NEUTRINO ASTROPHYSICS AT THE CROSS ROADS

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Nonstandard neutrino properties (masses, mixing, sterile states, electromagnetic interactions, and so forth) can have far-reaching ramifications in astrophysics and cosmology. We look at the most interesting cases in the light of the powerful current indications for neutrino oscillations.

1 Introduction

The indications for neutrino oscillations from the atmospheric and solar neutrino anomalies and from the LSND experiment are now so overwhelming that the discourse in neutrino physics has changed. One no longer asks if these particles indeed oscillate, one rather debates the most plausible pattern of masses and mixing angles, and if the existence of a sterile neutrino is required. Of course, all of the indications for oscillations are to various degrees preliminary, yet so intriguing that it is difficult to resist their charm.

It is a truism that astrophysics and cosmology play a unique role in neutrino physics, and conversely, that these light, weakly interacting particles are absolutely crucial for some of the most interesting astrophysical phenomena such as core-collapse supernovae and for the universe at large. Therefore, in this brief survey of current topics in neutrino astrophysics it behoves us to discuss what astrophysics and cosmology contribute to the current debate in neutrino physics and what the future perspectives are.

To this end we begin in Sec. 2 with an overview of the current indications for neutrino oscillations and possible global interpretations. Astrophysical neutrinos, i.e. those from the Sun and from cosmic-ray interactions in the upper atmosphere, play a dominant role in this context. In Sec. 3 we next turn to the cosmological arguments relevant to neutrino physics (dark matter, structure formation, cosmic microwave background, big-bang nucleosynthesis). Supernova (SN) neutrinos are the topic of Sec. 4 where we discuss the role of neutrino masses and oscillations in this environment, the interpretation of the SN 1987A neutrino burst, and what one could learn from a future galactic SN. The recent developments in high-energy neutrino astronomy are touched upon in Sec. 5, while in Sec. 6 astrophysical aspects of neutrino electromagnetic properties are briefly discussed in the light of the current evidence for oscillations. Finally, in Sec. 7 we summarize our conclusions.
2 Evidence for Neutrino Oscillations

2.1 Atmospheric Neutrinos

The current evidence for neutrino oscillations arises from the atmospheric neutrino anomaly, the solar neutrino problem, and the LSND experiment. It is probably fair to say that at present the most convincing indication comes from atmospheric neutrinos. We thus begin our short survey with this spectacular case that has changed the perception of this field.

The Earth is immersed in a diffuse flux of high-energy cosmic rays consisting of protons and nuclei. The upper atmosphere acts as a “beam dump” where these particles quickly lose their energy by the production of secondary pions (and some kaons) which subsequently decay according to the simple scheme

\[
\begin{align*}
\pi^+ &\rightarrow \mu^+ + \nu_\mu, \\
\mu^+ &\rightarrow e^+ + \nu_e + \bar{\nu}_\mu, \\
\pi^- &\rightarrow \mu^- + \bar{\nu}_\mu, \\
\mu^- &\rightarrow e^- + \bar{\nu}_e + \nu_\mu.
\end{align*}
\]

The expected unequal flavor distribution \(\nu_e : \nu_\mu : \nu_\tau \approx 1 : 2 : 0\) allows one to use the atmospheric neutrino flux to search for flavor oscillations. Of course, at energies beyond a few GeV the muons do not all decay before hitting the Earth so that the \(\nu_\mu / \nu_e\) flavor ratio increases with energy. Still, while the absolute neutrino flux predictions have large uncertainties, perhaps on the 20% level, the expected flavor ratio is thought to be nearly model independent and calculable for all relevant energies to within a few percent.

First events from atmospheric neutrinos were measured in two pioneering experiments in the mid-sixties, but it is only since the late eighties that several large underground detectors began to address the question of flavor oscillations in earnest. Around 1988 the Kamiokande water Cherenkov detector revealed a significantly reduced \(\nu_\mu / \nu_e\) flavor ratio—the atmospheric neutrino anomaly. There was no alternative explanation to oscillations, but a “smoking-gun” signature became available only with the high counting rates of SuperKamiokande which has taken data since April 1996.

For a given solid angle, the atmospheric neutrino flux from above should be equal to that produced in the atmosphere of the antipodes because the \(r^{-2}\) flux dilution with distance cancels a corresponding increase in surface area. However, SuperKamiokande observed a pronounced up-down-asymmetry in the multi-GeV sample (visible energy deposition in the detector exceeding 1.33 GeV). Using the zenith-angle range \(-1.0 \leq \cos \theta \leq -0.2\) as defining “up,” and the corresponding range \(0.2 \leq \cos \theta \leq 1.0\) for “down,” the \(\nu_e + \bar{\nu}_e\) flux shows a ratio \(\text{up/down} = 0.93_{-0.13}^{+0.06}\) while \(\nu_\mu + \bar{\nu}_\mu\) has 0.54_{-0.05}. It is
this up-down-asymmetry which gives one confidence that there is no simple explanation in terms of the neutrino production process in the atmosphere or the experimental flavor identification.

Neutrino oscillations, on the other hand, provide a simple and consistent interpretation. In the usual two-flavor formalism with a vacuum mixing angle $\Theta$, the appearance probability for the oscillation from a flavor $\nu$ to $\nu'$ is

$$P(\nu \to \nu') = \sin^2 2\Theta \sin^2 \left( 1.27 \frac{\Delta m^2_{\nu}}{eV^2} \frac{L}{\text{km}} \frac{\text{GeV}}{E_{\nu}} \right).$$

If the $\nu_\mu$'s oscillate into $\nu_\tau$'s with a nearly maximal mixing angle and if $\Delta m^2_{\nu}$ is of order $10^{-3} \text{eV}^2$, one obtains the observed behavior since the relevant energies are a few GeV and $L$ to the other side of the Earth is around $10^4 \text{km}$. The detailed 90% CL contours for the allowed range of mixing parameters from different signatures in Kamiokande and SuperKamiokande are summarized in Fig. 1. Meanwhile, more data have been taken, shifting the curve (1) to somewhat larger $\Delta m^2_{\nu}$ values with the boundaries shown in Table 1.

Equation (2) suggests that one should plot the data according to their $L/E_{\nu}$ as in Fig. 2. This representation provides perhaps the most convincing argument for the reality of atmospheric neutrino oscillations. The flat distribution of the $\nu_e$ points excludes $\nu_\mu \to \nu_e$ oscillations as a dominant channel.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure1.png}
\caption{Allowed mixing parameters at 90% CL from atmospheric neutrinos for $\nu_\mu \to \nu_\tau$ oscillations. They are based on the contained events in SuperKamiokande (1) and Kamiokande (2), the upward through-going muons in SuperKamiokande (3) and Kamiokande (4), and the stopping fraction of upward going muons in SuperKamiokande (5). (Figure reproduced with kind permission of T. Kajita.)}
\end{figure}
Figure 2: $L/E_\nu$ plot for the fully contained events at SuperKamiokande. The points show the ratio of measured counts over Monte Carlo expectation in the absence of oscillations. The dashed lines show the expectation for $\nu_\mu \to \nu_\tau$ oscillations with $\Delta m^2_{\nu} = 2.2 \times 10^{-3} \text{ eV}^2$ and $\sin^2 2\theta = 1$. (Figure reproduced with kind permission of T. Kajita.)

in agreement with the CHOOZ limits on this mode. Therefore, $\nu_\mu \to \nu_\tau$ or oscillations into a sterile channel $\nu_\mu \to \nu_s$ are favored.

A calculation of the $\nu_\mu \to \nu_s$ oscillation probability must include the refractive energy shift in the Earth. Recall that the neutrino weak potential is

$$V_{\text{weak}} = \frac{G_F n_B}{\sqrt{2}} \times \begin{cases} -2Y_n + 4Y_e & \text{for } \nu_e, \\ -2Y_n & \text{for } \nu_{\mu,\tau}, \\ 0 & \text{for } \nu_s, \end{cases} \quad (3)$$

where the upper sign refers to neutrinos, the lower sign to antineutrinos, $G_F$ is the Fermi constant, $n_B$ the baryon density, $Y_n$ the neutron and $Y_e$ the electron number per baryon (both about 1/2 in normal matter). Numerically we have

$$\frac{G_F n_B}{2\sqrt{2}} = 1.9 \times 10^{-14} \text{ eV} \frac{\rho}{\text{g cm}^{-3}}. \quad (4)$$

The dispersion relation is $E_\nu = V_{\text{weak}} + \sqrt{p^2_\nu + m^2_\nu}$ so that $V_{\text{weak}}$ should be compared with $m^2_\nu/2p_\nu$. For $\Delta m^2_\nu$ around $10^{-3} \text{ eV}^2$, $p_\nu$ of a few GeV, and $\rho$ of a few g cm$^{-3}$, the energy difference between $\nu_\mu$ and $\nu_s$ arising from $V_{\text{weak}}$
is about the same as that from $\Delta m^2_{2}/2p_{\nu}$. The resulting modification of the oscillation pattern can cause rather peculiar zenith-angle distributions, but the current data do not allow one to exclude the $\nu_{s}$ channel.

While the $\nu_{\tau}$ is quasi-sterile in the detector because of the large mass of the $\tau$-lepton, there is still an important difference to a $\nu_{s}$ because the $\nu_{\tau}$ produces pions in neutral-current collisions such as $\nu N \rightarrow N\nu\pi^{0}$ which can be seen by $\pi^{0} \rightarrow 2\gamma$. With better statistics and a dedicated analysis one may be able to distinguish the $\nu_{\tau}$ and $\nu_{s}$ oscillation channels.

The evidence for atmospheric neutrino oscillations is very compelling, yet an independent confirmation is urgently needed. Hopefully it will come from one of the long-baseline experiments where an accelerator neutrino beam is directed toward a distant detector. The most advanced project is the K2K experiment between KEK and Kamioka with a baseline of 250 km. Other projects include detectors in the Soudan mine at a distance of 730 km from Fermilab, or in the Gran Sasso Laboratory at 732 km from CERN.

2.2 Solar Neutrinos

The Sun, like other hydrogen-burning stars, liberates nuclear binding energy by the effective fusion reaction $4p + 2e^{-} \rightarrow ^{4}\text{He} + 2\nu_{e} + 26.73$ MeV so that its luminosity implies a $\nu_{e}$ flux at Earth of $6.6 \times 10^{10}$ cm$^{-2}$ s$^{-1}$. In detail,
the production of helium involves primarily the pp-chains—the CNO cycle is important in stars more massive than the Sun. The expected solar neutrino flux is shown in Fig. 3, where solid lines are for the three contributions which are most important for the measurements,

\[ \text{pp: \quad} p + p \rightarrow ^2\text{H} + e^+ + \nu_e \quad (E_\nu < 0.420 \text{ MeV}), \]

\[ \text{Beryllium: \quad} e^- + ^7\text{Be} \rightarrow ^7\text{Li} + \nu_e \quad (E_\nu = 0.862 \text{ MeV}), \]

\[ \text{Boron: \quad} p + ^7\text{Be} \rightarrow ^8\text{B} \rightarrow ^8\text{Be}^* + e^+ + \nu_e \quad (E_\nu \lesssim 15 \text{ MeV}). \quad (5) \]

A crucial feature of these reactions is that the beryllium and boron neutrinos both arise from $^7\text{Be}$ which may either capture a proton or an electron so that their relative fluxes depend on the branching ratio between the two reactions.

The solar neutrino flux has been measured in five different experiments with three different spectral response characteristics; the relevant energy range is indicated by the hatched bars above Fig. 3. The radiochemical gallium experiments GALLEX, SAGE and SAGE-GALLEX reach to the lowest energies and pick

**Figure 4: Solar neutrino fluxes measured in five experiments vs. theoretical predictions from a standard solar model.** (Figure courtesy of J. Bahcall.)
up fluxes from all source reactions. The Homestake chlorine experiment\textsuperscript{37} picks up beryllium and boron neutrinos, while the Kamiokande\textsuperscript{38,39} and SuperKamiokande\textsuperscript{38,39} water Cherenkov detectors see only the upper part of the boron flux. All of the experiments see a flux deficit relative to standard-solar model predictions as summarized in Fig. 4 and in a recent overview\textsuperscript{41}.

It has been widely discussed that there is no possibility to account for the measured fluxes by any apparent astrophysical or nuclear-physics modification of the standard solar models so that an explanation in terms of neutrino oscillations is difficult to avoid\textsuperscript{41,42}. Moreover, at something like the 99.8% CL one cannot account for the measurements by an energy-independent global suppression factor\textsuperscript{41}. Therefore, one cannot appeal to neutrino oscillations with an arbitrary \( \Delta m^2_{\nu} \) and a large mixing angle.

One viable possibility are vacuum oscillations with a large mixing angle and a \( \Delta m^2_{\nu} \) around \( 10^{-10} \) eV\(^2\), providing an oscillations length of order the Sun-Earth distance and thus an energy-dependent suppression factor. Second, one can have solutions with \( \Delta m^2_{\nu} \) in the neighborhood of \( 10^{-5} \) eV\(^2\) where the mass difference between the oscillating flavors (energy of order 1 MeV) can be canceled by the neutrino refractive effect in the Sun, leading to resonant or MSW oscillations\textsuperscript{43,44}, again with an energy-dependent suppression factor of the \( \nu_e \) flux. In this case one may have a nearly maximal mixing angle, or a small one as shown in Table 1. The large-angle MSW region does not provide a credible fit for \( \nu_e \rightarrow \nu_\mu,\tau \) oscillations while the other solutions are possible for the \( \nu_e \rightarrow \nu_\mu,\tau \) or \( \nu_e \rightarrow \nu_s \) channels, of course with somewhat different contours of preferred mixing parameters\textsuperscript{41}. It is noteworthy that the spectral distortion of the spectrum of recoil electrons measured at SuperKamiokande seems to single out the vacuum case as the preferred solution\textsuperscript{39}, although this must be considered a rather preliminary conclusion at present.

2.3 LSND

The LSND (Liquid Scintillation Neutrino Detector) experiment is the only case of a pure laboratory experiment which shows indications for neutrino oscillations\textsuperscript{45}. It utilizes a proton beam at the Los Alamos National Laboratory in the US. The protons are directed at a target where neutrinos arise from the same basic mechanism Eq. (1), upper line, that produces them in the atmosphere. From \( \pi^+ \) decay-in-flight one obtains a \( \nu_\mu \) beam of up to 180 MeV while the subsequent decay-at-rest of stopped \( \mu^+ \)'s provides a \( \bar{\nu}_\mu \) beam of less than 53 MeV. The beam should not contain any \( \bar{\nu}_e \)'s; they can be detected by \( \bar{\nu}_e p \rightarrow n e^+ \) in coincidence with \( np \rightarrow d \gamma (2.2 \text{ MeV}) \). For energies above 36 MeV, the 1993–95 data included 22 such events above an expected background of
4.6 ± 0.6; this excess is interpreted as evidence for \( \bar{\nu}_\mu \rightarrow \bar{\nu}_e \) oscillations.

The LSND data favor a large range of \( \nu_e-\nu_\mu \)-mixing parameters. After taking the exclusion regions of other experiments into account, one is left with a sliver of mixing parameters in the range indicated in Table 1. The KARMEN experiment is also sensitive in this range, but has not seen any events. This lack of confirmation, however, does not exclude the LSND evidence as the non-observation of only a few expected events is not a statistically persuasive conflict. Moreover, if one excludes the background-infested 20–36 MeV data in LSND one finds a much broader range of allowed mixing parameters than could have been probed by KARMEN \(^{18}\). Within 2–3 years all of the LSND area will be covered with high sensitivity by MiniBooNE \(^{17}\), a new experiment at Fermilab, which will settle this case.

2.4 Global Interpretation

In Table 1 we summarize the neutrino oscillation channels and mixing parameters indicated by the atmospheric and solar neutrino anomalies and the LSND experiment. Clearly there is no straightforward interpretation because there are too many indications! If only three different mass eigenstates \( m_i \), \( i = 1, 2, 3 \), exist, the mass splittings must satisfy

\[
\sum_{\text{Splittings}} \Delta m^2 = (m_3^2 - m_2^2) + (m_2^2 - m_1^2) + (m_1^2 - m_3^2) = 0, \tag{6}
\]

a trivial condition which is not met by the independent \( \Delta m^2 \) from Table 1. Some of the experiments may not be due to a single \( \Delta m^2 \) but rather to non-trivial three-flavor oscillation patterns \(^{49,50,51,52}\). Even then it appears that

| Experiment       | Favored Channel | \( \Delta m^2 \) [eV\(^2\)] | \( \sin^2 2\theta \) |
|------------------|----------------|-----------------------------|---------------------|
| LSND             | \( \bar{\nu}_\mu \rightarrow \bar{\nu}_e \) | 0.2–10                     | \( (0.2–3) \times 10^{-2} \) |
| Atmospheric      | \( \nu_\mu \rightarrow \nu_\tau \) | \( (1–8) \times 10^{-3} \)  | 0.85–1              |
|                  | \( \nu_\mu \rightarrow \nu_s \) | \( (2–7) \times 10^{-3} \)  | 0.85–1              |
| Solar            | \( \nu_e \rightarrow \text{anything} \) | \( (0.5–8) \times 10^{-10} \) | 0.5–1               |
| Vacuum           | \( \nu_e \rightarrow \text{anything} \) | \( (0.4–1) \times 10^{-5} \)  | \( 10^{-3}–10^{-2} \) |
| MSW (small angle)| \( \nu_e \rightarrow \nu_\mu \) or \( \nu_\tau \) | \( (3–30) \times 10^{-5} \)  | 0.6–1               |
one must ignore some of the experimental evidence or stretch the errors beyond plausible limits to accommodate all experiments in a three-flavor scheme.

If one has to throw out one of the indications, LSND is usually taken as the natural victim because there is no independent confirmation, and because the other cases simply look too strong to be struck from the list. Once LSND has been disposed of, a typical mass and mixing scheme may be as shown in Fig. 5 where the small-angle MSW solution has been taken for solar neutrinos.

\[ \begin{array}{c}
\mathbf{\nu}_\text{solar} \\
\mathbf{\nu}_\text{ATM} \\
\mathbf{\nu}_\mu \\
\mathbf{\nu}_\tau
\end{array} \]

Figure 5: Hierarchical mass and mixing scheme to account for solar and atmospheric neutrinos, the former by the small-angle MSW solution. The flavor content of each mass eigenstate is indicated by the fill-patterns. (Figure 53 reproduced with kind permission of A. Smirnov.)

However, the large mixing angle which is needed to account for the atmospheric neutrino anomaly suggests that more than one mixing angle may be large. Moreover, the spectral distortion observed in SuperKamiokande suggests that the solar vacuum solution may be preferred. Of course, the vastly different values for $\Delta m^2_{\nu}$ implied by atmospheric neutrinos and the solar vacuum solution looks unnatural. Shrugging off this objection, there are several workable schemes involving more than one large mixing angle, for example bi-maximal mixing or threefold maximal mixing.

It is also conceivable that the mass differences are not representative of the masses themselves, i.e. that all three flavors have, say, an eV-mass with small splittings as implied by solar and atmospheric neutrinos (degenerate mass pattern). Of course, such a scheme is very different from the hierarchical patterns that we know in the quark and charged-lepton sectors, but the large mixing angle or angles look very unfamiliar, too. If the neutrino masses are all Majorana, one may still evade bounds on the effective $\nu_e$ Majorana mass
\langle m_{\nu_e}^2 \rangle_{\text{eff}} \) relevant for neutrinoless $\beta\beta$ decay. For example, in the bi-maximal mixing case there is an exact cancellation so that \( \langle m_{\nu_e}^2 \rangle_{\text{eff}} = 0 \) in the limit where the mass differences can be neglected relative to the common mass scale.

At the present time there is no objective reason to ignore LSND. As a consequence, a very radical conclusion follows: there must be four independent mass eigenstates, i.e., at least one low-mass neutrino degree of freedom beyond the three sequential flavors. This fourth flavor $\nu_s$ would have to be sterile with regard to the standard weak interactions. Probably the most natural mass and mixing pattern is one like Fig. 6, but there are also other possibilities.

![Figure 6: Representative four-flavor mass and mixing scheme to account for all experimental evidence. (Figure 63 reproduced with kind permission of A. Smirnov.)](image)

Of course, it would be an extremely radical and unexpected finding if the oscillation experiments had not only turned up evidence for neutrino masses, but for an additional, previously unsuspected low-mass sterile neutrino. A confirmation of LSND by MiniBooNE would make this conclusion difficult to avoid so that this new experiment is perhaps the most urgent current effort in experimental neutrino physics.

## 3 Cosmology

### 3.1 Big-Bang Nucleosynthesis

Massive neutrinos and the existence of sterile neutrinos can have a variety of important cosmological consequences. One immediately wonders if a fourth neutrino flavor is not in conflict with the well-known big-bang nucleosynthesis (BBN) limit on the effective number of thermally excited primordial neutrino degrees of freedom. However, there are several questions. The first
and most obvious one is whether the observationally inferred light-element abundances strictly exclude a fourth flavor at the epoch of BBN. The unfortunate answer is that, while a fourth flavor clearly would make a very significant difference, BBN is not in a position to exclude this possibility with the sort of confidence that would be required to dismiss the sterile-neutrino hypothesis. 58

Second, a sterile neutrino need not attain thermal equilibrium in the first place. It is excited by oscillations in conjunction with collisions so that its contribution to the cosmic energy density at the BBN epoch depends on the mass difference and mixing angle with an active flavor. 59-62 If the atmospheric neutrino anomaly is due to $\nu_\mu \rightarrow \nu_\tau$ oscillations, the large mixing angle and large $\Delta m^2$ imply that the sterile neutrino would be fully excited at the time of BBN. On the other hand, for the small-angle MSW solution or the vacuum solution of the solar neutrino problem, it is barely excited so that the additional energy density is negligible. Therefore, of the different four-flavor patterns BBN favors those where $\nu_e - \nu_\tau$ oscillations solve the solar neutrino problem over those where $\nu_\mu - \nu_\tau$ oscillations explain the atmospheric neutrino anomaly.

Even this conclusion can be avoided if a lepton asymmetry of order $10^{-5}$ exists at the time of the primordial $\nu_\mu \rightarrow \nu_\tau$ oscillations. 66 It may be possible to create such asymmetries among the active neutrinos by oscillations between, say, $\nu_\tau$ ($\bar{\nu}_\tau$) and sterile states, although the exact requirements on the mass and mixing parameters are controversial in some cases. 67-73

Be that as it may, a sterile neutrino provides for a rich oscillation phenomenology in the early universe, but at the same time BBN is not quite enough of a precision tool to distinguish seriously between different four-flavor patterns. As it stands, BBN would benefit more from pinning down the neutrino mass and mixing pattern experimentally than the other way round.

3.2 Dark Matter

Irrespective of the possible existence of a sterile neutrino, it has become difficult to dispute that neutrinos have masses. Therefore, they could play an important role for the cosmological dark matter. Standard calculations in the framework of the big-bang cosmology reveal that the present-day universe contains about 100 cm$^{-3}$ neutrinos and antineutrinos per active flavor, leading to a cosmological mass fraction of

$$\Omega_\nu h^2 = \sum_{i=1}^{3} \frac{m_i}{93 \text{ eV}}, \quad (7)$$
where \( h \) is the Hubble constant in units of 100 km \( s^{-1} \) Mpc\(^{-1}\). The observed age of the universe together with the measured expansion rate reveals that \( \Omega h^2 \lesssim 0.4 \), leading to the most restrictive limit on the masses of all neutrino flavors. Once we believe the current indications for oscillations, the mass differences are so small that this limit reads \( m_\nu \lesssim 13 \text{ eV} \) for the common mass scale of all flavors, roughly identical with the world-averaged tritium endpoint limit on \( m_\nu \) of about 15 eV.

If the neutrino masses were in this range they could be the cosmic dark matter as first pointed out more than 25 years ago. However, it was quickly recognized that neutrinos do not make for a good universal dark matter candidate. The simplest counter-argument (“Tremaine-Gunn-limit”) arises from the phase space of spiral galaxies which cannot accommodate enough neutrinos to account for their dark matter unless the neutrino mass obeys a lower limit. For typical spiral galaxies it is \( m_\nu > 20 \text{ eV} \), for dwarf galaxies even \( m_\nu > 100–200 \text{ eV} \), difficult to reconcile with the cosmological upper limit.

### 3.3 Large-Scale Structure

The Tremaine-Gunn-limit is only the tip of the iceberg of evidence against neutrino dark matter. The most powerful argument arises from cosmic structure formation. At early times the universe was extremely smooth as demonstrated by the tiny amplitude of the temperature fluctuations of the cosmic microwave background radiation across the sky. The present-day distribution of matter, on the other hand, is very clumpy. There are stars, galaxies, clusters of galaxies, and large-scale coherent structures on scales up to about 100 Mpc. A perfectly homogeneous expanding universe stays that way forever. The standard theory for the formation of structure has it that the universe was initially almost, but not quite, perfectly homogeneous, with a tiny modulation of its density field. The action of gravity enhances the density contrast as time goes on, leading to the observed structures.

The outcome of this evolution depends on the initial spectrum of density fluctuations which is usually taken to be approximately flat, i.e. of the “Harrison-Zeldovich-type,” corresponding to the power-law-index \( n = 1 \). However, the effective spectrum relevant for structure formation is the processed spectrum which obtains at the epoch when the universe becomes matter dominated. As the matter which makes up the cosmic fluid can diffuse around, the smallest-scale density fluctuations will be wiped out. This effect is particularly important for weakly interacting particles which can diffuse far while they are relativistic. Low-mass particles stay relativistic for a long time and thus wipe out the primordial fluctuations up to large scales. Massive particles stay put.
earlier and thus have this effect only on small scales. One speaks of “hot dark matter” (HDM) if the particle masses are small enough that all fluctuations are wiped out beyond scales which later correspond to a galaxy. Conversely, “cold dark matter” (CDM) has this effect only on sub-galactic scales.

One way of presenting the results of calculations of structure formation is to show the expected power-spectrum of the present-day matter distribution (Fig. 7) which can be compared to the observed galaxy distribution. The theory of structure formation then predicts the form, but not the amplitude of the spectrum which can be fit either on large scales to the observed temperature fluctuations of the cosmic microwave background radiation as observed by the COBE satellite, or else on small scales to the observed galaxy distribution. Figure 7 illustrates that HDM (neutrinos) suppresses essentially all small-scale structure below a cut-off corresponding to a supercluster scale and thus does not seem to be able to account for the observations.

While cold dark matter works impressively well, it has the problem of producing too much clustering on small scales. Ways out include a primordial power spectrum which is not strictly flat (tilted dark matter), a mix of cold and hot dark matter, or the assumption of a cosmological constant. Currently there is a broad consensus that some variant of a CDM cosmology where structure
forms by gravitational instability from a primordial density fluctuations of approximately the Harrison-Zeldovich type is probably how our universe works.

Thus, while it is widely accepted that neutrinos are not the main dark-matter component, quite conceivably they contribute something like 20%, giving rise to a hot plus cold dark matter (HCDM) scenario which avoids the overproduction of small-scale structure of a pure CDM cosmology.\cite{85,86,87,88,89}

A HDM fraction exceeding about 20% is inconsistent with the size of voids in the galaxy distribution.\cite{90} It was claimed that the HCDM picture with about 20% HDM provides the best fit to all current large-scale structure data.\cite{89,91}

Moreover, if LSND is confirmed, especially with a $\Delta m^2_{\nu}$ of around 6 eV$^2$, there would be a cosmic HDM component of just the right magnitude.\cite{86} The LSND signal and the HCDM cosmologies have become closely intertwined issues.

However, important arguments against a HCDM scenario have appeared. First, a cosmological model with the critical amount of dark matter is hard to reconcile with all the evidence on the matter density; something like 30% looks far more convincing. Moreover, the high-redshift type Ia supernova Hubble diagram now indicates the existence of a cosmological constant $\Lambda$.\cite{92,93,94,95}

If correct, one is naturally led to a critical cosmological model with something like 5% baryonic matter, 25% CDM, and 70% “vacuum energy.” Likewise, the observed abundance of high-redshift ($z \sim 3$) galaxies is reproduced in this type of $\Lambda$CDM model, but not by HCDM.\cite{96} A small amount of HDM is still possible in a $\Lambda$CDM scenario, but not especially needed for anything.\cite{97}

The cosmic large-scale structure is sensitive to small neutrino masses,
whether or not they are needed. Put another way, the unknown common mass scale which is left open by oscillation experiments has a measurable impact on the power spectrum of the large-scale matter distribution. For example, the upcoming Sloan Digital Sky Survey \(^{68}\) will produce precision data where a neutrino mass as small as 0.1 eV makes a noticeable difference \(^{69}\), even though a statistically meaningful neutrino mass limit may not lie far below 1 eV. This is illustrated in Fig. 8 where the expected Sloan sensitivity to the power spectrum of bright red galaxies is compared with theoretical predictions in a universe with the critical mass in dark matter \((\Omega_M = 1)\) and a low-density universe \((\Omega_M = 0.2)\), each time with or without a 1 eV neutrino.

In the long-term future, weak lensing of galaxies by large-scale structure may provide even more precise information on cosmological parameters. An ultimate sensitivity to a neutrino mass as low as 0.1 eV has been suggested \(^{100}\).

3.4 Cosmic Microwave Background Radiation

Another sensitive probe of large-scale structure is the cosmic microwave background radiation (CMBR), and more specifically the power-spectrum of its temperature fluctuations across the sky. The anticipated sky maps of the future MAP \(^{101}\) and PLANCK \(^{102}\) satellite missions have already received advance praise as the “Cosmic Rosetta Stone” \(^{103}\) because of the wealth of cosmological precision information they are expected to reveal \(^{104,105,106,107}\).

CMBR sky maps are characterized by their fluctuation spectrum \(C_\ell = \langle a_{\ell m} a^*_{\ell m} \rangle\) where \(a_{\ell m}\) are the coefficients of a spherical-harmonic expansion. Figure 9 (solid line) shows \(C_\ell\) for standard cold dark matter (SCDM) with \(N_{\text{eff}} = 3\) for the effective number of neutrino degrees of freedom. Sterile neutrinos increase the radiation content and thus modify this pattern in a characteristic way illustrated by the dotted line, which corresponds to \(N_{\text{eff}} = 4\).

While this shift appears small, the lower panel of Fig. 9 shows that for \(\ell \gtrsim 200\) it is large on the scale of the expected measurement precision. It is fundamentally limited by the “cosmic variance” \(\Delta C_\ell / C_\ell = \sqrt{2/(2\ell + 1)}\), i.e. by the fact that at our given location in the universe we can measure only \(2\ell + 1\) numbers \(a_{\ell m}\) to obtain the expectation value \(\langle a_{\ell m} a^*_{\ell m} \rangle\). The actual sensitivity will be worse, but the cosmic variance gives us an optimistic idea of what one may hope to achieve. The true sensitivity to \(\Delta N_{\text{eff}}\) is further limited by our lack of knowledge of several other cosmological parameters. Even then it is safe to assume that we are sensitive to \(|\Delta N_{\text{eff}}| \lesssim 0.3\), and much better with prior knowledge of other parameters \(^{105}\). Thus it appears that the CMBR is a more powerful tool to measure \(N_{\text{eff}}\) than the standard BBN argument, although a more pessimistic assessment was put forth in a more recent analysis \(^{107}\).
If LSND is right, some of the neutrinos have eV masses which imprint themselves on the CMBR fluctuation spectrum. For example, if the atmospheric neutrino anomaly is due to $\nu_\mu - \nu_\tau$-oscillations, we will have approximately $N_{\text{eff}} = 4$, and two of these states will have an eV-range mass. The CMBR imprint of this scenario is illustrated with the dashed curve in Fig. 9 where $\Omega_\nu = 0.2$. With $\Omega_\nu h^2 = 2m_\nu/93$ eV and taking $h = 0.5$ this implies $m_\nu \approx 2.4$ eV, well within the range suggested by LSND.

The range of $\Delta N_{\text{eff}}$ and the HDM fraction that can be determined by the future CMBR sky maps, together with large-scale galaxy surveys, cannot be foretold with certainty, but surely these cosmological precision observables...
are significantly affected by the currently debated neutrino mass and mixing patterns. Cosmology may be our best bet to pin down the overall neutrino mass scale which is left undetermined by oscillation experiments.

4 Supernova Physics

4.1 Kinematical Mass Limits

When SN 1987A exploded on 23 February 1987 in the Large Magellanic Cloud at a distance of about 50 kpc (165,000 lyr), it produced the third case of a measured neutrino signal from an astrophysical source after the Sun and the Earth’s atmosphere. Therefore, we turn to the role of masses and mixings for SN neutrinos in general, and for the SN 1987A burst in particular.

A type II SN explosion marks the end of the life of a massive star ($M \gtrsim 8\, \text{M}_\odot$) which has developed a degenerate iron core, surrounded by several burning shells. As the core reaches its Chandrasekhar limit of 1–2 $\text{M}_\odot$ (solar masses) it becomes unstable and collapses down to nuclear density ($3 \times 10^{14}\, \text{g cm}^{-3}$) where the equation of state stiffens and the implosion is halted. At this point a shock wave forms which ejects the mantle of the progenitor star—the SN explosion is the reversed core implosion.

At about nuclear density and a temperature of several 10 MeV the newly formed neutron star is opaque to neutrinos which are thus emitted from a shell at about unit optical depth, the “neutrino sphere,” crudely with a thermal spectrum. One expects that the total binding energy is roughly equipartioned between all (anti)neutrino flavor degrees of freedom and that it is emitted within several seconds. This picture agrees well with the SN 1987A observations in the Kamiokande and IMB water Cherenkov detectors and the Baksan Scintillator Telescope, which were all primarily sensitive to the positrons from the $\bar{\nu}_e + p \rightarrow n + e^+$ capture reaction.

A neutrino mass can manifest itself by a time-of-flight dispersion of the SN burst. The neutrino arrival time from a distance $D$ is delayed by

$$\Delta t = 2.57 \, s \left( \frac{D}{50\, \text{kpc}} \right) \left( \frac{10\, \text{MeV}}{E_\nu} \right)^2 \left( \frac{m_\nu}{10\, \text{eV}} \right)^2.$$ 

As the $\bar{\nu}_e$’s from SN 1987A were registered within a few seconds and had energies in the 10 MeV range, the $m_\nu$ limit is around 10 eV. Detailed analyses reveal that the pulse duration is consistently explained by the SN cooling time and that $m_\nu \lesssim 20\, \text{eV}$ is implied at something like 95% CL.
The high-statistics observation of a future galactic SN with a large detector like SuperKamiokande allows one to improve the $m_{\nu_e}$-sensitivity to about 3 eV because one can use the fast rise-time of the signal as a dispersion measure rather than the overall burst duration itself\cite{122}. On the other hand, the neutral-current signal in a large water Cherenkov detector like SuperKamiokande or SNO provides a direct handle on $m_{\nu_\mu}$ and $m_{\nu_\tau}$ of no better than 30 eV\cite{124,125,126,127}. Even with a future neutral-current detector like OMNIS it is not realistically possible to probe $m_{\nu_\mu}$ and $m_{\nu_\tau}$ down to a few eV\cite{128,129}.

\subsection{SN 1987A and Flavor Oscillations}

While the SN 1987A limit on $m_{\nu_e}$ is not truly interesting for the current debate, the event energies bear on the large-angle solutions of the solar neutrino problem, and especially on the vacuum solution. In typical numerical simulations one finds for the average energies for the different flavors\cite{130}

$$
\langle E_{\nu} \rangle = \begin{cases} 
10-12 \text{ MeV} & \text{for } \nu_e, \\
14-17 \text{ MeV} & \text{for } \bar{\nu}_e, \\
24-27 \text{ MeV} & \text{for } \nu_{\mu,\tau} \text{ and } \bar{\nu}_{\mu,\tau},
\end{cases}
$$

so that $\langle E_{\nu_e} \rangle : \langle E_{\bar{\nu}_e} \rangle : \langle E_{\text{others}} \rangle \approx \frac{2}{3} : \frac{1}{3} : \frac{1}{2}$. Large mixing angle oscillations between $\bar{\nu}_e$ and $\bar{\nu}_\mu$ would partially swap their fluxes and thus “stiffen” the $\bar{\nu}_e$ spectrum observable at Earth\cite{122,131,132,133,134}. (We take $\bar{\nu}_\mu$ to stand for either $\bar{\nu}_{\mu}$ or $\bar{\nu}_{\tau}$.) Therefore, some of the SN 1987A events would have been oscillated $\bar{\nu}_\mu$’s which should have been correspondingly more energetic.

A maximum-likelihood analysis of the $\bar{\nu}_e$ spectral temperature and the neutron-star binding energy inferred from the Kamiokande\cite{117} and IMB\cite{118} data (Fig. 10) reveals that even in the no-oscillation case there is only marginal overlap with the theoretical expectation of Eq. (10). The observed neutrinos were softer than predicted, especially at Kamiokande. Including a spectral swap exacerbates this problem in that the energies should have been even higher. In Fig. 10 we show 95% likelihood contours for the inferred $\bar{\nu}_e$ spectral temperature $T_{\bar{\nu}_e} = \langle E_{\bar{\nu}_e} \rangle / 3$ and the neutron-star binding energy $E_b$ for maximum $\bar{\nu}_e$-$\bar{\nu}_\mu$-mixing and for several values of $\tau = T_{\bar{\nu}_e}/T_{\nu_e}$. Even for moderate spectral differences a maximum mixing between $\bar{\nu}_e$ and the other flavors causes a conflict with the SN 1987A data\cite{133,134}.

It may be premature to exclude the solar vacuum solution on these grounds as the spectral differences may have been overestimated. They arise because of flavor-dependent opacities. The electron-flavored neutrinos are trapped by
Figure 10: Best-fit values for the spectral $\bar{\nu}_e$ temperature $T_{\bar{\nu}_e}$ and the neutron-star binding energy $E_b$, as well as contours of constant likelihood corresponding to 95% confidence regions. They are based on a joint analysis between the Kamiokande and IMB data, assuming maximum mixing and the indicated values for $\tau = T_{\mu}/T_{\nu_e}$, where $\tau = 1$ corresponds to no oscillations. The hatched region represents the predictions of Eqs. (8) and (10).

$\nu_e n \rightarrow pe^-$ and $\bar{\nu}_e p \rightarrow ne^+$. The other flavors interact by neutral-current collisions which have smaller cross sections so that these particles emerge from deeper and hotter layers. They escape from their “transport sphere” where collisions are no longer effective, but most critical for their spectrum is the “energy sphere” where they last exchanged energy with the medium. Electron scattering $\nu e^- \rightarrow e^- \nu$ was taken to dominate for energy-exchange and $e^+e^- \rightarrow \nu\bar{\nu}$ for pair production. However, the dominant pair-process is nucleonic bremsstrahlung $NN \rightarrow NN\nu\bar{\nu}$, the dominant energy-exchange processes are recoils and inelasticities in $\nu N \rightarrow N\nu$ scattering. Including these effects clearly makes the $\bar{\nu}_\mu$ spectrum more similar to $\bar{\nu}_e$. A preliminary estimate suggests that the remaining spectral differences may be small enough to avoid a conflict between SN 1987A and the solar vacuum solution. Since neutrino oscillations can be crucial for the interpretation of the signal from a future galactic SN, one should indeed spend more effort at understanding details of the spectra formation process.

An interesting case which does not depend on the spectral differences is the “prompt $\nu_e$ burst,” originating from the deleptonization of the outer core layers at about 100 ms after bounce when the shock wave breaks through the
edge of the collapsed core. This “deleptonization burst” propagates through the mantle and envelope of the progenitor star so that resonant oscillations take place for a large range of mixing parameters between $\nu_e$ and some other flavor, notably for some of those values where the MSW effect operates in the Sun. In a Cherenkov detector one can see this burst by $\nu_e$-$e$-scattering which is forward peaked, but one would have expected only a fraction of an event from SN 1987A. The first event in Kamiokande may be attributed to this signal, but this interpretation is statistically insignificant. The experimental signal of the prompt $\nu_e$ burst from a future galactic SN is closely intertwined with the mixing parameters which solve the solar neutrino problem.

4.3 Flavor Oscillations and Supernova Physics

Flavor oscillations can have interesting ramifications for SN physics itself, independently of neutrino flux measurements at Earth. As galactic SNe are rare (one every few decades or even less) it is not guaranteed that we will observe neutrinos from another SN anytime soon. Therefore, it is even more important to use the SN phenomenon itself as a laboratory for neutrino physics.

For example, flavor oscillations can help with the explosion. The standard scenario of a type II SN explosion has it that a shock wave forms near the edge of the core when its collapse halts at nuclear density and that this shock wave ejects the mantle of the progenitor star. However, in typical numerical

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure11.png}
\caption{Mixing parameters between $\nu_e$ and $\nu_\mu$ or $\nu_\tau$ where a spectral swap would help explode supernovae and where it would prevent r-process nucleosynthesis.}
\end{figure}
calculations the shock wave stalls so that this “prompt explosion” scenario does not seem to work. In the “delayed explosion” picture the shock wave is revived by neutrino heating, perhaps in conjunction with convection, but even then it appears difficult to obtain a successful or sufficiently energetic explosion. The efficiency of neutrino heating can increase by resonant flavor oscillations which swap the $\nu_e$ flux with, say, the $\nu_\tau$ one. Therefore, what passes through the shock wave as a $\nu_e$ was born as a $\nu_\tau$ at the proto neutron star surface. It has on average higher energies and thus is more effective at transferring energy. In Fig. 1 the shaded range of mixing parameters is where SNe are helped to explode, assuming a “normal” neutrino mass spectrum with $m_{\nu_e} < m_{\nu_\tau}$. Below the shaded region the resonant oscillations take place beyond the shock wave and thus do not affect the explosion.

A few seconds after core bounce the shock wave has long taken off, leaving behind a relatively dilute “hot bubble” above the neutron-star surface. This region is one suspected site for the r-process heavy-element synthesis, which requires a neutron-rich environment. The neutron-to-proton ratio, which is governed by the beta reactions $\nu_e + n \rightarrow p + e^-$ and $\bar{\nu}_e + p \rightarrow n + e^+$, is shifted to a neutron-rich phase if $\langle E_{\nu_e} \rangle < \langle E_{\bar{\nu}_e} \rangle$ as for standard neutrino spectra. Resonant oscillations can again swap the $\nu_e$ flux with another one, inverting this hierarchy of energies. In the hatched range of mixing parameters shown in Fig. 1 the r-process would be disturbed in conflict with the upper range of LSND-inspired mass differences. On the other hand, oscillations $\nu_e \rightarrow \nu_s$ into a sterile neutrino could actually help the r-process by depleting the neutron-stealing $\nu_e$ flux.

### 4.4 Pulsar Kicks by Oscillations?

Radio pulsars often move with velocities of several 100 km s$^{-1}$, a phenomenon yet to be explained. The acceleration probably takes place in the context of their formation in a core-collapse SN, i.e. they likely receive a kick at birth. One explanation appeals to a “neutrino rocket” because the momentum carried by the neutrino burst is so large that an emission anisotropy as small as 1% suffices to account for a recoil of about 300 km s$^{-1}$. However, even such a small anisotropy is difficult to explain.

Pulsars tend to have strong magnetic fields which may well be suspected to cause the asymmetry. The neutrino refractive index depends on the direction of the neutrino momentum relative to $\mathbf{B}$. For suitable conditions, resonant neutrino oscillations occur between the neutrinospheres of $\nu_e$ and $\nu_\tau$, deforming the effective $\nu_\tau$ sphere. The $\nu_\tau$’s would thus emerge from regions of varying effective temperature and thus, it was argued, would be emitted anisotropi-
This argument was then taken up in several papers with modified neutrino oscillation scenarios. Unfortunately, this intriguing idea does not work for plausible magnetic field strengths. The oscillations take place in the “atmosphere” of the neutron star, while the neutrino flux is fixed much deeper inside. The atmosphere adjusts itself to transport the neutrino flux, not the other way round. Neutrino oscillations in the atmosphere leave the overall flux unchanged except for a higher-order backreaction effect which obtains because of the anisotropically modified atmospheric structure. It may still be that a neutrino rocket effect is responsible for the pulsar kicks, but the cause for the anisotropy remains unclear and if it is related to nonstandard neutrino properties.

4.5 Neutrino Mass Limit from Neutron-Star Stability?

In a thought-provoking paper, it was recently claimed that neutron stars provided a lower neutrino mass limit of $m_\nu \gtrsim 0.4$ eV. Two-neutrino exchange between fermions gives rise to a long-range force. A neutrino may also pass around several fermions, so to speak, producing a much smaller potential. This multibody neutrino exchange, it was argued, would be a huge effect in neutron stars because combinatorial factors among many neutrons win out against the smallness of the potential for a given set of them. One way out is to suppress the long-range nature of neutrino exchange by a nonzero $m_\nu$.

This idea triggered a series of papers where it was shown that a proper resummation of a seemingly divergent series of terms leads to a well-behaved and small “neutron-star self-energy,” invalidating the claim of a lower neutrino mass limit. As naively expected, there is no mysterious long-range force from neutrino exchange, but these papers are still interesting reading for anyone interested in questions of neutrino physics in media.

5 Neutrino Astronomy

5.1 Neutrino Telescopes

For twenty years after the first observation of solar neutrinos at the Homestake detector, neutrino astronomy remained a one-experiment field. The SN 1987A neutrino observations mark a turning point—the number of experiments and observatories has multiplied since about that time, with more than a dozen previous, operating or projected neutrino detectors measuring solar and atmospheric neutrinos or searching for a new SN burst. The neutrino sky at low energies is dominated by these sources with a solar $\nu_e$ flux of around...
6.6 \times 10^{10} \text{ cm}^{-2} \text{ s}^{-1} \text{ in the MeV range and that from a SN at a distance of 10 kpc of around } 3 \times 10^{12} \text{ cm}^{-2} \text{ s}^{-1} \text{ in the 10–100 MeV range during the burst of a few seconds. At around 1 GeV the atmospheric neutrino flux for all flavors together and integrated over all angles is } dN_\nu/d\ln E_\nu \approx 0.7 \text{ cm}^{-2} \text{ s}^{-1}, \text{ dropping with energy approximately as } E_\nu^{-2}.

A new development is the emergence of huge neutrino telescopes with the goal of observing astrophysical sources of neutrinos with energies in the TeV range and beyond \cite{171,172,173}. The existence of cosmic rays with energies reaching beyond \(10^{20} \text{ eV}\) proves that they must have been accelerated somewhere, but the nature of the accelerators remains mysterious. Protons are deflected in the micro-Gauss galactic magnetic field so that the cosmic rays hitting the Earth do not point back to their sources, a problem not shared by neutrinos. High-energy neutrinos are expected from “cosmic beam dumps” whenever the protons interact with matter or even photons to produce pions— the Earth’s atmosphere as a neutrino source is the simplest case in point.

Estimates of the expected neutrino fluxes vary, but certainly one needs detectors far exceeding the size of SuperKamiokande. For a useful neutrino Cherenkov telescope one probably needs a cubic-kilometer of water or ice instrumented with photomultipliers which can be placed on a grid with a typical spacing of order 30 m. There are now several such utopian-sounding projects on their way. A small but functioning instrument has been deployed in Lake Baikal \cite{174} but probably it will not grow to the km\(^3\) scale. Two Mediterranean projects, NESTOR \cite{175} and ANTARES \cite{176} are in the R&D and feasibility-study phase. At present the most advanced detector with a realistic km\(^3\) perspective is AMANDA \cite{177} at the South Pole (Fig. 1). The antarctic ice is used both as a Cherenkov medium and as a mechanical support structure for strings of photomultipliers which are frozen into 2 km deep holes.

The main focus of these exciting projects is neutrino astronomy, i.e. to study the sky in a new form of radiation and to learn about the nature of the astrophysical sources. However, high-energy neutrino astronomy has several important ramifications of direct particle-physics interest.

5.2 Search for Particle Dark Matter

First, one may search for dark matter in the form of weakly interacting massive particles (WIMPs), especially in the guise of the supersymmetric neutralinos. The case for these particles has become stronger as massive neutrinos no longer seem tenable as a main dark-matter constituent. Galactic WIMPs are accreted by the Sun or Earth where they annihilate with each other, leading to a secondary GeV–TeV neutrino flux. Depending on details of the assumed su-
Figure 12: Schematic view of the AMANDA South Pole high-energy neutrino telescope \cite{77}.
(Figure reproduced with permission of F. Halzen.)
persymmetric model, this “indirect” method to search for particle dark matter is competitive with the direct laboratory experiments.\textsuperscript{178,179}

5.3 Tau-Neutrinos from Astrophysical Sources

Neutrinos produced in cosmic beam dumps should have the same flavor content as those produced in the atmosphere. If atmospheric neutrinos indeed oscillate, so do the ones from high-energy astrophysical sources. If the $\nu_\mu \rightarrow \nu_\tau$ oscillation channel is what explains the atmospheric anomaly, then the astrophysical beam dumps produce a flux which includes high-energy $\nu_\tau$'s.

One signature in a Cherenkov detector are so-called double-bang events\textsuperscript{180} which consist of a big hadronic shower from the initial $\nu_\tau$ interaction, a muon-like $\tau$-track, and then a second big particle cascade when the $\tau$ decays. This could be 100 m downstream from the first interaction if the primary energy was in the PeV ($10^{15}$ eV) range as expected from active galactic nuclei (AGNs) as neutrino sources.\textsuperscript{181} However, such signatures may be difficult to detect in a first-generation telescope like AMANDA.

The Earth is opaque to neutrinos with energies above something like 100 TeV, but $\nu_\tau$'s can still make it to the detector from below.\textsuperscript{182} The main idea is that a $\tau$ produced in a charged-current interaction of the primary $\nu_\tau$ decays back into a $\nu_\tau$ before losing much energy, thereby piling up $\nu_\tau$'s at energies around 100 TeV. Moreover, this effect would manifest itself by a flat zenith-angle dependence of source intensity at the highest energies.\textsuperscript{182}

The atmospheric neutrino anomaly has rather immediate consequences for high-energy neutrino astronomy!

5.4 Neutrino Masses

Besides AGNs, gamma-ray bursts are one of the favored suspects for producing the highest energy cosmic rays and for producing high-energy neutrinos.\textsuperscript{183} Their pulsed nature allows one to search for neutrino masses by time-of-flight dispersion in analogy to the SN 1987A mass limit. Since typical gamma-ray bursts are at cosmological distances of order 1000 Mpc, one gains enormously in Eq. (9) relative to SN 1987A, but of course the final mass sensitivity depends on the time-structure (perhaps as short as milliseconds) and the observed neutrino energies.

If neutrinos with energies as high as $10^{22}$ eV are copiously produced in astrophysical sources, and if eV-mass neutrinos exist as a hot-dark matter component and are locally clustered, then high-energy particle cascades would be initiated which could produce, as secondary products, the highest-
energy observed cosmic rays which have energies beyond $10^{20}$ eV. The universe is opaque for protons above $4 \times 10^{19}$ eV, the Greisen-Zatsepin-Kuzmin cutoff, due to photo-pion production on the cosmic microwave radiation. Therefore, the highest-energy cosmic rays, if they are protons, must have a local source, but the observed events do not point toward any plausible structure which might serve as such. Neutrinos thus offer one of many speculative explanations for the puzzle of the highest-energy cosmic rays.

6 Neutrino Electromagnetic Properties

6.1 Form Factors

A survey of neutrino astrophysics would be incomplete without a discussion of neutrino electromagnetic properties which could have several important astrophysical consequences. The most general neutrino interaction with the electromagnetic field is:

$$\mathcal{L}_{\text{int}} = -F_1 \bar{\psi} \gamma_\mu \psi A^\mu - G_1 \bar{\psi} \gamma_\mu \gamma_5 \psi \partial_\mu F^{\mu\nu} - \frac{1}{2} \bar{\psi} \sigma_{\mu\nu} (F_2 + G_2 \gamma_5) \psi F^{\mu\nu},$$

(11)

where $\psi$ is the neutrino field, $A^\mu$ the electromagnetic vector potential, and $F^{\mu\nu}$ the field-strength tensor. The form factors are functions of $Q^2$ with $Q$ the energy-momentum transfer. In the $Q^2 \to 0$ limit $F_1$ is a charge, $G_1$ an anapole moment, $F_2$ a magnetic, and $G_2$ an electric dipole moment.

Charge neutrality implies $F_1(0) = 0$. What remains is a charge radius which, like the anapole moment, vanishes in the $Q^2 \to 0$ limit. Therefore, it provides for a contact interaction and as such a correction to processes with $Z^0$ exchange. As astrophysics provides no precision test for the effective strength of neutral-current interactions, these form factors are best probed in laboratory experiments.

Therefore, the only astrophysically interesting possibility are magnetic and electric dipole and transition moments. If the standard model is extended to include neutrino Dirac masses, the magnetic dipole moment is $\mu_\nu = 3.20 \times 10^{-19} \mu_B m_\nu / eV$ where $\mu_B = e/2m_e$ is the Bohr magneton. An electric dipole moment $e_\nu$ violates CP, and both are forbidden for Majorana neutrinos. Flavor mixing implies electric and magnetic transition moments for both Dirac and Majorana neutrinos, but they are even smaller due to a GIM cancelation. Neutrino electromagnetic form factors which are large enough to be of experimental or astrophysical interest require a more radical extension of the standard model, for example the existence of right-handed currents.
6.2 Astrophysical Limits

Assuming that neutrinos have nonstandard electric or magnetic dipole or transition moments, how large can they be? Astrophysics, not laboratory experiments, provides the most restrictive limits. Dipole or transition moments allow for several interesting processes (Fig. 13). For the purpose of deriving limits, the most important case is $\gamma \rightarrow \nu \bar{\nu}$ which is kinematically possible in a plasma because the photon acquires a dispersion relation which roughly amounts to an effective mass. Even without anomalous couplings, the plasmon decay proceeds because the charged particles of the medium provide an effective neutrino-photon interaction\(^{192,193,194}\). Put another way, even standard neutrinos have nonvanishing electromagnetic form factors in a medium\(^{195,196}\).

The standard plasma process dominates the neutrino production in white dwarfs or the degenerate helium core of globular-cluster red giants. The presence of a direct neutrino-photon coupling by a dipole or transition moment enhances the neutrino losses, delaying the ignition of helium. Observations of globular-cluster stars thus reveal a limit\(^{197,198,199,200,201}\)

$$\mu_\nu \lesssim 3 \times 10^{-12} \mu_B,$$  \hfill (12)

applicable to magnetic and electric dipole and transition moments for Dirac and Majorana neutrinos. Of course, the final-state neutrinos must be lighter than the photon plasma mass of around 10 keV for the relevant conditions. A
slightly weaker bound obtains from the white-dwarf luminosity function. Right-handed (sterile) states are produced in electromagnetic spin-flip collisions if neutrinos have Dirac dipole or transition moments. The duration of the SN 1987A neutrino signal precludes excessive cooling by sterile states, yielding a limit on $\mu_{\nu_e}^{\text{(Dirac)}}$ which is numerically equivalent to Eq. (12) 203, 204.

The corresponding laboratory limits are much weaker. The most restrictive bound is $\mu_{\nu_e} < 1.8 \times 10^{-10} \mu_B$ at 90% CL from a measurement of the $\bar{\nu}_e$-e-scattering cross section involving a reactor source. A significant improvement should become possible with the MUNU experiment 205, but it is unlikely that the globular-cluster limit can be reached anytime soon.

A neutrino mass eigenstate $\nu_i$ may decay to another one $\nu_j$ by the emission of a photon, where the only contributing form factors are the magnetic and electric transition moments. The inverse radiative lifetime is found to be

$$\tau_{\gamma}^{-1} = \frac{|\mu_{ij}|^2 + |\epsilon_{ij}|^2}{8\pi} \left( \frac{m_i^2 - m_j^2}{m_i} \right)^3$$

$$= 5.308 \, \text{s}^{-1} \left( \frac{\mu_{\text{eff}}}{\mu_B} \right)^2 \left( \frac{m_i^2 - m_j^2}{m_i^2} \right)^3 \left( \frac{m_i}{\text{eV}} \right)^3,$$  \hspace{0.5cm} (13)

where $\mu_{ij}$ and $\epsilon_{ij}$ are the transition moments while $|\mu_{\text{eff}}|^2 \equiv |\mu_{ij}|^2 + |\epsilon_{ij}|^2$.

![Figure 14: Astrophysical limits on neutrino dipole moments. The light-shaded and dark-shaded exclusion range is from the absence of excessive cosmic diffuse background photons. The dashed line represents the approximation formula in Eq. (14), bottom line.](image-url)
Radiative neutrino decays have been constrained from the absence of decay photons of reactor $\bar{\nu}_e$ fluxes, the solar $\nu_e$ flux, and the SN 1987A neutrino burst. For $m_\nu \equiv m_i \gg m_j$ these limits can be expressed as

$$\frac{\mu_{\text{eff}}}{\mu_B} \lesssim \begin{cases} 
0.9 \times 10^{-1} \,(eV/m_\nu)^2 & \text{Reactor ($\bar{\nu}_e$)}, \\
0.5 \times 10^{-1} \,(eV/m_\nu)^2 & \text{Sun ($\nu_e$)}, \\
1.5 \times 10^{-8} \,(eV/m_\nu)^2 & \text{SN 1987A (all flavors)}, \\
1.0 \times 10^{-11} \,(eV/m_\nu)^{9/4} & \text{Cosmic background (all flavors)}. 
\end{cases} \quad (14)$$

In this form the SN 1987A limit applies for $m_\nu \lesssim 40$ eV. The decay of cosmic background neutrinos would contribute to the diffuse photon backgrounds, excluding the shaded areas in Fig. 14. They are approximately delineated by the dashed line, corresponding to the analytic expression in Eq. (14). More restrictive limits obtain for certain masses above 3 eV from the absence of emission features from several galaxy clusters.

For low-mass neutrinos the $m_\nu^3$ phase-space factor in Eq. (13) is so punishing that the globular-cluster limit is the most restrictive one for $m_\nu$ below a few eV, i.e. in the mass range which today appears favored from neutrino oscillation experiments. Turning this around, if neutrino mass differences are indeed as small as currently believed, the globular-cluster limit implies that radiative neutrino decays do not have observable consequences.

### 6.3 Spin and Spin-Flavor Precession

Neutrinos with magnetic or electric dipole moments spin-precess in external magnetic fields, an effect which may have a number of astrophysical consequences for $\mu_\nu$-values below the globular-cluster limit of Eq. (13). For example, solar neutrinos can precess into sterile and thus undetectable states in the Sun’s magnetic field. The same for SN neutrinos in the galactic magnetic field where an important effect obtains for $\mu_\nu \gtrsim 10^{-12} \mu_B$. Moreover, the high-energy sterile states emitted by spin-flip collisions from the inner SN core could precess back into active ones and cause events with anomalously high energies in SN neutrino detectors, an effect which probably requires $\mu_\nu(\text{Dirac}) \lesssim 10^{-12} \mu_B$ from the SN 1987A signal. For the same $\mu_\nu$-range one may expect an anomalous rate of energy transfer to the shock wave in a SN, helping with the explosion.

The refractive energy shift in a medium for active neutrinos relative to sterile ones creates a barrier to spin precessions. The neutrino mass difference has the same effect if the precession is between different flavors through...
a transition moment\textsuperscript{231}. Combining the effects one arrives at spin-flavor precession in a medium. The mass difference and the refractive term can cancel, leading to resonant oscillations in the spirit of the MSW effect\textsuperscript{233, 234, 235, 236}.

Large magnetic fields exist in SN cores so that spin-flavor precession could play an important role, with possible consequences for the explosion mechanism, r-process nucleosynthesis, or the measurable neutrino signal\textsuperscript{237, 238, 239, 240, 241}.

The downside of this richness of phenomena is that there are so many unknown parameters (electromagnetic neutrino properties, masses, mixing angles) as well as the unknown magnetic field strength and distribution that it is difficult to come up with reliable limits or requirements on neutrino properties. The SN phenomenon is probably too complicated to serve as a laboratory to pin down electromagnetic neutrino properties, but it clearly is an environment where these properties could have far-reaching consequences.

Resonant spin-flavor precessions can explain all solar neutrino data\textsuperscript{242, 243}, but require somewhat large toroidal magnetic fields in the Sun since the neutrino magnetic (transition) moments have to obey the globular-cluster limit of Eq. 12. The main original motivation for magnetically induced oscillations was an apparent correlation between the Homestake solar neutrino data and indicators of solar magnetic activity. Very recent re-analyses reveal that there is no significant correlation with Sun spots\textsuperscript{244}, but also that the hypothesis of a constant flux should be rejected with a significance level of 0.1–6%, depending on the test\textsuperscript{245}. For Majorana neutrinos, the spin-flavor precession amounts to transitions between neutrinos and antineutrinos. The observation of antineutrinos from the Sun would be a diagnostic for this effect\textsuperscript{246, 247, 248}, and probably the only convincing one.

7 Conclusions

As it stands, the most titillating question of neutrino physics no longer is if these elusive particles have masses at all, but rather if a fourth, hitherto unsuspected and otherwise noninteracting degree of freedom exists to reconcile all current indications for neutrino oscillations. If shockingly this were the case, the mass differences suggested by LSND would imply that neutrinos are significant as a hot dark matter component, corresponding to an eV-mass for one or two flavors, which is what nowadays one means with a “cosmologically significant neutrino mass.” Sterile neutrinos and a cosmological hot dark matter component have become closely intertwined issues.

Oscillation experiments reveal only mass differences, leaving a common offset from zero undetermined. Even if LSND is right, the common mass scale may exceed the indicated mass difference, and if LNSD is wrong and
sterile neutrinos do not exist, the sequential neutrinos could still have nearly
degenerate eV-masses and play a role for hot dark matter. Fixing the common
mass scale may soon become the major challenge of neutrino physics.

There are few realistic opportunities to achieve this goal. While neutrinoless \(\beta\beta\) decay experiments and precise tritium endpoint \(\beta\)-spectra remain
crucial, cosmology likely will play a key role for this task. The cosmological
precision information expected from the MAP and PLANCK microwave back-
ground missions and from large-scale redshift surveys are in principle sensitive
to sub-0.1 eV masses. Whether or not they will actually pin down such a small
mass remains to be seen, but surely they cannot ignore it as one of about a
dozen nontrivial cosmological parameters which are not fixed by other data.

A direct kinematical mass limit from signal dispersion of a future galactic
supernova could get down to about 3 eV for \(\nu_e\), probably not good enough for
the questions at hand. If high-energy neutrinos from pulsed sources such as
gamma-ray bursts are observed in upcoming neutrino telescopes one may get
down to much smaller masses.

The atmospheric neutrino anomaly requires a large mixing angle, suggest-
ing that all mixing angles in the neutrino sector could be large, in blunt contrast
to what is observed in the quark sector. A large mixing angle between \(\nu_e\) and
other flavors radically changes the interpretation of the SN 1987A neutrino
signal and that from a future galactic SN. Therefore, it is of paramount im-
portance to develop a better theoretical understanding of the neutrino spectra
formation in SNe to see if swapping flavors by oscillations indeed has significant
and observable effects. Apart from this important issue it does not look as if
neutrino oscillations had much to do with SN physics itself, i.e. with the ex-
losion mechanism, pulsar kicks, or r-process nucleosynthesis, except perhaps
if sterile neutrinos exist.

The large mixing angle implied by atmospheric neutrinos definitely means
that the neutrinos from “cosmic beam dumps” have a modified flavor spectrum,
presumably containing a large fraction of \(\nu_\tau\)’s, which produce unique signatures
in high-energy neutrino telescopes.

Neutrino physics and neutrino astrophysics are at the cross roads. On the
one hand, it is now almost impossible to deny that neutrinos oscillate and thus
presumably have small masses. On the other hand, unless a sterile neutrino
truly exists, there is a sense that neutrino masses are too small to be of very
much cosmological or astrophysical interest. Neutrino astrophysics could turn
out to be more interesting than one would have originally suspected, or more
boring, depending on whether sterile states exist or not.

Either way, it may not be long until the neutrino mass and mixing pattern
has been reconstructed. The main beneficiary may be neutrino astronomy. As
we better understand the behavior of the neutrino beam from distant sources, neutrino astronomy will return to its roots and focus on the physics of the sources rather than worrying about the behavior of the radiation. It may not be long until flavor oscillations in neutrino astronomy are as commonplace a phenomenon as the Faraday effect in radio astronomy!

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