Astrophysical and Cosmological Neutrinos

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Summary. — We review the role of massive neutrinos in astrophysics and cosmology, assuming that the oscillation interpretation of solar and atmospheric neutrinos is correct. In particular, we discuss cosmological mass limits, neutrino flavor oscillations in the early universe, leptogenesis, and neutrinos in core-collapse supernovae.
1. – Introduction

The connection between neutrinos and astrophysics and cosmology is one of the traditional pillars of astroparticle physics. On the one hand side the intrinsic properties of neutrinos are difficult to measure; the “heavenly laboratories” provide invaluable complementary information [1, 2, 3]. On the other hand side neutrinos dominate the dynamics of the radiation dominated universe and of core-collapse supernovae and are important cooling agents even for ordinary stars. Knowing the intrinsic neutrino properties is crucial for our understanding of various astrophysical and cosmological phenomena.

Yet the focus of neutrino astrophysics and cosmology is changing in the light of what is beginning to be the established wisdom. Pure laboratory experiments will soon overtake solar and atmospheric neutrinos at measuring the mixing parameters. While precision cosmology continues to provide the most restrictive limit on neutrino masses, the importance of astrophysics and cosmology as neutrino laboratories is probably diminishing. Rather, the measured neutrino properties will be crucial input information in the astrophysical context. For example, neutrinos are likely to become important as “astrophysical messengers” as the upcoming generation of km$^2$-scale high-energy neutrino telescopes begins to open a new window to the universe [4, 5]. The measured mixing parameters will be crucial for interpreting the neutrino signals.

In my lectures I will focus on two other themes that connect neutrinos with astrophysics. One is the role of neutrino masses in cosmology. The theory of structure formation in conjunction with recent and upcoming galaxy redshift surveys provide the most restrictive limits on neutrino masses (Sec. 2). Further, the anticipated laboratory confirmation of the large-mixing-angle oscillation solution of the solar neutrino problem in conjunction with big-bang nucleosynthesis will for the first time fix the number density of cosmic neutrinos and thus establish a unique connection between neutrino masses and the cosmic hot dark matter fraction (Sec. 3). Of course, massive neutrinos most likely have less to do with cosmic dark matter than with the baryonic matter by virtue of the leptogenesis mechanism for creating the baryon asymmetry of the universe (Sec. 4).

My second topic are supernova (SN) neutrinos. The current and upcoming large-scale neutrino detectors have a variety of primary goals, but usually double as SN observatories. The foreseeable experiments may cover the neutrino sky for several decades so that one may well expect an eventual SN observation. While the detection of about 20 neutrinos from SN 1987A was a milestone of neutrino astronomy, a high-statistics observation would allow one to extract a wealth of information. Section 5 is devoted to the mechanism of core-collapse SN explosions, the expected neutrino signature, and the relevant detectors while Sec. 6 is given over to neutrino masses and oscillations.

The atmospheric neutrino anomaly indicates $\nu_\mu \to \nu_\tau$ oscillations [6] with the mixing parameters shown in table I. The recent SNO results have largely established active-active oscillations as a solution of the solar neutrino problem [7, 8, 9]. The LMA parameters are strongly favored, but the LOW case may still be viable. There remains the unconfirmed evidence for flavor transformations from the LSND experiment [10]. If interpreted in terms of oscillations, the mixing parameters from all three sources are mu-
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Table I. – Experimental evidence for neutrino flavor oscillations.

| Evidence   | Channel       | $\Delta m^2$ [eV$^2$] | $\sin^2 2\Theta$ |
|------------|---------------|-----------------------|------------------|
| Atmospheric| $\nu_\mu \to \nu_\tau$ | $(1.6-3.9) \times 10^{-3}$ | 0.92-1          |
| Solar:     | LMA           | $\nu_e \to \nu_\mu\tau$ | $(0.2-2) \times 10^{-4}$ | 0.2-0.6       |
| Solar:     | LOW           | $\nu_e \to \nu_\mu\tau$ | $1.3 \times 10^{-7}$     | 0.92           |
|            | $\bar{\nu}_\mu \to \bar{\nu}_e$ | 0.2-10     | $(0.2-3) \times 10^{-2}$ |

tually inconsistent. Even including a putative sterile neutrino no longer provides a good
global fit because all three cases prefer active-active over active-sterile oscillations [11, 12].
MiniBooNE at Fermilab is expected to confirm or refute LSND within two years [13]. As
there is no straightforward global interpretation I will follow the widespread practice of
discarding LSND. If it is due to neutrino conversions after all, something fundamentally
new is going on and much of what follows may have to be revised. I will always assume
that there are three neutrino mass eigenstates separated by the atmospheric and solar
mass differences. One obvious open question is the value of the overall neutrino mass
scale $m_\nu$. What is cosmology’s contribution to clarifying this issue?

2. – Neutrino Dark Matter and Cosmic Structure Formation

The cosmic number density of neutrinos and anti-neutrinos per flavor is $n_{\nu\bar{\nu}} = \frac{3}{4} n_\gamma$
with $n_\gamma$ the number density of cosmic microwave photons, and assuming that a possible
asymmetry between the $\nu$ and $\bar{\nu}$ distributions is negligible [14]. With $T_\gamma = 2.728$ K this
translates into $n_{\nu\bar{\nu}} = 112$ cm$^{-3}$. If neutrinos have masses one finds a cosmic mass fraction

$$\Omega_{\nu\bar{\nu}} h^2 = \sum_{\text{flavors}} \frac{m_\nu}{92.5 \text{ eV}},$$

where $h$ is the Hubble constant in units of 100 km s$^{-1}$ Mpc$^{-1}$. The requirement that
neutrinos not overclose the universe implies the traditional limit $\sum m_\nu < 40$ eV, an
argument that was first advanced in a classic paper by Gershtein and Zeldovich [15].

Later Cowsik and McClelland [16] speculated that massive neutrinos could actually
provide the dark matter in clusters of galaxies. However, it soon became clear that
neutrinos were not a good dark matter candidate for two reasons. One is the limited
phase space for neutrinos gravitationally bound to a galaxy [17]. Therefore, galactic
dark-matter neutrinos must obey a lower mass limit of some 30 eV for typical spirals,
and even 100–200 eV for dwarf galaxies (“Tremaine-Gunn-limit”). Today the most re-
strictive laboratory limits on the overall neutrino mass scale come from the Mainz [18]
and Troitsk [19] tritium experiments. The current limit is [20]

$$m_\nu < 2.2 \text{ eV} \quad \text{at 95\% CL}.$$
It applies to all mass eigenstates if the mass differences are given by the atmospheric and solar oscillation interpretation. Equation (2) is so restrictive that neutrinos as galactic dark matter are completely ruled out.

Secondly, cosmic structure formation provides powerful information on the neutrino mass fraction $\Omega_{\nu \bar{\nu}}$. The observed structure in the distribution of galaxies is thought to arise from the gravitational instability of primordial density fluctuations. The small masses of neutrinos imply that they stay relativistic for a long time after their decoupling (“hot dark matter”), allowing them to stream freely, thereby erasing the primordial density fluctuations on small scales [14, 21]. While this effect does not preclude neutrino dark matter, it implies a top-down scenario for structure formation where large structures form first, later fragmenting into smaller ones. It was soon realized that the predicted properties of the large-scale matter distribution did not agree with observations [22].

Today it is widely accepted that the universe has critical density and that its mass-energy inventory sports several nontrivial components. Besides some 5% baryonic matter (most of it dark) there are some 25% cold dark matter in an unidentified physical form and some 70% of a negative-pressure component (“dark energy”). And because neutrinos do have mass, they contribute at least 0.1%. This fraction is based on $m_3 = 50$ meV, the smallest value consistent with atmospheric neutrino oscillations.

An upper limit on the neutrino dark matter fraction can be derived from the power spectrum $P_M(k)$ of the cosmic matter distribution. Neutrino free streaming suppresses the small-scale structure by an approximate amount [23]

$$\frac{\Delta P_M}{P_M} \approx -8 \frac{\Omega_{\nu}}{\Omega_M}$$

where $\Omega_M$ is the cosmic matter fraction, excluding the dark energy. This effect is illustrated in fig. 1 where $P_M(k)$ measured by the 2dF Galaxy Redshift Survey is compared with the predictions for a cold dark matter cosmology with neutrino fractions $\Omega_{\nu} = 0$, 0.01, and 0.05, respectively [24]. Based on the 2dFGRS data one finds $\sum m_\nu < 1.8–3.0$ eV, depending on the assumed priors for other cosmological parameters, notably the Hubble constant, the overall matter fraction $\Omega_M$, and the tilt of the spectrum of primordial density fluctuations [24, 25]. For a reasonable set of priors one may adopt

$$\sum_{\text{flavors}} m_\nu < 2.5 \text{ eV}$$

at a statistical confidence level of 95%. This limit corresponds approximately to the dot-dashed ($\Omega_{\nu} = 0.05$) curve in fig. 1. Neutrinos may still contribute as much as 5% of the critical density, about as much as baryons.

Within the standard theory of structure formation, the largest systematic uncertainty comes from the unknown biasing parameter $b$ that relates the power spectrum of the galaxy distribution to the underlying matter distribution, $P_{\text{Gal}}(k) = b^2 P_M(k)$. The biasing parameter is one of the quantities which must be taken into account when fitting all large-scale structure data to observations of the galaxy distribution and of the
temperature fluctuations of the cosmic microwave background radiation. In the future the Sloan Digital Sky Survey will have greater sensitivity to the overall shape of $P_{\text{Gal}}(k)$ on the relevant scales, allowing one to disentangle more reliably the impact of $b$ and $\Omega_\nu$ on $P_M(k)$. It is foreseen that one can then reach a sensitivity of $\sum m_\nu \sim 0.65 \text{ eV}$ [23].

For degenerate neutrino masses eq. (4) corresponds to a limit on the overall mass scale of $m_\nu < 0.8 \text{ eV}$, more restrictive than the laboratory limit eq. (2). However, the KATRIN project for improving the tritium endpoint sensitivity is foreseen to reach 0.3 eV [26], similar to the anticipated sensitivity of future cosmological observations. If both methods yield a positive signature, they will mutually re-enforce each other. If they both find upper limits, again they will be able to cross-check each other’s constraints.

3. – How Many Neutrinos in the Universe?

The laboratory limits or future measurements of $m_\nu$ and of a hot dark matter component can be related to each other if the cosmic neutrino density $n_{\nu\bar{\nu}}$ is known. However, the cosmic neutrinos can not be measured with foreseeable methods so that one depends on indirect arguments for determining $n_{\nu\bar{\nu}}$. Even if we accept that there are exactly three neutrino flavors as indicated by the $Z^0$ decay width and that they were once in thermal equilibrium does not fix $n_{\nu\bar{\nu}}$. Each flavor is characterized by an unknown chemical potential $\mu_\nu$ or a degeneracy parameter $\xi_\nu = \mu_\nu/T$, the latter being invariant under cosmic expansion. While the small cosmic baryon-to-photon ratio $\sim 10^{-9}$ suggests that the degeneracy parameters of all fermions are very small, for neutrinos this is an assumption and not an established fact.
In the presence of a degeneracy parameter $\xi_{\nu}$, the number and energy densities of relativistic neutrinos plus anti-neutrinos of one flavor are

$$n_{\nu\bar{\nu}} = T_{\nu}^3 \frac{3\zeta_3}{2\pi^2} \left[ 1 + \frac{2\ln(2)}{3\zeta_3} + \frac{\xi_{\nu}^4}{72\zeta_3} + O(\xi_{\nu}^6) \right],$$

(5)

$$\rho_{\nu\bar{\nu}} = T_{\nu}^4 \frac{7\pi^2}{120} \left[ 1 + \frac{30}{7} \left( \frac{\xi_{\nu}}{\pi} \right)^2 + \frac{15}{7} \left( \frac{\xi_{\nu}}{\pi} \right)^4 \right].$$

(6)

Therefore, if chemical potentials are taken to be the only uncertainty, $n_{\nu\bar{\nu}}$ can only be larger than the standard value. In this sense the structure-formation limits on the hot dark matter fraction provide a conservative limit on $m_{\nu}$. Conversely, a laboratory limit on $m_{\nu}$ does not limit the hot dark matter fraction.

Big-bang nucleosynthesis (BBN) is affected by $\rho_{\nu\bar{\nu}}$ in that a larger neutrino density increases the primordial expansion rate, thereby increasing the neutron-to-proton freeze-out ratio $n/p$ and thus the cosmic helium abundance. Therefore, the observed primordial helium abundance provides a limit on $\rho_{\nu\bar{\nu}}$ which corresponds to some fraction of one effective extra neutrino species. In addition, an electron neutrino chemical potential modifies $n/p \propto \exp(-\xi_{\nu_e})$. Depending on the sign of $\xi_{\nu_e}$ this effect can increase or decrease the helium abundance and can compensate for the effect of other flavors [27]. If $\xi_{\nu_e}$ is the only chemical potential, BBN provides the limit

$$-0.01 < \xi_{\nu_e} < 0.07.$$  

(7)

Including the compensation effect, the only upper limit on the radiation density comes from precision measurements of the power spectrum of the temperature fluctuations of the cosmic microwave background radiation and from large-scale structure measurements. A recent analysis yields the allowed regions [28]

$$-0.01 < \xi_{\nu_e} < 0.22, \quad |\xi_{\nu_{\mu,\tau}}| < 2.6,$$

(8)

in agreement with similar results of [29] and [30].

However, the neutrino oscillations imply that the individual flavor lepton numbers are not conserved and that in full thermal equilibrium all neutrinos are characterized by one single chemical potential $\xi_{\nu}$. If flavor equilibrium is achieved before $n/p$ freeze-out, the restrictive BBN limit on $\xi_{\nu_e}$ applies to all flavors, i.e. $|\xi_{\nu}| < 0.07$, implying that the cosmic number density of neutrinos is fixed to within about 1%. In that case the relation between $\Omega_{\nu}$ and $m_{\nu}$ is uniquely given by the standard expression eq. (1).

The approach to flavor equilibrium in the early universe by neutrino oscillations and collisions was recently studied [31, 32, 33, 34]. Assuming the atmospheric and solar LMA solutions, an example for the cosmic flavor evolution is shown in fig. 2. The detailed treatment is complicated and involves subtleties related to the large weak potential caused by the neutrinos themselves. The intriguing phenomenon of synchronized flavor
oscillations [35, 36] plays an interesting role, that coincidentally could also be important in the context of supernova neutrinos [37].

![Cosmological evolution of neutrino degeneracy parameters for the initial values ξνe = ξντ = 0 and ξνμ = −0.1. The mixing parameters were chosen according to the atmospheric and solar LMA solutions with Θ13 = 0. (Figure from Ref. [32] with permission.]

The practical bottom line, however, is rather simple. Effective flavor equilibrium before n/p freeze-out is reliably achieved if the solar oscillation parameters are in the favored LMA region. In the LOW region, the result depends on the value of the small but unknown third mixing angle Θ13. In the SMA and VAC regions, which are now heavily disfavored, equilibrium is not achieved. Establishing LMA as the correct solution of the solar neutrino problem, for example by the Kamland reactor experiment [38], amounts in our context to counting the number of cosmological neutrinos and thus to establishing a unique relationship between mν and Ων¯ν.

4. – Leptogenesis

Neutrino masses in the sub-eV range can play an interesting albeit indirect role for creating the baryon asymmetry of the universe (BAU) in the framework of leptogenesis scenarios [39]. The main ingredients are those of the usual see-saw mechanism for small neutrino masses. Restricting ourselves to a single family, the relevant parameters are the heavy Majorana mass M of the ordinary neutrino’s right-handed partner and a Yukawa coupling gν between the neutrinos and the Higgs field Φ. The observed neutrino then has a Majorana mass

\[ m_\nu = \frac{g_\nu^2 \langle \Phi \rangle^2}{M} \]

that can be very small if M is large, even if the Yukawa coupling gν is comparable to that for other fermions. Here, \( \langle \Phi \rangle \) is the vacuum expectation value of the Higgs field which also gives masses to the other fermions.
The heavy Majorana neutrinos will be in thermal equilibrium in the early universe. When the temperature falls below their mass, their density is Boltzmann suppressed. However, if at that time they are no longer in thermal equilibrium, their abundance will exceed the equilibrium distribution. The subsequent out-of-equilibrium decays can lead to the net generation of lepton number. CP-violating decays are possible by the usual interference of tree-level with one-loop diagrams with suitably adjusted phases of the various couplings. The generated lepton number excess will be re-processed by standard-model sphaleron effects which respect $B-L$ but violate $B+L$. It is straightforward to generate the observed BAU by this mechanism.

The requirement that the heavy Majorana neutrinos freeze out before they get Boltzmann suppressed implies an upper limit on the combination of parameters $g_\nu^2/M$ that also appears in the see-saw formula for $m_\nu$. The out-of-equilibrium condition thus implies an upper limit on $m_\nu$. Detailed scenarios for generic neutrino mass and mixing schemes have been worked out, see Ref. [40] for a recent review and citations of the large body of pertinent literature. The bottom line is that neutrino mass and mixing schemes suggested by the atmospheric and solar oscillation data are nicely consistent with plausible leptogenesis scenarios. Of course, it is an open question of how one would go about to verify or falsify leptogenesis as the correct baryogenesis mechanism. Still, it is intriguing that massive neutrinos may have a lot more to do with the baryons than with the dark matter of the universe!

5. – Core-Collapse Supernovae

5.1. Explosion Mechanism. – A core-collapse supernova (SN) is perhaps the only system besides the early universe that is dynamically dominated by neutrinos. Moreover, stellar collapse neutrinos have been observed once from SN 1987A and may be observed again from a future galactic SN. A core-collapse SN marks the evolutionary end of a massive star ($M \gtrsim 8 M_\odot$) that has reached the usual onion structure with several burning shells, an expanded envelope, and a degenerate iron core. The core mass grows by nuclear burning at its edge until it reaches the Chandrasekhar limit. The collapse can not ignite nuclear fusion because iron is the most tightly bound nucleus. Therefore, the collapse continues until the equation of state stiffens at about nuclear density ($3 \times 10^{14} \text{ g cm}^{-3}$). At this “bounce” a shock wave forms, moving outward and expelling the stellar mantle and envelope, i.e. the explosion is a reversed implosion. Within the expanding nebula, a compact object remains in the form of a neutron star or perhaps sometimes a black hole. The kinetic energy of the explosion carries about 1% of the liberated gravitational binding energy of about $3 \times 10^{53} \text{ erg}$, the remaining 99% going into neutrinos. The main phases of this sequence of events are illustrated in fig. 3.

The bounce-and-shock explosion mechanism [41, 42, 43, 44] is essentially a hydrodynamical phenomenon. However, realistic numerical simulations have difficulties exploding for a physical reason. The shock wave forms within the iron core. As it moves outward it dissipates energy by the dissociation of iron. The nuclear binding energy of 0.1 $M_\odot$ iron is about $1.7 \times 10^{51} \text{ erg}$ and thus comparable to the explosion energy. Therefore,
Fig. 3. – Schematic picture of the core collapse of a massive star \( (M \gtrsim 8 \, M_\odot) \) and the beginning of a SN explosion. There are four main phases: 1. Collapse. 2. Prompt-shock propagation and break-out, release of prompt \( \nu_e \) burst. 3. Matter accretion and mantle cooling. 4. Kelvin-Helmholtz cooling of proto-neutron star. The curves mark the evolution of several characteristic radii: The stellar iron core \( (R_{Fe}) \). The neutrino sphere \( (R_\nu) \) with diffusive transport inside and free streaming outside. The “inner core” \( (R_{ic}) \) for \( t \gtrsim 0.1 \) s is the region of subsonic collapse, later it is the settled, compact inner region of the nascent neutron star. The SN shock wave \( (R_{\text{shock}}) \) is formed at core bounce, stagnates for as much as several 100 ms, and then propagates outward. The shaded area is the neutrino source region.

The shock wave stalls without driving off the stellar mantle and envelope. The standard scenario holds that the stagnating shock is “re-juvenated” by energy deposition so that enough pressure builds up behind the shock to set it back into motion. This “delayed explosion scenario” was first proposed in the early 1980s by Bethe and Wilson [45]. One source of energy deposition behind the shock wave is energy absorption from the nearly freely streaming neutrinos which originate from the neutrino sphere near the neutron-star surface. Continued mass accretion and convection below the shock wave also deposit energy and thus contribute to the shock revival.

The main recent progress in numerical SN calculation has been the implementation of efficient Boltzmann solvers so that a realistic neutrino transport scheme can be self-consistently coupled with the hydrodynamical evolution [46, 47, 48]. Such state-of-the-art spherically symmetric calculations do not lead to successful explosions. (The Livermore group does obtain robust explosions [49]. In their spherically symmetric calculations they include a mixing-length treatment of “neutron finger convection,” thereby enhancing the early neutrino luminosity and thus the energy deposition behind the shock [50]. Their
results agree with the findings of other groups that diffusive neutrino transport alone is not enough to trigger the explosion.) The spherically symmetric calculations are not self-consistent in that the regions below the shock wave are convectively unstable. Forthcoming calculations will reveal if convection, perhaps coupled with more accurate neutrino interaction rates, will lead to successful explosions.

5.2. Expected Neutrino Signal. – The expected neutrino fluxes and spectra are illustrated by fig. 4. The $\nu_e$ lightcurve shows a conspicuous spike early on, representing the prompt neutrino burst which occurs when the shock wave reaches the region of neutrino trapping in the iron core. The dissociation of iron allows for the quick neutronization of a layer of the proto neutron star. Of course, most of the lepton number remains trapped and slowly escapes by neutrino diffusion. The subsequent broad shoulder up to about 300 ms, best visible in the right panel with linear scales, represents the accretion phase where material keeps falling in and powers the neutrino emission. After this phase the shock wave has driven off the stellar mantle. The subsequent long and flat tail represents the neutron star cooling by neutrino emission. The duration of the accretion phase depends on how long it takes to revive the shock wave. In the absence of a confirmed robust explosion mechanism the exact duration of the accretion phase is not known. Moreover, it may well depend on details of the star, e.g. on the progenitor mass.

In order to observe effects of neutrino oscillations, the fluxes and/or spectra must be different between the different flavors. In the simulation of fig. 4 the neutrino luminosity is virtually equipartitioned among the flavors after about 100 ms. In the recent Oakridge simulation [47], which includes a state-of-the-art Boltzmann solver, the equipartition is

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Fig. 4. – SN neutrino luminosities and average energies from a simulation with the Livermore code. The $\nu_e$ line represents each of $\nu_\mu$, $\bar{\nu}_\mu$, $\nu_\tau$ and $\bar{\nu}_\tau$. Left panel: Logarithmic luminosity and time scales. Right panel: Linear scales. (Figures from Ref. [49] with permission.)
also nearly perfect between \( \nu_e \) and \( \bar{\nu}_e \), but the \( \nu_x \) luminosity is less than half that after 50 ms out to 600 ms when this simulation terminates. Similar results are found in a Garching simulation [51]. Therefore, “equipartition” probably should be taken to mean “equal to within about a factor of two.” A recent systematic study of flavor-dependent neutrino transport also concludes that one can not expect the flux equipartition to be exact throughout the SN evolution [52].

The neutrino average energies obey the well-known hierarchy \( \langle E_{\nu_e} \rangle < \langle E_{\bar{\nu}_e} \rangle < \langle E_{\nu_x} \rangle \) which is explained by the different trapping processes, \( \beta \) processes for the electron flavor and elastic scattering on nucleons for the rest. Therefore, the different flavors originate in layers with different temperatures. A physical understanding of the neutrino spectra can be developed without large-scale numerical simulations [53]. While the flavor hierarchy of average energies appears to be generic, the differences are likely smaller than previously thought after all relevant processes have been included, notably nucleon bremsstrahlung and energy transfer by recoils [51, 52, 53, 54]. The average neutrino energies increase for the first few seconds. This is a generic effect because the neutrino-emitting regions heat up by accretion and by the contraction of the neutron star.

Numerical neutrino light curves can be compared with the SN 1987A data where the measured energies are found to be “too low.” For example, the numerical simulation of fig. 4 yields time-integrated values \( \langle E_{\nu_e} \rangle \approx 13 \) MeV, \( \langle E_{\bar{\nu}_e} \rangle \approx 16 \) MeV, and \( \langle E_{\nu_x} \rangle \approx 23 \) MeV. On the other hand, the data imply \( \langle E_{\bar{\nu}_e} \rangle = 7.5 \) MeV at Kamiokande and 11.1 MeV at IMB [55]. Even the 95% confidence range for Kamiokande implies \( \langle E_{\bar{\nu}_e} \rangle < 12 \) MeV. Flavor oscillations would increase the expected energies and thus enhance the discrepancy [55]. It has remained unclear if these and other anomalies of the SN 1987A neutrino signal should be blamed on small-number statistics, or point to a serious problem with the SN models or the detectors.

5.3. Possibility of a Future Detection. – A serious confrontation of SN theory with neutrino data has to await a high-statistics observation. Detectors for measuring the neutrinos from a galactic SN have almost continuously operated since 1980 when the Baksan Scintillator Telescope (BST) started. For a SN at a distance of 10 kpc with neutrino fluxes and spectra roughly like those of fig. 4, BST would register about 70 events. The neutrinos from SN 1987A in the Large Magellanic Cloud at a distance of 50 kpc were measured in Kamiokande [56], IMB [57], and BST [58] with a few events each. Today, much larger detectors exist, although BST keeps running. Super-Kamiokande would measure about 8000 events from a SN at 10 kpc. A simulated light curve based on the SN model of fig. 4 is shown in fig. 5. Of course, a SN observation by Super-Kamiokande depends on the reconstruction of the detector after its damage by the accident of 12 November 2001. Fortunately, there are other large observatories. The Sudbury Neutrino Observatory (SNO) would register about 800 events [59], ignoring flavor oscillations. The Large Volume Detector (LVD) in the Gran Sasso Laboratory is a scintillation detector that would register about 400 events [60]. A similar number of events would be expected in the KamLAND scintillation experiment which took up operations [61] and about 200 in MiniBooNE [62]. The Borexino solar neutrino experiment, that will soon be ready, is
smaller and would register about 100 events [63]. The AMANDA South Pole neutrino telescope also works as a SN neutrino detector in that the correlated noise of all photomultipliers caused by the Cherenkov light of the SN neutrinos produces a significant signal, especially when AMANDA is enlarged to the km$^3$ IceCube [64].

The dominant signal is usually the charged-current reaction $\bar{\nu}_e p \rightarrow n e^+$. SNO has a unique $\nu_e$ detection capability from the CC deuterium dissociation $\nu_e d \rightarrow p p e^-$. Neutral-current reactions which are sensitive to all flavors include elastic scattering on electrons, the deuterium dissociation $\nu d \rightarrow np\nu$ in SNO, the excitation of $^{16}$O in water Cherenkov detectors, and the corresponding excitation of $^{12}$C in scintillation detectors, notably in LVD and KamLAND, where the $\gamma$-rays from the subsequent de-excitation can be measured. Another possibility is proton elastic scattering that can cause a measurable scintillation signal [65]. Specific neutral-current detectors for SN neutrinos have been proposed on the basis of the reaction $\nu + (A, Z) \rightarrow (A-1, Z) + n + \nu$ where the neutron will be measured. For example, lead or iron could be used as targets in the proposed OMNIS detector [66, 67]. This sort of detector would be complementary to Super-Kamiokande and SNO in that it is primarily sensitive to the heavy-flavor neutrinos.

At present one debates the possibility of building even larger detectors for precision neutrino long-baseline oscillation experiments, proton decay, and high-statistics solar, atmospheric and SN neutrino detection. A typical size could be a megatonne of water or scintillator. This option is discussed under the name of Hyper-Kamiokande in Japan [68], under UNO in the US [69], and is also debated in Europe [70]. Such a detector could produce as many as $10^5$ events from our fiducial SN at 10 kpc.
The most crucial question is the SN rate in our galaxy because even the largest foreseen detectors will barely cover the local group. A SN in the largest member, Andromeda (M31), at a distance of about 750 kpc would yield only about 30 events in a megatonne. One approach to estimating the galactic SN rate is to use the average rates derived from external galaxies. The most recent estimate is \(2 \pm 1\) core-collapse SNe per century \([71]\), about a factor of 2 smaller than previous estimates \([72, 73]\). Another approach relies on the historical SN record, extrapolated to the entire galaxy. Because of obscuration by dust, only SNe out to a few kpc have been observed. The rate of core-collapse SNe is then estimated to be 3–4 per century \([73, 74]\), with a large Poisson uncertainty from the small number of observed cases (5 SNe during the past millennium). Given the vagaries of small-number statistics, these estimates agree with each other, and with circumstantial evidence such as the estimated population of progenitor stars or the neutron-star formation rate.

Large neutrino detectors are motivated by many goals so that it is not unrealistic to expect another few decades of coverage. Within that time one may well observe a galactic SN even though they are rare. A high-statistics neutrino light curve would allow one to observe directly the collapse dynamics. For example, the early accretion-powered neutrino emission could be clearly distinguished from the subsequent neutron-star cooling phase \([49]\). One of the most energetic astrophysical phenomena would be caught in the act, allowing one to unravel the underlying physics.

6. – Neutrino Masses and Oscillations

6.1. Time-of-Flight Dispersion of Neutrino Burst. – In principle, neutrino masses can be measured by the dispersion of a neutrino burst from a pulsed source, notably a SN. The time-of-flight delay of massive neutrinos with energy \(E_{\nu}\) is

\[
\Delta t = \frac{m_{\nu}^2}{2E_{\nu}^2} D
\]

where \(D\) is the distance to the source. Therefore, if a neutrino burst has the intrinsic duration \(\Delta t\) and the energies are broadly distributed around some typical energy \(E_{\nu}\), one is approximately sensitive to masses

\[
m_{\nu} > 10 \text{ eV} \left( \frac{E_{\nu}}{10 \text{ MeV}} \right) \left( \frac{\Delta t}{s} \right)^{1/2} \left( \frac{10 \text{ kpc}}{D} \right)^{1/2}.
\]

The measured \(\bar{\nu}_e\) burst of SN 1987A was characterized by \(E \approx 20 \text{ MeV}, \Delta t \approx 10 \text{ s},\) and \(D \approx 50 \text{ kpc},\) leading to the well-known limit \(m_{\nu} \lesssim 20 \text{ eV} [75].\) A recent re-analysis yields a somewhat more restrictive limit \([76]\). Either way, these results are only of historical interest because the tritium and cosmological limits are now much more restrictive.

A high-statistics observation of a future galactic SN would yield more restrictive limits because the relevant time scale \(\Delta t\) is the fast rise time of around 100 ms rather than the overall burst duration of several seconds. Therefore, one is sensitive to smaller masses than the SN 1987A burst, despite the shorter baseline. Detailed Monte-Carlo
simulations imply that Super-Kamiokande would be sensitive to about $m_\nu \gtrsim 3$ eV, almost independently of the exact distance [77]. (At a larger distance one gains baseline but loses statistics, two effects that virtually cancel for a given detector size.)

Conceivably this sensitivity could be improved if a gravitational wave signal could be detected preceding the neutrinos, signifying the instant of the stellar collapse and bounce [78, 79]. In this case one may be sensitive to about 1 eV. It is also conceivable that a SN collapses to a black hole some short time after the original collapse. In this case the neutrino signal would terminate within $\Delta t \lesssim 0.5$ ms, thereby defining a very short time scale. Super-Kamiokande would be sensitive to $m_\nu \gtrsim 1.8$ eV [80, 81]. With a megatonne detector one could measure SN neutrinos throughout the local group of galaxies. From Andromeda at a distance of 750 kpc one would get around 30 events. Using the overall signal duration for $\Delta t$ yields a sensitivity of a few eV.

The only conceivable time-of-flight technique that could probe the sub-eV range involves Gamma-Ray Bursts (GRBs) which have been speculated to be strong neutrino sources. If the neutrino emission shows time structure on the millisecond scale, and assuming a cosmological distance of 1 Gpc, one would be sensitive to neutrino masses $m_\nu \gtrsim 0.1$ eV $E_\nu$/GeV. Therefore, observing millisecond time structure in sub-GeV neutrinos from a GRB would be sensitive to the sub-eV mass scale [82, 83].

6.2. Flavor Oscillations. – Neutrino oscillations are now firmly established so that the SN neutrino fluxes and spectra in a detector can be very different from those emitted by the source. This is especially true if the favored solar LMA case applies. The relevant mass difference implies that matter effects are important in the SN, and also in the Earth if the neutrinos enter the detector “from below.” The large “solar” mixing angle $\theta_{12}$ implies that oscillations will be important in both the $\nu_e$ and $\bar{\nu}_e$ channel. In the LMA case it is unavoidable, for example, that oscillations influence the SN 1987A signal interpretation, and that the detectors saw different spectra due to different Earth-crossing paths [55, 84, 85, 86]. Taking oscillations into account will be crucial for extracting information about the source if a future galactic SN is observed [87, 88, 89].

Taking the numerical source model of Fig. 4, the time-integrated spectra at Super-Kamiokande and SNO are shown in Fig. 6 for different nadir angles which determines the Earth-crossing path. The oscillation parameters were chosen for the LMA case with $\Delta m^2_{12} = 2 \times 10^{-5}$ eV$^2$, $\Delta m^2_{13} = 3.2 \times 10^{-3}$ eV$^2$, $\sin^2 \theta_{12} = 0.87$, and $\sin^2 \theta_{23} = 1.0$. The unknown third mixing angle was chosen small as $\sin^2 \theta_{13} = 1.0 \times 10^{-6}$. Figure 6 illustrates that dramatic modifications of the spectra can be expected for certain cases.

It is difficult to anticipate everything about future data. If a galactic SN is observed, what we can learn about neutrino oscillations depends on the detectors operating at that time and their geographical location. It will also depend on the true source properties regarding flavor-dependent spectra and fluxes, and what is already known about the neutrino oscillation parameters as input information at that time. Many authors have studied these questions [90, 91, 92, 93, 94, 95, 96, 97, 98]. It appears that one may well distinguish between large and small values of the elusive $\theta_{13}$ and between normal or inverted mass hierarchies, or even accurately pin down $\Delta m^2_{12}$. An important caveat
is that these studies usually made unrealistic assumptions about the source spectra and fluxes [52]. More realistic cases remain to be investigated. Still, a SN neutrino observation would complement the upcoming efforts of precision determination of neutrino oscillation parameters in long-baseline experiments [99, 100, 101].

7. – Summary and Conclusions

The compelling detection of flavor oscillations in the solar and atmospheric neutrino data have triggered a new era in neutrino physics. In the laboratory one will proceed with precision experiments aimed at measuring the details of the mixing matrix. Future tritium decay experiments may well be able to probe the overall neutrino mass scale down to the 0.3 eV range, but if the absolute masses are smaller, it will be very difficult to measure them, and the overall mass scale may remain the most important unknown quantity in neutrino physics for a long time to come.

Cosmological large-scale structure data at present provide the most restrictive limit on neutrino masses of $\sum m_\nu < 2.5$ eV, corresponding to $m_\nu < 0.8$ eV in a degenerate mass scenario. A rigorous relationship between the cosmic hot dark matter fraction $\Omega_{\nu\bar{\nu}}$ and $m_\nu$ depends on the cosmic neutrino density $n_{\nu\bar{\nu}}$. If the solar LMA solution is correct, big-bang nucleosynthesis constrains $n_{\nu\bar{\nu}}$ without further assumptions about the neutrino chemical potentials. In the LMA case neutrinos reach de-facto flavor equilibrium before the epoch of weak-interaction freeze out.

While neutrinos provide only a small fraction of the cosmic mass density, the mass and mixing schemes suggested by the oscillation experiments are nicely consistent with leptogenesis scenarios for creating the cosmic baryon asymmetry. Therefore, massive neutrinos may be more closely related to the cosmic baryons than the dark matter.
Experimental neutrino physics is developing fast. It is foreseeable that a broad range of neutrino detectors will operate for a long time, perhaps decades, with a variety of goals from precision oscillation physics to proton decay. Typically such detectors are also sensitive to SN neutrinos so that one may expect the neutrino sky to be covered perhaps for decades. Even though galactic SNe are rare, the neutrinos from about a thousand SNe are on their way, and the expected time scale of neutrino coverage makes a future high-statistics SN observation a realistic possibility.

The existing neutrino mass limits are already so low that one can not expect improving them with a galactic SN observation. Turning this around, we already know so much about neutrinos that we can be sure that the observed SN neutrino light curve faithfully represents the behavior of the source, without modifications by dispersion effects. The same can not be said about the spectra which are inevitably modified by flavor oscillations. Depending on the details of the source spectra, detection channel, and location on Earth, flavor oscillations can have a large impact on the observations, and may even allow one to disentangle some of the neutrino mixing scenarios.

Arguably the most important benefit of a high-statistics measurement of stellar collapse would be observing the dynamics of a cataclysmic astrophysical event that could never be observed in any other way. Whether or not numerical SN simulations will soon converge on a theoretical standard model for the collapse and explosion mechanism, the importance of its independent verification or falsification by a detailed neutrino light curve can not be overstated. As we learn more about neutrinos, their future role is likely that of astrophysical messengers, showing us phenomena that otherwise would remain forever invisible.

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