Signatures of supernova neutrino oscillations

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After a brief review of our current understanding of neutrino flavor conversions inside a core collapse supernova, we analyze the signatures of these neutrino oscillations that can be observed at future large neutrino detectors. We examine the observability of model-independent signatures like the neutronization burst suppression, multiple spectral splits, earth matter effects, and shock wave effects. We also indicate some indirect oscillation signals, and comment on the effect of oscillations on supernova astrophysics. Finally we point out the features in the neutrino spectra that experiments should look for, even irrespective of their theoretical interpretations.

1 Introduction

Neutrinos emitted from a core collapse supernova and arriving at a detector on the earth undergo flavor conversions in three distinct regions. Inside the star, the collective effects due to neutrino-neutrino interactions [1, 2] and the MSW matter effects due to neutrino-matter interactions [3] drive the flavor transformations. Between the star and the earth, the neutrino mass eigenstates travel independently so that there are no effective flavor conversions. If the neutrinos have to pass through the earth before reaching the detector, further neutrino oscillations due to the MSW matter effects take place inside the earth.

Our understanding of the neutrino flavor conversions inside the star has undergone significant changes in the last decade, and some gaps are yet to be comprehensively filled in. The analyses around the turn of the century were carried out under the assumption that the flavor conversions mainly take place in the MSW resonance regions H and L in the mantle, around densities of $\rho_H \sim 10^{3-4}$ g/cc and $\rho_L = 1 - 10$ g/cc, respectively. The neutrino spectra arriving at the earth are then sensitive to the neutrino mass hierarchy and to the value of θ_{13}, for $\sin^2\theta_{13}$ as low as 10^{-5} [4, 5]. The flavor conversion probabilities are independent of the primary fluxes in this scenario.

However the neutrino-neutrino forward scattering interactions just outside the neutrinosphere, where $\rho \sim 10^{6-10}$ g/cc, can trigger self-induced flavor conversions [6] and give rise to significant flavor transformations [7]. These *collective effects* manifest themselves in the form of qualitatively new phenomena like synchronized oscillations [8], bipolar/pendular oscillations [9], and spectral splits [10, 11]. These collective flavor conversions are possible even for $\sin^2\theta_{13}$ as low as 10^{-10} or even lower, since the pendular oscillations can be triggered by even a small instability [12]. However the neutrino flavor conversion probabilities now depend strongly on the primary neutrino fluxes. Initial investigations into the collective effects suggested that, while these collective oscillations would be virtually ineffective for normal hierarchy (NH), in the inverted hierarchy (IH) they would result in the complete swapping of $\bar{\nu}_e$ and $\bar{\nu}_\mu$ spectra. In addition, the ν_e and ν_μ spectra would be completely swapped for $E > E_c$ and unaffected for $E < E_c$ for a critical energy E_c [10]. The sharp change in the spectrum at $E = E_c$ is the spectral split.

It was later realized [13] that the phenomenon of a single spectral split at $E = E_c$ is a valid outcome only under special circumstances, for example, when $L_{\nu_e} \approx L_{\bar{\nu}_e} \gtrsim L_{\nu_\mu}$. In the general case, multiple spectral splits would take place, i.e. both $\nu_e \leftrightarrow \nu_y$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_y$ swaps ($\nu_y = \cos\theta_{23}\nu_\mu + \sin\theta_{23}\nu_\tau$) occur, in sharply separated energy regions. In addition, three-flavor effects [14, 15] tend to give rise to even $\nu_e \leftrightarrow \nu_x$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_x$ swaps ($\nu_x = -\sin\theta_{23}\nu_\mu + \cos\theta_{23}\nu_\tau$). The swapped and unswapped energy regions depend on primary fluxes and mass hierarchy. Combined with the MSW effects, these collective effects can give rise to many distinctive features in the ν_e and $\bar{\nu}_e$ spectra at the detector [16].

Most of the initial results with the inclusion of collective effects, both analytical and numerical, were obtained using the so-called single-angle approximation. With more detailed numerical simulations, it has become apparent that the *multi-angle* effects [7] can also have an impact: they can smoothen the flavor conversion features [17] and may even suppress the flavor conversions themselves [18, 19, 20, 21, 22]. The collective effects due to neutrino-matter interactions also come into play when multi-angle effects are included. The analytical understanding of these multi-angle effects is still a work in progress.

The aim of this talk is to analyze possible signatures of supernova neutrino oscillations at future large neutrino detectors, The first step is to determine what are the observables to look for. As we shall see, these observables can be identified with the knowledge of some broad features of the collective as well as MSW effects on neutrino oscillations. The feasibility of the relevant observations, and their interpretation in terms of neutrino mixing parameters, is where the detailed understanding of the neutrino flavor conversions becomes crucial. We shall start with a review of these flavor conversions in Sec. 2, and examine the relevant observables at the detectors in Sec. 3. We shall conclude in Sec. 4 with a summary and an outlook towards future.

2 Flavor conversions of supernova neutrinos

The simulations of supernova explosions still give rather varied predictions for the primary neutrino fluxes [23]. While they all agree on approximately thermal flavor spectra, and on the hierarchy of average energies $\langle E^0_{\nu_e}\rangle < \langle E^0_{\bar{\nu}_e}\rangle < \langle E^0_{\nu_\mu,\nu_\tau}\rangle \approx \langle E^0_{\bar{\nu}_\mu,\bar{\nu}_\tau}\rangle$, they differ in the actual values of these average energies. One typically has $\langle E^0_{\nu_e}\rangle \approx 10 - 12$ MeV, $\langle E^0_{\bar{\nu}_e}\rangle \approx 12 - 15$ MeV and $\langle E^0_{\nu_\mu}\rangle \approx 15 - 20$ MeV. These values also depend on the mass of the progenitor star. The relative luminosities of the flavors are also uncertain; though all the simulations agree on equal luminosities $L_{\nu_\mu}, L_{\nu_\tau}, L_{\bar{\nu}_\mu}$ and $L_{\bar{\nu}_\tau}$, and approximately equal luminosities L_{ν_e} and $L_{\bar{\nu}_e}$. The total energy released in neutrinos is $\sim 10^{53}$ erg. While discussing supernova neutrinos, it is convenient to talk in terms of the *flavors* ν_x and ν_y mentioned above. In the limit $\theta_{13} \to 0$, the states ν_x and ν_y are also mass eigenstates. Clearly, $\langle E^0_{\nu_x}\rangle \approx \langle E^0_{\nu_y}\rangle \approx \langle E^0_{\bar{\nu}_x}\rangle \approx \langle E^0_{\bar{\nu}_y}\rangle$ and $L_{\nu_x} \approx L_{\nu_y} \approx L_{\bar{\nu}_x} \approx L_{\bar{\nu}_y}$.

A neutrino ensemble may be described in the language of the occupation number matrices $\varrho(\vec{p})$ in the flavor basis [24, 2]. Its evolution may be described in terms of the effective

Hamiltonian $H(\vec{p}) = H_{vac}(\vec{p}) + H_{\rm MSW} + H_{\nu\nu}(\vec{p})$, where

$$\begin{aligned} H_{vac}(\vec{p}) &= M^2/(2p) , \\ H_{\rm MSW} &= \sqrt{2} G_F n_{e^-} diag(1,0,0) , \\ H_{\nu\nu}(\vec{p}) &= \sqrt{2} G_F \int \frac{d^3q}{(2\pi)^3} (1-\cos\theta_{pq}) \big(\varrho(\vec{q}) - \bar{\varrho}(\vec{q})\big) . \end{aligned} \quad (1)$$

Here θ_{pq} is the angle between the vectors $\vec{p}$ and $\vec{q}$. The first term $H_{vac}(\vec{p})$ arises from the neutrino mixing in vacuum, the second term $H_{\rm MSW}$ from the neutrino-matter interactions, and the last term $H_{\nu\nu}(\vec{p})$ from the neutrino-neutrino interactions. The equations of motion are

$$i \frac{\partial}{\partial t} \varrho(\vec{p}) = [H(\vec{p}), \varrho(\vec{p})] . \quad (2)$$

Note that $H(\vec{p})$ and $\varrho(\vec{p})$ are 3×3 matrices. The term $H_{\nu\nu}$ depends on ϱ itself, and hence makes the equations of motion nonlinear. In addition, the integration over $\vec{q}$ is rather complicated (and numerically time-consuming) due to the presence of $\varrho(\vec{q})$ terms in the integrand.

A simplifying assumption often used is the so-called *single-angle* approximation, where all the neutrinos at a given location are taken to be subject to the same average $\nu\nu$ potential, irrespective of their momentum $\vec{p}$. This is equivalent to an effective averaging of the factor of $(1-\cos\theta_{pq})$. As we shall see below, this approximation is enough to bring out many qualitative features of the evolution of the neutrino ensemble, however for a complete understanding of the flavor conversions, the complete multi-angle treatment is essential.

2.1 Oscillations due to collective effects

2.1.1 Collective effects with single-angle approximation

With the single-angle approximation, the neutrino-neutrino term in the Hamiltonian is dominant just outside the neutrinosphere, where the neutrinos have started free-streaming. The only effect of the neutrino-matter term in this region is the suppression of the effective mixing angle θ_{13}. In an iron-core supernova, the collective phenomena of synchronized oscillations, bipolar/pendular oscillations and spectral splits occur sequentially [17], followed by the MSW flavor conversions that occur mainly in the resonance layers in the mantle. The suppressed θ_{13} in matter implies that the flavor conversions are extremely small in the synchronization phase. However even a small θ_{13} is enough to cause a nonlinear instability in certain situations and start significant flavor conversions. This culminates in the formation of one or more spectral splits [10, 13].

The three-flavor effects can be roughly factorized into two-flavor effects that take place in a sequential manner [25]. The $\nu_e \leftrightarrow \nu_y$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_y$ pendular oscillations and spectral swaps are complete first, while the $\nu_e \leftrightarrow \nu_x$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_x$ pendular oscillations and spectral swaps occur later. These later swaps [14] (sometimes referred to as the *solar* swaps) are more likely to be incomplete or non-adiabatic, however they can sometimes effectively reverse the earlier $\nu_e \leftrightarrow \nu_y$ swaps, at least partially [15].

The net effect of collective oscillations is then a series of alternate swapped and unswapped regions in the ν_e and $\bar{\nu}_e$ spectra. The swaps may be incomplete in some cases. The locations and widths of the swapped/unswapped regions depend on the primary spectra. Assuming the equality of the luminosities of primary ν_e and $\bar{\nu}_e$ spectra (observed in most of the simulations), the numerical investigation of the pattern of these swapped regions leads to the following observations [26]: When the electron flavor dominates in the primary fluxes, i.e. $L_{\nu_e} > L_{\nu_\mu}$, one obtains (i) no spectral split for NH, and (ii) single spectral splits arising from $\nu_e \leftrightarrow \nu_y$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_y$ swaps in the neutrino and antineutrino channels, respectively, for IH. On the other hand, when the non-electron flavors dominate, i.e. $L_{\nu_e} < L_{\nu_\mu}$, one gets (i) single spectral splits arising from $\nu_e \leftrightarrow \nu_y$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_y$ swaps in the neutrino and antineutrino channels, respectively, for NH, and (ii) up to two spectral splits, both in the neutrino and antineutrino channels: arising from $\nu_e \leftrightarrow \nu_y, \nu_e \leftrightarrow \nu_x$ and $\bar{\nu}_e \leftrightarrow \bar{\nu}_y, \bar{\nu}_e \leftrightarrow \bar{\nu}_x$ swaps, respectively, for IH. The incompleteness of some of the swaps may lead to a possible energy dependence in the swapped energy regime.

2.1.2 Multi-angle effects on collective oscillations

Though most of the qualitative features of the collective effects may be obtained with a single-angle approximation, the numerical multi-angle simulations have indicated that the multi-angle effects can be significant in certain situations. Typically, multi-angle effects smear the sharp features in the spectra [17]. Large neutrino densities but low $\nu - \bar{\nu}$ asymmetry may give rise to additional instabilities that would have been absent with the single-angle approximation, further leading to multi-angle decoherence [18]. Very high neutrino densities also tend to delay the onset of pendular oscillation [19]. Moreover, with multi-angle effects included, the role of matter density is not restricted to the suppression of effective θ_{13}. Indeed, large matter densities, as may be possible in the accretion phase, also tend to cause multi-angle decoherence [20]. It is possible that such large densities may lead to a complete suppression of collective oscillations deep inside the core during the accretion phase [21]. The multi-angle effects thus change the picture of when the collective oscillations start and how they develop. Recent multi-angle numerical simulations [19] seem to suggest that while the onset of large oscillations may be delayed by large multi-angle effects of the neutrino or matter background, the final spectra look like smeared versions of the single-angle predictions.

The task of analytically understanding all the features of multi-angle effects seems intractable, primarily because of the nonlinear nature of the differential equations describing the neutrino evolution. However the observation that the equations of motion may be linearized at the onset of pendular oscillations allows one to analytically examine the conditions for the onset, even with the inclusion of multi-angle effects. Such an examination for azimuthally symmetric neutrino emission reveals the following interesting features [22], which are helpful in understanding some of the numerical observations above: (i) The neutrino background potential μ and matter background potential λ appear through the combination $\bar{\lambda} = \lambda + \epsilon\mu$, where ϵ is the fractional lepton number asymmetry. (ii) When $\mu \gg \bar{\lambda}$ or $\bar{\lambda} \gg \mu$, pendular oscillations cannot start. Indeed, the instability that would start significant oscillations cannot form unless the matter potential and neutrino potential are similar in magnitude. Nontrivial angular distributions may, however, give rise to additional instabilities [27].

Most of the work so far in understanding of multi-angle effects has been numerical and exploratory, and the jury is still out on the extent of these effects. The net effects of collective oscillations cannot be directly observed as neutrinos subsequently pass through MSW resonance regions, where they may undergo further flavor conversions.

2.2 Oscillations due to the MSW effect

After the neutrinos exit the region where the collective oscillations occur, further flavor conversions occur mainly in the MSW resonance regions [4, 5]. Here the conversion probabilities are independent of the spectra themselves, and are well understood in terms of the neutrino mixing parameters and density profiles. In particular, the flavor conversion in the H resonance is completely adiabatic (non-adiabatic) for $\sin^2\theta \gtrsim 10^{-3}$ ($\lesssim 10^{-5}$), while the L resonance is always completely adiabatic [5]. The neutrino fluxes exiting the star is an ensemble of decoherent neutrino mass eigenstates in vacuum, so the flavor combination during the propagation between the star and the earth is unchanged. The neutrino fluxes F of ν_e and $\bar{\nu}_e$ arriving at the earth may be written in terms of the primary fluxes F^0 and the survival probabilities p and $\bar{p}$ of ν_e and $\bar{\nu}_e$, respectively:

$$F_{\nu_e} = pF^0_{\nu_e} + (1-p)F^0_{\nu_x} \ , \qquad F_{\bar{\nu}_e} = \bar{p}F^0_{\bar{\nu}_e} + (1-\bar{p})F^0_{\nu_x} \ . \tag{3}$$

Though p and $\bar{p}$ are in general functions of energy, they are approximately constant with energy for *small* θ_{13} ($\sin^2\theta_{13} \lesssim 10^{-5}$) and *large* θ_{13} ($\sin^2\theta_{13} \gtrsim 10^{-3}$). At intermediate θ_{13} values, the energy dependence is more complicated, however we shall not consider such a situation here.

The value of p can be directly related to the neutrino mixing pattern during the ~ 10 ms *neutronization* burst of ν_e that occurs immediately after the core bounce. During the later accretion and cooling phases, unless the primary fluxes have widely different energies, it is virtually impossible to determine p or $\bar{p}$ given only the final ν_e and $\bar{\nu}_e$ spectra. However it may be possible to distinguish between zero and nonzero values of p or $\bar{p}$ through earth matter effects. Another phenomenon that allows us to decipher p or $\bar{p}$ values is the time variation in these quantities when the shock wave passes through the MSW resonance regions. Both the earth matter effects and shock wave effects are instances of the neutrino-matter interactions affecting neutrino survival probabilities. Below we review their essential features.

2.2.1 Earth matter effects

If F_{ν_1} and F_{ν_2} are the fluxes of ν_1 and ν_2 arriving at the earth, and the net ν_e flux after the neutrinos have travelled a distance L through the earth matter is $F^D_{\nu_e}(L)$, then [5]

$$F^D_{\nu_e}(L) - F^D_{\nu_e}(0) = (F_{\nu_2} - F_{\nu_1}) \sin 2\theta^{\oplus}_{12} \sin(2\theta^{\oplus}_{12} - 2\theta_{12}) \sin^2\left(\frac{\Delta m^2_{\oplus} L}{4E}\right) \ , \tag{4}$$

where $(\Delta m^2_{\oplus}, \theta^{\oplus}_{12})$ are the values of the solar Δm^2 and mixing angle in earth matter. For antineutrinos, the right hand side changes sign.

The vanishing of neutrino survival probability p corresponds to $F_{\nu_1} = F_{\nu_2}$ and similarly $\bar{p} = 0$ corresponds to $F_{\bar{\nu}_1} = F_{\bar{\nu}_2}$. Therefore, nonzero earth matter effects require $p \neq 0$ for neutrinos and $\bar{p} \neq 0$ for antineutrinos.

2.2.2 Shock wave effects

When the shock wave passes through the MSW resonance regions, the sharp density fluctuations in the shock wave may cause the adiabatic resonances to become non-adiabatic [29]. When the shock wave is at density ρ, it affects the neutrinos with the resonant energy

$$E_{\rm res} \approx 25 \text{ MeV} \frac{600}{Y_e \rho \text{ (g/cc)}} \ . \tag{5}$$

The resonant energies increase with time, and hence the nonadiabatic regions shift to higher energies with time [30]. This will result in a time-dependent value for p and $\bar{p}$ in NH and IH, respectively, during the time of propagation of the shock wave through the resonance region $\rho_H \sim 10^{3-4}$ g/cc, i.e. around 4-5 seconds after the core collapse.

The turbulence that follows the shock wave may, if large enough, cause flavor depolarization. In the extreme case, when complete three-flavor deleptonization occurs, the fluxes of all the neutrino species – or all the antineutrino species, depending on the hierarchy – become identical [32, 33]. For low turbulence amplitude and large θ_{13}, the features of the shock effect may survive [33]. Since the extent of turbulence created during the supernova explosion is still largely uncertain, it is not possible to make a concrete statement about about the net effects of turbulence at this point of time.

3 Neutrino signals at detectors

In this section, we shall point out the features of neutrino spectra at the detectors that will act as signatures of the neutrino oscillations. We shall further comment on the feasibility of robust identification of these signatures, and what we can learn about the neutrino mixing pattern as well as the dynamics of supernova explosion.

We shall consider three main categories of large neutrino detectors: water Cherenkov, carbon-based scintillators, and liquid argon detectors. The major interaction in the first two detectors is the inverse beta decay

$$\bar{\nu}_e + p \to e^+ + n ,$$

which helps reconstruct the $\bar{\nu}_e$ spectrum. While the energy resolution of the water Cherenkov detectors is typically a factor of 5-10 worse than that of the liquid scintillators [34], it is easier to make larger water Cherenkov detectors, so they typically have the advantage of larger statistics. The liquid argon detector is the best detector for observing the ν_e spectrum [35], the corresponding charged current (CC) reaction being

$$\nu_e + ^{40}\text{Ar} \to ^{40}\text{K}^* + e^- .$$

The rule-of-thumb estimate for the number of events observed through the above reactions is ~ 300 per kt in the 10 s duration of the neutrino pulse, for a supernova at 10 kpc. The neutral current (NC) interaction

$$\nu + p \to \nu + p$$

can also be identified through the small proton recoil [36], which can be measured at scintillation detectors. There are also sub-leading interactions like

- the forward scattering $\nu + e^- \to \nu + e^-$ that occurs in all the above detectors,
- $\nu_e + ^{16}\text{O} \to X + e^-$ in water Cherenkov, and
- the NC reaction $\nu + ^{12}\text{C} \to \nu + X + \gamma(15.11\ \text{MeV})$ in scintillator detectors,

which will not be discussed here.

The charged and neutral current interactions of neutrinos with a heavy nucleus like lead can release one or two neutrons free from the nuclei, through the reactions

- CC: $\nu_e +^{208}\text{Pb} \rightarrow^{207}\text{Bi} + n + e^-$, $\quad \nu_e +^{208}\text{Pb} \rightarrow^{206}\text{Bi} + 2n + e^-$, and
- NC: $\nu +^{208}\text{Pb} \rightarrow^{207}\text{Pb} + n$, $\quad \nu +^{208}\text{Pb} \rightarrow^{206}\text{Pb} + 2n$.

The threshold for the emission of two neutrons is higher than that for single neutron emission, thus allowing some energy discrimination on a statistical basis. Though no large detector of this type is under consideration, the HALO detector [37] has started operating already at the SNOLAB.

We shall focus on the leading charged current reactions at the large detectors above, which enable the reconstruction of the ν_e and $\bar{\nu}_e$ spectra. Following are some of the features of these spectra that can act as smoking gun signals of specific neutrino mixing scenarios.

3.1 Neutronization burst

The ~ 10 ms burst of ν_e that occurs immediately after the core bounce has a well-predicted flux [38] that is relatively free of model uncertainties. If the hierarchy is normal and $\sin^2\theta_{13} \gtrsim 10^{-3}$, the survival probability $p \approx \sin^2\theta_{13}$, so the burst signal is suppressed by a factor of $\sin^2\theta_{13}$ [5]. Such an extreme suppression (almost vanishing) of the neutronization burst signal is therefore a clear signature of this mixing scenario. However the robust identification of this signal needs a liquid Ar detector with sufficient time resolution to be able to separate the neutronization burst signal from the accretion phase fluxes that follow it.

A unique situation occurs if the progenitor of the supernova is not an iron-core star, but a O-Ne-Mg one. In such stars, the MSW resonances occur deep inside the region where the neutrinos are still undergoing collective oscillations [39]. Neutrinos of all energies then undergo the MSW resonance with the same adiabaticity [40]. In this case, the MSW resonance helps in preparing the neutrino ensemble for the spectral split. For NH this results in a single spectral split, while for IH this results in two sequential spectral splits [41]. The positions of these splits can be determined from the initial spectra and the non-adiabaticities at the resonances [42]. Recent multi-angle simulations indicate that the multi-angle effects do not change the results significantly [43]. The distortion of the neutrino spectra during the neutronization burst of a O-Ne-Mg supernova is thus unique, can identify the hierarchy even at extremely small θ_{13} values, and could be instrumental in identifying a supernova with such a light progenitor, in case an optical observation is not possible. Some more intriguing features of the neutronization burst phase of such a supernova have recently been reported in the numerical simulations [44].

3.2 Spectral split and Earth matter effects

Though the survival probability of ν_e or $\bar{\nu}_e$ changes sharply at the spectral splits, the observed signal is often diluted by the small difference between the swapping spectra. Moreover if the split is at lower energies, the smaller cross sections make the detection of the spectral split difficult. However if the primary fluxes are dominated by non-electron flavors, the splits can be at higher energies and may manifest themselves as shoulders in the ν_e or $\bar{\nu}_e$ spectra [26].

Earth matter effects provide a more practical way of determining a nonzero value of p or $\bar{p}$, since they introduce modulations of known frequency in the spectrum. Time-dependent changes in relative luminosities observed at two detectors, only one of which is shadowed by the earth, are indicators of earth matter effects [45]. On the other hand, the modulations in the ν_e or $\bar{\nu}_e$ spectra allow one to detect Earth matter effects even at a single detector [46]. While the former method needs two detectors with large fiducial masses (e.g. megaton water

Cherenkov, IceCube), the latter method needs detectors with a good energy resolution (e.g. liquid scintillator or liquid Ar).

A few remarks are in order. Typically, earth effects will be present only in a part of the spectrum due to the presence of spectral swaps. However this feature may be hard to observe. Multi-angle decoherence, turbulent effects, or small differences in primary spectra may result in the earth effects being unobservably small, so nothing may be inferred from their non-observation. However a positive identification of Earth effects would be enough to shortlist neutrino mixing patterns [26]. Relating the presence of Earth effects to the specific neutrino mixing scenarios needs a more complete understanding of the collective effects that we have at the moment.

3.3 Shock wave effects

The propagation of the shock wave through the MSW resonance region inside the star can give rise to time-dependent changes in p for NH, and in $\bar{p}$ for IH [29]. This would result in time-dependent dip/peak features in observables like $N_{\nu_e,\bar{\nu}_e}(E)$, $\langle E_{\nu_e,\bar{\nu}_e}\rangle$, $\langle E^2_{\nu_e,\bar{\nu}_e}\rangle$, etc.. Sharp changes in these observables in ν_e ($\bar{\nu}_e$) spectra at $t \gtrsim 3-4$ s testify for NH (IH) and $\sin^2\theta_{13} \gtrsim 10^{-5}$ [30, 31]. This may even allow the tracking of shock wave while it is still inside the mantle [30]. Note that probing the propagation of the shock wave at such early stages – before it breaks up the envelope of the star – is not possible through any other means (apart from possibly gravitational waves, but their detection is even harder).

If the multi-angle effects cause decoherence at such late times, or if the turbulence that follows the shock wave is large enough to cause flavor depolarization [32, 33], the spectra of all flavors may become identical and no shock effects will be observed. Thus, the non-observation of shock effects does not convey any concrete information. However, a positive observation of these effects can pinpoint the neutrino mixing pattern.

3.4 Indirect oscillation signals

So far we have focused on the oscillation signals through charge-current interactions, which are the primary reactions that allow us to reconstruct the ν_e and $\bar{\nu}_e$ spectra. However even more detailed information can be gained from complementary observables. Let us indicate some such observables in this subsection.

The detection of low energy protons recoiled from the $\nu p \to \nu p$ interaction is possible at a scintillator detector, with a threshold of ~ 0.2 MeV. The recoil spectrum of protons above this threshold can be reliably reconstructed with the superior energy resolution of such a detector. This would allow us to reconstruct the high energy tail of the sum of fluxes of all neutrinos. Clearly this will convey no direct information on neutrino oscillations. However, since the primary fluxes of non-electron neutrinos will be the major contributors to this tail, a fit to this tail would allow a measurement of the average energy $\langle E^0_{\nu_x}\rangle$ to a good accuracy [36]. This would allow a better interpretation of the primary signal observed through the charged-current interactions.

A QCD phase transition may take place in the core of the star, a few tens of a second after the core collapse. This would cause a sudden compactification of the progenitor core. At water Cherenkov detectors like Super-Kamiokande or IceCube, this will result in a prominent burst of $\bar{\nu}_e$ [47]. If a black hole is formed during the neutrino emission process, the neutrino signal will suddenly cease. Though not directly relevant to neutrino oscillations, these neutrino signals will provide information about supernova astrophysics that is not possible through any other means.

The neutrinos emitted from all the supernova exploded in the universe till now form the diffused supernova neutrino background (DSNB). The measurement of this background should be possible within a few years at large water Cherenkov or scintillation detectors. Such a measurement would test the predictions from astrophysics and cosmology. Collective oscillations affect predictions of the DSNB fluxes by up to $\sim 50\%$ [48], and the shock wave effects can further change these predictions by $10 - 20\%$ [49]. Thus, neutrino oscillations inside the star strongly influence the predictions of the DSNB flux.

If significant neutrino oscillations take place deep inside the core, they can also affect the abundances of heavy elements in the ejecta of supernovae. The r-process nucleosynthesis that is responsible for the production of heavy elements is influenced by the densities of ν_e and $\bar{\nu}_e$ in the relevant region, since these two species take part directly in the nucleosynthesis process. In the absence of collective effects, neutrino flavor conversions occur only in the resonance layers that are out in the mantle and hence the r-process cannot be affected. With the collective effects, however, there is a possibility of neutrino oscillations in a region deeper than the r-process region. These oscillations would tend to increase the average ν_e energy, thus the ν_e cross section with nuclei, suppressing the production of heavy elements. Oscillations are thus in general detrimental to successful r-process nucleosynthesis. However the exact amount of suppression depends strongly on astrophysical conditions and no concrete predictions can be made at this stage [50].

The shock wave propagation can also be affected by neutrino oscillations if they take place deep inside the core, a possibility opened up by the collective effects. Recent explorations into this question [51] indicate no significant impact on the explosion mechanism, however this is still work in progress.

4 Concluding remarks

The neutrino signal from the explosion of a core collapse supernova carries information on primary neutrino fluxes, neutrino mixing parameters, and the shock wave propagation. This information may be extracted by various complementary probes like the neutronization burst, earth matter effects, and shock wave effects. The vanishing of neutronization burst serves for a robust determination of NH and large θ_{13}, however it needs a liquid Ar detector with a good time resolution. The spectral splits are rather difficult to identify, however the identification of earth matter effects, which manifest themselves in terms of spectral modulations, vouches for a nonzero value for the survival probabilities p and $\bar{p}$. Interpreting p and $\bar{p}$ in terms of the neutrino mixing parameters needs a better understanding of multi-angle collective effects than we have at present. The shock wave effects, that result in time-dependent sharp changes in the spectra, are independent of collective effects and can identify the hierarchy, as long as θ_{13} is not too small and turbulent convections behind the shock wave do not give rise to complete depolarization. While the charged current events form the primary signal that helps us reconstruct the ν_e and $\bar{\nu}_e$ spectra, the proton recoil signal from the neutral current events aids the reconstruction of the primary flux of non-electron neutrinos.

With the help of the above signals, one can hope to solve the *inverse supernova neutrino problem*, which consists of observing (i) the $\nu_e/\bar{\nu}_e$ spectra (ii) NC events, (iii) time variation of the signal, and (iv) earth matter effects, and drawing conclusions about (i) the primary fluxes (ii) the shock propagation and (iii) the neutrino mixing parameters, especially the mass hierarchy. The task is not impossible, but there are many gaps yet to be filled. The major source of the gaps are the uncertainties in primary fluxes, which prevent us from a good reconstruction of the survival probabilities p and $\bar{p}$ as a function of energy, and our incomplete understanding of flavor oscillations inside the star. The details of collective oscillations including the multi-angle effects and the extent of turbulence are two issues that still remain to be resolved.

Nevertheless, it is crucial to make the following measurements: (i) reconstruction of ν_e and $\bar{\nu}_e$ spectra through CC events, (ii) NC spectra through proton recoil at scintillation detectors, (iii) single- and double-neutron events at Pb detectors, (iv) time modulation of observables like flux, average energy, and higher moments, (v) time dependent ratios of relative luminosities at large detectors, (vi) oscillatory spectral modulations from earth effects, (vii) other non-thermal features in the spectrum. Detectors should focus on the above measurements irrespective of the theoretical motivation or interpretation available presently. This has two reasons. First, a core-collapse supernova explosion in our galaxy is a rare enough event that when it happens once, the opportunity to extract whatever data from it should not be missed. Second, the history of neutrinos is full of surprises in the data that the theory had not anticipated at all.

The recent indications of a large θ_{13} [52, 53] imply some interesting consequences for the supernova neutrino analysis. In such a case, the H resonance is adiabatic, except possibly when the shock wave is propagating through the resonance region. The shock wave effects would then be prominent, the hierarchy determination easier, and shock tracking more feasible. Moreover, since the flavor transformations in the resonance regions are now known (modulo the effects of turbulence), the spectra just after the collective effects can be reconstructed from the one observed. In addition if earth effects are observed, one would know if p and $\bar{p}$ vanish or not. This would further help reconstruct the spectra before collective effects.

Of course, we first need a galactic supernova.

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References

[1] J. T. Pantaleone, Phys. Lett. B **287**, 128 (1992).

[2] G. Raffelt and G. Sigl, Astropart. Phys. **1**, 165 (1993).

[3] L. Wolfenstein, Phys. Rev. D **17**, 2369 (1978); S. P. Mikheev and A. Y. Smirnov, Sov. J. Nucl. Phys. **42**, 913 (1985) [Yad. Fiz. **42**, 1441 (1985)].

[4] S. P. Mikheev and A. Y. Smirnov, Sov. Phys. JETP **64**, 4 (1986) [Zh. Eksp. Teor. Fiz. **91**, 7 (1986)].

[5] A. S. Dighe and A. Y. Smirnov, Phys. Rev. D **62**, 033007 (2000).

[6] S. Samuel, Phys. Rev. D **48**, 1462 (1993); V. A. Kostelecký and S. Samuel, Phys. Lett. B **318**, 127 (1993); V. A. Kostelecký and S. Samuel, Phys. Rev. D **52**, 621 (1995); S. Samuel, Phys. Rev. D **53**, 5382 (1996); R. F. Sawyer, Phys. Rev. D **72**, 045003 (2005); R. F. Sawyer, Phys. Rev. D **79**, 105003 (2009).

[7] H. Duan, G. M. Fuller, J. Carlson, Y. -Z. Qian, Phys. Rev. Lett. **97**, 241101 (2006); H. Duan, G. M. Fuller and Y. Z. Qian, Phys. Rev. D **74** (2006) 123004; H. Duan, G. M. Fuller, J. Carlson and Y. Z. Qian, Phys. Rev. D **74** (2006) 105014.

[8] S. Pastor and G. G. Raffelt, Phys. Rev. Lett. **89**, 191101 (2002).

[9] S. Hannestad, G. G. Raffelt, G. Sigl and Y. Y. Y. Wong, Phys. Rev. D **74**, 105010 (2006) [Erratum-ibid. D **76**, 029901 (2007)].

[10] G. G. Raffelt and A. Y. Smirnov, Phys. Rev. D **76**, 081301 (2007); G. G. Raffelt and A. Y. Smirnov, Phys. Rev. D **76**, 125008 (2007).

[11] H. Duan, G. M. Fuller and Y. Z. Qian, Phys. Rev. D **76**, 085013 (2007).

[12] B. Dasgupta, G. G. Raffelt, I. Tamborra, Phys. Rev. **D81**, 073004 (2010).

[13] B. Dasgupta, A. Dighe, G. G. Raffelt and A. Y. Smirnov, Phys. Rev. Lett. **103**, 051105 (2009).

[14] A. Friedland, Phys. Rev. Lett. **104**, 191102 (2010).

[15] B. Dasgupta, A. Mirizzi, I. Tamborra and R. Tomas, Phys. Rev. D **81**, 093008 (2010).

[16] H. Duan and J. P. Kneller, J. Phys. G **36**, 113201 (2009); H. Duan, G. M. Fuller, Y. -Z. Qian, Ann. Rev. Nucl. Part. Sci. **60**, 569-594 (2010).

[17] G. L. Fogli, E. Lisi, A. Marrone and A. Mirizzi, JCAP **0712**, 010 (2007).

[18] A. Esteban-Pretel, S. Pastor, R. Tomàs, G. G. Raffelt and G. Sigl, Phys. Rev. D **76**, 125018 (2007).

[19] H. Duan and A. Friedland, Phys. Rev. Lett. **106**, 091101 (2011).

[20] A. Esteban-Pretel, A. Mirizzi, S. Pastor, R. Tomas, G. G. Raffelt, P. D. Serpico and G. Sigl, Phys. Rev. D **78**, 085012 (2008).

[21] S. Chakraborty, T. Fischer, A. Mirizzi, N. Saviano and R. Tomas, arXiv:1104.4031 [hep-ph]; S. Chakraborty, T. Fischer, A. Mirizzi, N. Saviano and R. Tomas, Phys. Rev. D **84**, 025002 (2011).

[22] A. Banerjee, A. Dighe, G. Raffelt, Phys. Rev. **D84**, 053013 (2011).

[23] T. Fischer, S. C. Whitehouse, A. Mezzacappa, F. K. Thielemann and M. Liebendorfer, Astron. Astrophys. **517** (2010) A80; L. Hudepohl, B. Muller, H. T. Janka, A. Marek and G. G. Raffelt, Phys. Rev. Lett. **104** (2010) 251101; T. D. Brandt, A. Burrows, C. D. Ott and E. Livne, Astrophys. J. **728**, 8 (2011).

[24] A. D. Dolgov, Sov. J. Nucl. Phys. **33**, 700 (1981) [Yad. Fiz. **33**, 1309 (1981)].

[25] B. Dasgupta and A. Dighe, Phys. Rev. D **77**, 113002 (2008).

[26] S. Choubey, B. Dasgupta, A. Dighe and A. Mirizzi, arXiv:1008.0308 [hep-ph].

[27] A. Mirizzi and P. D. Serpico, arXiv:1110.0022 [hep-ph];

[28] I. Tamborra, G. G. Raffelt, L. Huedepohl and H. -T. Janka, arXiv:1110.2104 [astro-ph.SR].

[29] R. C. Schirato and G. M. Fuller, arXiv:astro-ph/0205390; G. L. Fogli, E. Lisi, D. Montanino and A. Mirizzi, Phys. Rev. D **68**, 033005 (2003).

[30] R. Tomas, M. Kachelriess, G. Raffelt, A. Dighe, H. T. Janka and L. Scheck, JCAP **0409**, 015 (2004).

[31] G. L. Fogli, E. Lisi, A. Mirizzi and D. Montanino, JCAP **0504**, 002 (2005); G. L. Fogli, E. Lisi, A. Mirizzi and D. Montanino, JCAP **0606**, 012 (2006); S. Choubey, N. P. Harries and G. G. Ross, Phys. Rev. D **74**, 053010 (2006); B. Dasgupta and A. Dighe, Phys. Rev. D **75**, 093002 (2007); J. P. Kneller, G. C. McLaughlin and J. Brockman, Phys. Rev. D **77**, 045023 (2008); J. Gava, J. Kneller, C. Volpe and G. C. McLaughlin, Phys. Rev. Lett. **103**, 071101 (2009).

[32] A. Friedland and A. Gruzinov, arXiv:astro-ph/0607244.

[33] J. P. Kneller and C. Volpe, Phys. Rev. **D82**, 123004 (2010).

[34] T. Akiri *et al.* [The LBNE Collaboration], arXiv:1110.6249 [hep-ex].

[35] I. Gil Botella and A. Rubbia, JCAP **0310**, 009 (2003); I. Gil Botella and A. Rubbia, JCAP **0408**, 001 (2004).

[36] B. Dasgupta and J. .F. Beacom, Phys. Rev. **D83**, 113006 (2011).

[37] http://www.snolab.ca/halo/index.html

[38] M. Kachelriess, R. Tomas, R. Buras, H. T. Janka, A. Marek and M. Rampp, Phys. Rev. D **71**, 063003 (2005).

[39] H. Duan, G. M. Fuller, J. Carlson and Y. Z. Qian, Phys. Rev. Lett. **100**, 021101 (2008); C. Lunardini, B. Muller and H. T. Janka, Phys. Rev. D **78**, 023016 (2008).

[40] Y. Y. Y. Wong, AIP Conf. Proc. **655**, 240 (2003).

[41] H. Duan, G. M. Fuller and Y. Z. Qian, Phys. Rev. D **77**, 085016 (2008);

[42] B. Dasgupta, A. Dighe, A. Mirizzi and G. G. Raffelt, Phys. Rev. D **77**, 113007 (2008).

[43] J. F. Cherry, G. M. Fuller, J. Carlson, H. Duan and Y. Z. Qian, Phys. Rev. D **82**, 085025 (2010).

[44] J. F. Cherry, M. -R. Wu, J. Carlson, H. Duan, G. M. Fuller, Y. -Z. Qian, arXiv:1108.4064 [astro-ph.HE]; J. F. Cherry, M. -R. Wu, J. Carlson, H. Duan, G. M. Fuller, Y. -Z. Qian, arXiv:1109.5195 [astro-ph.HE].

[45] A. S. Dighe, M. T. Keil and G. G. Raffelt, JCAP **0306**, 005 (2003); B. Dasgupta, A. Dighe and A. Mirizzi, Phys. Rev. Lett. **101**, 171801 (2008).

[46] A. S. Dighe, M. T. Keil and G. G. Raffelt, JCAP **0306**, 006 (2003); A. S. Dighe, M. Kachelriess, G. G. Raffelt and R. Tomas, JCAP **0401**, 004 (2004).

[47] B. Dasgupta, T. Fischer, S. Horiuchi, M. Liebendorfer, A. Mirizzi, I. Sagert, J. Schaffner-Bielich, Phys. Rev. **D81**, 103005 (2010).

[48] S. Chakraborty, S. Choubey, B. Dasgupta and K. Kar, JCAP **0809**, 013 (2008);

[49] S. Galais, J. Kneller, C. Volpe and J. Gava, Phys. Rev. D **81**, 053002 (2010).

[50] H. Duan, A. Friedland, G. C. McLaughlin and R. Surman, J. Phys. G **38**, 035201 (2011).

[51] B. Dasgupta, E. P. O'Connor and C. D. Ott, arXiv:1106.1167 [astro-ph.SR]; O. Pejcha, B. Dasgupta and T. A. Thompson, arXiv:1106.5718 [astro-ph.HE].

[52] K. Abe *et al.* [T2K Collaboration], arXiv:1106.2822 [hep-ex]; Z. Isvan [MINOS Collaboration], arXiv:1110.1900 [hep-ex].

[53] G. L. Fogli, E. Lisi, A. Marrone, A. Palazzo and A. M. Rotunno, arXiv:1106.6028 [hep-ph]; T. Schwetz, M. Tortola and J. W. F. Valle, arXiv:1108.1376 [hep-ph].