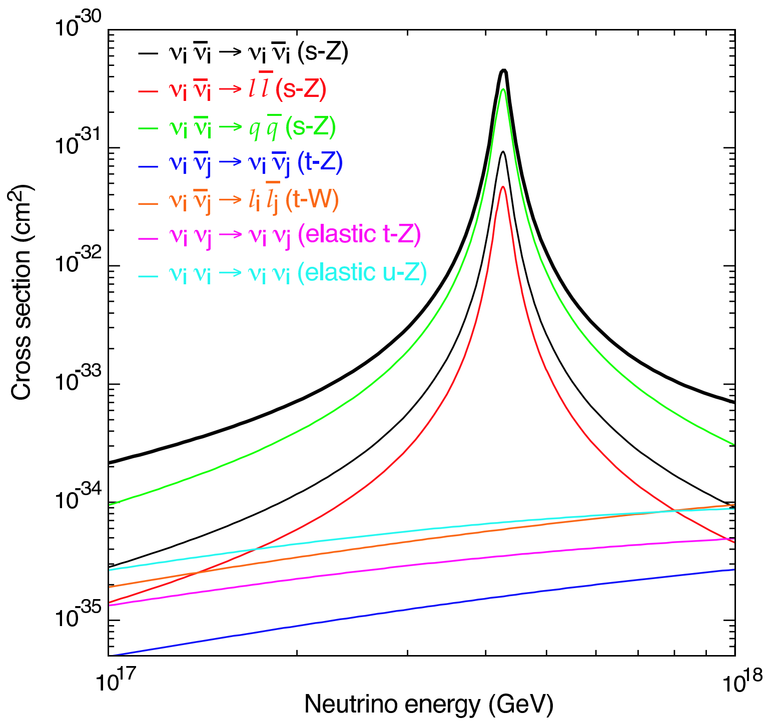

Figure 7: Total neutrino annihilation cross section and the different contributing channels as a function of the neutrino energy assuming a relic neutrino mass of $m_\nu = 10^{-5}$ eV and zero redshift.[11)]

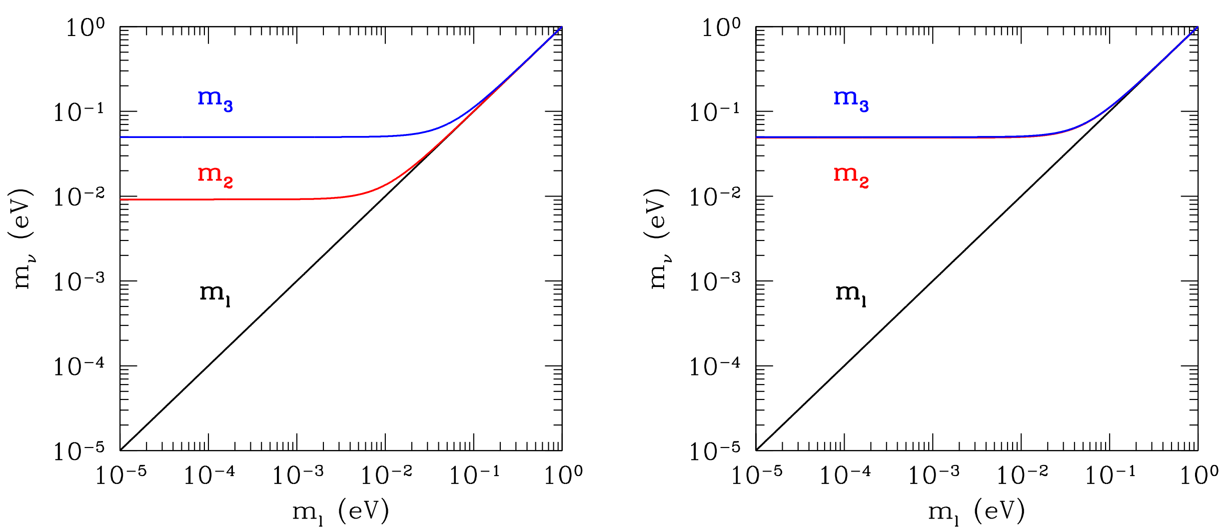

Figure 8: Favored values for the neutrino masses as functions of the lightest neutrino mass, m_ℓ, in the three neutrino scenario for normal hierarchy (left panel, $m_\ell = m_1$) and the inverted hierarchy (right panel, $m_\ell = m_3$).[11)]

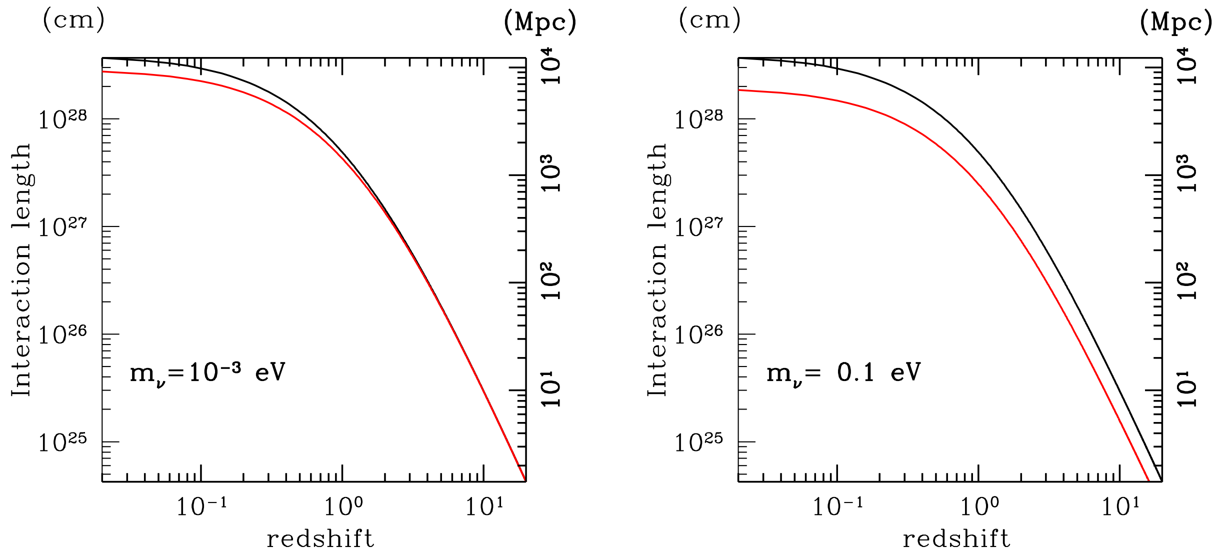

Figure 9: Interaction lengths versus redshift at the Z^0 resonance for neutrino masses $m_\nu = 10^{-3}, 10^{-1}$ eV (left, and right panels). The left-hand scales are in centimeters, the right-hand scales in megaparsecs. In the center and right panels, the upper (black) line is for the Dirac-neutrino case; the lower (red) line applies to Majorana neutrinos.[11)]

ever to detect the attenuation of neutrinos on the relic-neutrino background, we shall have to make use of astrophysical or cosmological neutrinos sources traversing the Universe over cosmic time scales. The expansion of the Universe over the propagation time of the neutrinos entails three important effects: the evolution of relic-neutrino density, the redshift of the incident neutrino energy, and the redshift of the relic-neutrino temperature.

The decrease of interaction lengths with increasing redshift shown in Figure 9 reveals that for redshifts in the range from one to ten, the interaction length matches the distance to the AGNs we consider as plausible UHE neutrino sources.[e] The absorption lines that result from a full calculation, including the effects of the relics' Fermi motion and the evolution of the Universe back to redshift $z = 20$, are shown in Figure 10 for two values of the lightest neutrino mass, $m_\ell = 10^{-5}$ and 10^{-3} eV. Although the lines are distorted and displaced from their natural shapes and positions by redshifting and Fermi motion, they would nevertheless confirm our expectations for the relic neutrino background and give important information about the neutrino spectrum. In particular, the ν_e/ν_μ ratio, shown in Figure 11, is a marker for the normal or inverted mass hierarchy.

The observation of cosmic-neutrino absorption lines will open the way—at least in principle—to new insights about neutrino properties and the thermal history of the

[e]. . . though not perhaps with the energies required here!

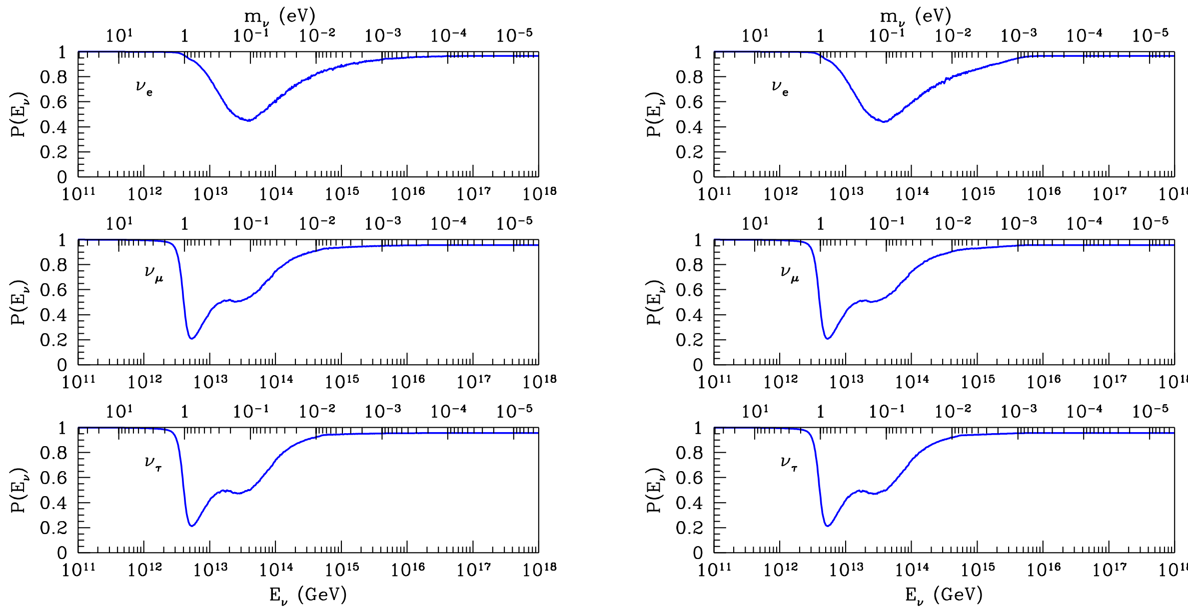

Figure 10: Survival probabilities for ν_e, ν_μ, and ν_τ as a function of the neutrino energy, after integration back to redshift $z = 20$, taking into account the Fermi smearing induced by the thermal motion of the relic neutrinos. The results apply for a normal hierarchy with lightest neutrino mass $m_\ell = 10^{-5}$ eV (left panel) or $m_\ell = 10^{-3}$ eV (right panel). [11]

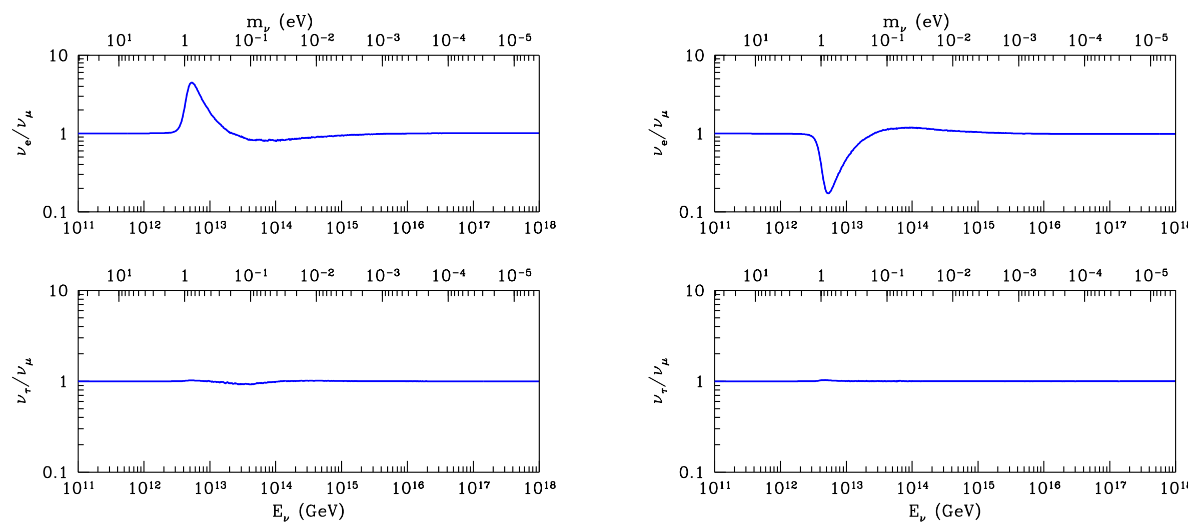

Figure 11: Flux ratios ν_e/ν_μ and ν_τ/ν_μ at Earth, for normal (left panel) and inverted (right panel) mass hierarchies with $m_\ell = 10^{-5}$ eV, after integration back to redshift $z = 20$ and a thermal averaging over the relic-neutrino momentum distribution. The scale at the top shows the neutrino mass defined as $m_\nu = M_Z^2/2E_\nu$ that would be inferred if neutrino energies were not redshifted.[11]

universe. Our calculations, with their successive inclusion of potentially significant effects, show that how the tale unfolds will depend on factors we cannot foresee. The earlier in redshift the relevant cosmic-neutrino sources appear, the lower the present-day energy of the absorption lines and the denser the column of relics the super-high-energy neutrinos must traverse. In particular, the appearance of dips at energies much lower that we expect points to early—presumably nonacceleration—sources, that could give us insight into fundamental physics at early times and high energy scales. On the other hand, integration over a longer range in redshift means more smearing and distortion of the absorption lines.

If it can be achieved at all, the detection of neutrino absorption lines will not be done very soon, and the interpretation is likely to require many waves of observation and analysis. Nevertheless, observations of cosmic-neutrino absorption lines offer the possibility to establish the existence of another relic from the big bang and, conceivably, they may open a window on periods of the thermal history of the universe not readily accessible by other means.

7. Neutrino Coannihilation on Dark-Matter Relics?

If the nonbaryonic dark matter that makes up some 20% of the mass-energy of the Universe is composed of particle relics, might it be possible to observe evidence of neutrino–dark-matter coannihilations? Clearly the answer depends on the nature of the dark-matter particles. An instructive example is provided by neutralino dark matter in the framework of supersymmetric models, so I have been analyzing this case with Gabriela Barenboim and Olga Mena.[22)]

If the lightest supersymmetrical particle is the neutralino χ_1^0, with a mass in the neighborhood of 150 GeV, then we immediately have two pieces of good news. First, the resonant energy for sneutrino formation in the reaction $\nu\chi_1^0 \to \tilde{\nu}$ is typically $E_\nu^{\mathrm{res}} \approx 400$ GeV, so we are assured of a reasonable neutrino flux, perhaps even from cosmic-ray interactions in the atmosphere.[f] Second, the $\tilde{\nu}$-formation cross section is typically about 10% of the $\nu\bar{\nu} \to Z$ annihilation cross section, so is not small.

There ends the good news, at least for the Universe at large. Whereas we expect that stable neutrinos should be the most abundant particles in the Universe after the photons of the cosmic microwave background, the neutralino gas is on average very tenuous. The neutralino fraction of the Universe (identified with the dark-matter fraction) is $\Omega_\chi h^2 = \varrho_\chi h^2/\varrho_c$, where the numerical value of the critical density is $\varrho_c = 1.05 \times 10^{-5} h^2$ GeV cm^{-3}. Consequently the mass density of relic neutralinos is $\varrho_\chi = 1.05 \times 10^{-5} \cdot \Omega_\chi h^2 = \bar{\mathcal{N}}_\chi M(\chi)$, where $\bar{\mathcal{N}}_\chi$ is the mean number density of relic neutralinos throughout the Universe. For neutralino masses in the 150-GeV range, $\bar{\mathcal{N}}_\chi \lesssim 10^{-8}$ cm^{-3}, some ten orders of magnitude smaller than the current relic neutrino

[f]The Rome group has studied the Lorentz-boosted case of ultrahigh-energy neutralinos incident on the relic neutrinos.[23)]

density and 31 orders of magnitude smaller than the density of electrons in water. As a result, the interaction length for resonant sneutrino formation is some 10^{15} Mpc, so $\nu\chi \rightarrow \tilde{\nu}$ coannihilation is utterly irrelevant in the Universe at large.

Our location in the Universe may not be privileged, but it is not average. Dark matter clusters in galaxies. A useful benchmark is the spherically symmetric universal profile proposed by Navarro, Frenk, and White,[24] which yields a mean dark-matter density (neglecting local influences) of $\varrho_\chi \approx 0.3$ GeV cm^{-3} at our distance from the galactic center, for a number density $\mathcal{N}_\chi =$ a few $\times\ 10^{-3}$ cm^{-3}, five orders of magnitude enhancement over the relic density in the Universe at large. Even considering possible enhancements of the relic density in the solar system, the rate of interactions produced by atmospheric neutrinos in Earth's atmosphere is negligible.

There is one remaining hope, not for neutrino observatories but for gamma-ray telescopes. The galaxy as a whole contains some 10^{65} neutralinos in the scenario we are describing. With a conservative neutrino flux, we might expect 10^{24} sneutrino events in a year. Some fraction of these will decay inelastically and give rise to a γ spectrum in the few-GeV range. Alas, the number of coannihilations viewed by a detector near Earth is only $\approx 10^{-21}$ cm^{-2} y^{-1}.

8. Gravitational Lensing of Neutrinos

Surely neutrinos—in common with other forms of matter and energy—experience gravitational interactions. Where is the observational evidence to support this assertion? No analogue of the classic demonstration of the deflection of starlight by the Sun is in prospect. We do not know any continuous intense point source of neutrinos, and the angular resolution of neutrino telescopes—$\approx 5^\circ$ for Super-Kamiokande in the few-MeV range and the 0.5° goal for km^3-scale ultrahigh-energy neutrino telescopes—is poorly matched to the anticipated 1.75-arcsecond deflection of neutrinos from a distant source. We must therefore look elsewhere to demonstrate that neutrinos have normal gravitational interactions..

The Supernova 1987A neutrino burst, recorded within 3 h of the associated optical display after a 166 000-year voyage, provides circumstantial evidence that neutrinos and photons follow the same trajectories in the gravitational field of our galaxy, to an accuracy better than 0.5%.[25,26] On the assumption that neutrinos do have normal gravitational interactions, weak lensing induced by relic neutrinos could suppress the large-scale structure power spectrum on small scales.[27,28] The SN1987A argument, though telling, is indirect, and weak lensing by relic neutrinos has not yet been observed. Can we imagine a more direct manifestation of gravity's influence on neutrinos?

Raghavan has recently advocated[29] a neutrino analogue of the Pound–Rebka experiment[30] to demonstrate the blue shift of neutrinos falling in a gravitational field, applying the Mössbauer effect to recoilless resonant capture of antineutrinos.

At this meeting, Minakata[31)] has suggested that the method might lead to a tabletop measurement of the neutrino parameters Δm^2_{31} and θ_{13}.[32)]

With my colleagues Olga Mena and Irina Mocioiu, I have been looking into the possibility of observing the lensing of supernova neutrinos by the black hole at the galactic center.[33)] The improbably ideal case of a supernova explosion on the other side of the galaxy, symmetric to our position, would be characterized by a prodigious amplification of the neutrino flux at Earth. The signals from lensed supernovae that are not quite so impeccably positioned would be characterized by noticeable amplification and by dispersion in arrival time that would stretch the apparent duration of the neutrino burst to several times its natural length. We might be richly rewarded for attending to rare events, for

> "It is a part of probability that many improbable things will happen."
> — George Eliot (after Aristotle's *Poetics*), *Daniel Deronda*

9. From Neutrino Astronomy to Particle Physics

As we move into the era of km^3-scale neutrino telescopes,[34,35)] the baseline strategy of detecting muons produced by charged-current interactions in or near the instrumented volume seems certain to achieve the first, astronomical, goal of this enterprise: prospecting for distant neutrino sources. In this talk, I have tried to show by example how more ambitious detection strategies might be rewarded with sensitivity to information of interest to particle physics. For the determination of neutrino properties, flavor tagging would be immensely valuable, and we can imagine obtaining information about neutrino sources and neutrino decays that it not otherwise available. The ability to record and to characterize the energy of neutral-current events may give decisive sensitivity to the onset of new phenomena.

I have also exhibited two reminders of the importance of being attentive to surprises. If neutrino sources are associated with star-bearing regions, and so turn on at redshifts smaller than 20, then cosmic-neutrino absorption spectroscopy is probably a very distant dream. But if non-acceleration sources, such as the decay of cosmic strings at very early times, are important neutrino producers, then we might find absorption lines at considerably lower energies than the rule of thumb $E_\nu^{Z\text{res}} = M_Z^2/2m_\nu$ would suggest. Gravitational lensing of neutrinos might lead to characteristic signatures of anomalously intense supernova bursts of anomalously long duration. Finally, neutrinos interacting with dark-matter relics might produce signals for large-area γ-ray telescopes.

10. Acknowledgements

I salute with pleasure the hard work and perceptive insights of my extraterrestrial-neutrino collaborators, Hallsie Reno, Terry Walker, Raj Gandhi, Ina Sarcevic, Marcela Carena, Magda Lola, Debajyoti Choudhury, Vu Anh Tuan, Gabriela Barenboim, Olga Mena, and Irina Mocioiu. Fermilab is operated by Universities Research Association Inc. under Contract No. DE-AC02-76CH03000 with the U.S. Department of Energy. I am grateful for the hospitality of the CERN Particle Theory Group during the preparation of this talk. Finally, I thank Milla Baldo-Ceolin for animating this splendid week in Venice.

11. References

1) Y. M. Andreev, V. S. Berezinsky and A. Y. Smirnov, *Phys. Lett. B*$\mathbf{84}$ (1979) 247.
2) E. Eichten, I. Hinchliffe, K. D. Lane and C. Quigg, *Rev. Mod. Phys.* **56** (1984) 579 [Addendum-*ibid.* **58** (1986) 1065].
3) C. Quigg, M. H. Reno and T. P. Walker, *Phys. Rev. Lett.* **57** (1986) 774.
4) M. H. Reno and C. Quigg, *Phys. Rev. D*$\mathbf{37}$ (1988) 657.
5) R. Gandhi, C. Quigg, M. H. Reno and I. Sarcevic, *Astropart. Phys.* **5** (1996) 81 [arXiv:hep-ph/9512364].
6) R. Gandhi, C. Quigg, M. H. Reno and I. Sarcevic, *Phys. Rev. D*$\mathbf{58}$ (1998) 093009 [arXiv:hep-ph/9807264].
7) M. H. Reno, *Nucl. Phys. Proc. Suppl.* **143** (2005) 407 [arXiv:hep-ph/0410109].
8) J. Pumplin, D. R. Stump, J. Huston, H. L. Lai, P. Nadolsky and W. K. Tung, *JHEP* **0207** (2002) 012 [arXiv:hep-ph/0201195].
9) C. Adloff *et al.* [H1 Collaboration], *Eur. Phys. J. C* **30** (2003) 1 [arXiv:hep-ex/0304003]; *ibid.* **21** (2001) 33 [arXiv:hep-ex/0012053]. S. Chekanov *et al.* [ZEUS Collaboration], *Eur. Phys. J. C* **21** (2001) 443 [arXiv:hep-ex/0105090].
10) T. Ahmed *et al.* [H1 Collaboration], *Phys. Lett. B*$\mathbf{324}$ (1994) 241; M. Derrick *et al.* [ZEUS Collaboration], *Phys. Rev. Lett.* **75** (1995) 1006 [arXiv:hep-ex/9503016].
11) G. Barenboim, O. Mena Requejo and C. Quigg, *Phys. Rev. D*$\mathbf{71}$ (2005) 083002 [arXiv:hep-ph/0412122].
12) M. Carena, D. Choudhury, S. Lola and C. Quigg, *Phys. Rev. D*$\mathbf{58}$ (1998) 095003 [arXiv:hep-ph/9804380].
13) G. Barenboim and C. Quigg, *Phys. Rev. D*$\mathbf{67}$ (2003) 073024 [arXiv:hep-ph/0301220].
14) P. C. de Holanda and A. Y. Smirnov, JCAP **0302** (2003) 001 [arXiv:hep-ph/0212270].
15) J. F. Beacom and N. F. Bell, *Phys. Rev. D*$\mathbf{65}$ (2002) 113009 [arXiv:hep-ph/0204111].
16) J. F. Beacom, N. F. Bell, D. Hooper, S. Pakvasa and T. J. Weiler, *Phys. Rev. Lett.* **90** (2003) 181301 [arXiv:hep-ph/0211305].
17) T. J. Weiler, *Phys. Rev. Lett.* **49** (1982) 234.

18) E. Roulet, *Phys. Rev. D***47** (1993) 5247.
19) D. Fargion, B. Mele and A. Salis, *Astrophys. J.* **517** (1999) 725 [arXiv:astro-ph/9710029].
20) T. J. Weiler, *Astropart. Phys.* **11** (1999) 303 [arXiv:hep-ph/9710431].
21) B. Eberle, A. Ringwald, L. Song and T. J. Weiler, *Phys. Rev. D***70** (2004) 023007 [arXiv:hep-ph/0401203].
22) G. Barenboim, O. Mena Requejo and C. Quigg, "Neutrino Coannihilation on Dark-Matter Relics?" FERMILAB–PUB–06/050–T.
23) A. Datta, D. Fargion and B. Mele, *JHEP* **0509** (2005) 007 [arXiv:hep-ph/0410176].
24) J. F. Navarro, C. S. Frenk and S. D. M. White, *Astrophys. J.* **490** (1997) 493.
25) M. J. Longo, *Phys. Rev. Lett.* **60** (1988) 173.
26) L. M. Krauss and S. Tremaine, *Phys. Rev. Lett.* **60** (1988) 176.
27) A. R. Cooray, *Astron. Astrophys.* **348** (1999) 31 [arXiv:astro-ph/9904246].
28) J. Lesgourgues, L. Perotto, S. Pastor and M. Piat, arXiv:astro-ph/0511735.
29) R. S. Raghavan, "Recoilless resonant capture of antineutrinos," arXiv:hep-ph/0511191; "Recoilless resonant capture of antineutrinos from tritium decay," arXiv:hep-ph/0601079.
30) R. V. Pound and G. A. Rebka, Jr., *Phys. Rev. Lett.* **4** (1960) 337.
31) H. Minakata, "Do Neutrinos Violate *CP*?," talk presented at this workshop, axpd24.pd.infn.it/NO-VE2006/talks/NOVE_Minakata.ppt.
32) H. Minakata and S. Uchinami, "Recoilless resonant absorption of monochromatic neutrino beam for measuring Δm^2_{31} and θ_{13}," arXiv:hep-ph/0602046.
33) O. Mena Requejo, I. Mocioiu, and C. Quigg, "Gravitational Lensing of Supernova Neutrinos" FERMILAB–PUB–06/051–T.
34) P. O. Hulth, "From AMANDA to IceCube," talk presented at this workshop, axpd24.pd.infn.it/NO-VE2006/talks/NOVE_Hulth.ppt.
35) J. J. Aubert, "Neutrino Telescopes," talk presented at this workshop, axpd24.pd.infn.it/NO-VE2006/talks/NOVE_Aubert.ppt.