

Diffuse neutrino flux from supernovae: Upper limit on the electron neutrino component from the nonobservation of antineutrinos at SuperKamiokande

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I derive an upper bound on the electron neutrino component of the diffuse supernova neutrino flux from the constraint on the antineutrino component at SuperKamiokande. The connection between antineutrino and neutrino channels is due to the similarity of the muon and tau neutrino and antineutrino fluxes produced in a supernova, and to the conversion of these species into electron neutrinos and antineutrinos inside the star. The limit on the electron neutrino flux is $5.5 \text{ cm}^{-2} \text{ s}^{-1}$ above 19.3 MeV of neutrino energy, and is stronger than the direct limit from Mont Blanc by 3 orders of magnitude. It represents the minimal sensitivity required at future direct searches, and is intriguingly close to the reach of the Sudbury Neutrino Observatory (SNO) and of the ICARUS experiment. The electron neutrino flux will have a lower bound if the electron antineutrino flux is measured. Indicatively, the first can be smaller than the second at most by a factor of 2–3 depending on the details of the neutrino spectra at production.

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The observation of the diffuse flux of neutrinos from core collapse supernovae has become a concrete possibility in the last years. If achieved, it would represent a milestone of neutrino astronomy, as it would allow faster progress with respect to searches for an individual supernova signal, and would provide several tests of neutrino properties, of the physics of core collapse and of the cosmological rate of supernovae.

Currently, the best experimental result on the diffuse supernova neutrino flux (DSN ν F from here on) is the upper limit on its electron antineutrino ($\bar{\nu}_e$) component, put by SuperKamiokande (SK) above 19.3 MeV of neutrino energy [1]. This limit comes from the nonobservation of events due to $\bar{\nu}_e$ absorption on protons in water above the background. At 90% confidence level, it reads:

$$\Phi_{\bar{e}}(E > 19.3 \text{ MeV}) < 1.2 \text{ cm}^{-2} \text{ s}^{-1}. \quad (1)$$

This bound approaches the range of theoretical predictions of the flux [2–9], suggesting that a positive signal may be seen in the near future. Such signal will add to the information given by the $\bar{\nu}_e$ data from SN1987A [10,11].

The effort to detect or constrain the electron neutrino (ν_e) component of the DSN ν F is at least as important as the search in the $\bar{\nu}_e$ channel. Indeed, it has the potential to give the very first data on *neutrinos* (instead than antineutrinos) from supernovae, with crucial first tests of theory—e.g., of neutrino transport in dense media—in the neutrino channel. In spite of its importance, however, the search of the ν_e component of the DSN ν F has had less progress than that of $\bar{\nu}_e$. This is because in water an exclusive (one flavor only) ν_e signal is not possible, and nonwater experiments are forced to smaller volumes than SK, resulting in a lower sensitivity. The best experimental information on the ν_e flux is the bound:

$$\Phi_e(25 \text{ MeV} < E < 50 \text{ MeV}) < 6.8 \cdot 10^3 \text{ cm}^{-2} \text{ s}^{-1}, \quad (2)$$

at 90% confidence level, found in 1992 at the Liquid Scintillator Detector (LSD) of Mont Blanc from the nonobservation of ν_e absorption on ^{12}C [12]. New, stronger bounds will come from the Sudbury Neutrino Observatory (SNO) [13] and from the upcoming experiment for Imaging Cosmic And Rare Underground Signals (ICARUS) [14].

In this paper I show that the ν_e diffuse flux is *already* strongly restricted by the SK result itself. Indeed, the limit on $\bar{\nu}_e$, Eq. (1), can be easily translated into an analogous result for ν_e . This new bound improves greatly on the LSD limit. Clearly, it has the same information content as the direct bound (1), and so it is important mainly because it defines the minimum sensitivity required for future direct searches in the ν_e channel.

Briefly, what allows to place a limit on the ν_e flux from the result on $\bar{\nu}_e$ is the combination of two elements. The first is the fact that the fluxes of muon and tau neutrinos and antineutrinos ($\nu_\mu, \nu_\tau, \bar{\nu}_\mu, \bar{\nu}_\tau$) produced inside the star are equal to first approximation. The second element is that neutrinos undergo flavor conversion due to masses and mixings on their way from production to detection. This implies that at least part of the detected ν_e ($\bar{\nu}_e$) flux originates from the produced ν_μ and ν_τ ($\bar{\nu}_\mu, \bar{\nu}_\tau$). Their origin from species that have similar fluxes at production provides the connection between the ν_e and $\bar{\nu}_e$ channels. This same connection will also give a *lower* bound on the ν_e component of the DSN ν F if the $\bar{\nu}_e$ flux is detected, thus giving important guidance to direct searches of the ν_e flux.

Some generalities are in order. In a core collapse supernova neutrinos are present as a thermal gas, right outside the core, from where they diffuse out carrying away about 99% of the $\mathcal{O}(10^{53})$ ergs of gravitational energy that is liberated in the collapse. This energy is shared among neutrinos and antineutrinos of the three flavors, e, μ, τ , in comparable amounts.

Before flavor conversion, the energy spectrum of each flavor of neutrinos is expected to be nearly thermal with average energy of 10–20 MeV [15]. The nonelectron species, ν_μ , ν_τ , $\bar{\nu}_\mu$ and $\bar{\nu}_\tau$, interact with matter only via neutral current exchange, which is flavor independent and, at the lowest order in momentum transfer, parity preserving [16]. This implies that ν_μ and ν_τ ($\bar{\nu}_\mu$ and $\bar{\nu}_\tau$) have equal fluxes and so they can be considered as a single species, ν_x ($\bar{\nu}_x$). The spectra of ν_x and $\bar{\nu}_x$ are similar, with differences due to weak magnetism corrections to the neutrino-nucleus cross section. These cause the ν_x to have a $\sim 7\%$ lower average energy [17]. One also expects softer spectrum of ν_e with respect to $\bar{\nu}_e$, because the interaction of ν_e ($\bar{\nu}_e$) with matter is dominated by capture on neutrons (protons) and matter is highly neutronized deep inside the star. In summary, the hierarchy of average energies $E_{0\bar{x}} \sim E_{0x} \gtrsim E_{0e} \gtrsim E_{0\bar{e}}$ is expected at the point of decoupling of neutrinos from matter, and is supported by numerical calculations (e.g. [18–20]). Indicative values of the average energies are: $E_{0e} \sim 9\text{--}16$ MeV, $E_{0\bar{e}} \sim 12\text{--}18$ MeV, $E_{0x} \sim 15\text{--}22$ MeV.

Between the production and the detection points, the neutrino fluxes are modified by two effects: the redshift of energy and the flavor conversion (oscillations) in the star and in the Earth. The latter is due to the interplay of neutrino masses, flavor mixing and coherent interaction of neutrinos with the medium (refraction). In cases where the conversion has no zenith nor energy dependence, the diffuse fluxes of ν_e and of $\bar{\nu}_e$ in a detector in a given energy bin $E_1 - E_2$ have the form:

$$\Phi_e = p\phi_e + (1 - p)\phi_x, \quad \Phi_{\bar{e}} = \bar{p}\phi_{\bar{e}} + (1 - \bar{p})\phi_{\bar{x}}, \quad (3)$$

where p (\bar{p}) is the ν_e ($\bar{\nu}_e$) survival probability, i.e. the probability that a neutrino produced as ν_e ($\bar{\nu}_e$) is detected as ν_e ($\bar{\nu}_e$). Here ϕ_w is the diffuse flux of ν_w ($w = e, \bar{e}, x, \bar{x}$) in the detector in absence of conversion. It depends on the redshift z , on the ν_w flux originally produced in a individual star, $dN_w(E')/dE'$, and on the comoving rate of supernovae, R_{SN} :

$$\phi_w \equiv \frac{c}{H_0} \int_{E_1}^{E_2} dE \int_0^{z_{\max}} \frac{dN_w(E')}{dE'} \frac{R_{SN}(z) dz}{\sqrt{\Omega_m(1+z)^3 + \Omega_\Lambda}} \quad (4)$$

(see e.g. [7]). Here E and $E' = (1+z)E$ are the neutrino energy at Earth and at the production point; $\Omega_m \simeq 0.3$ and $\Omega_\Lambda \simeq 0.7$ are the fraction of the cosmic energy density in matter and dark energy respectively; c is the speed of light and H_0 is the Hubble constant.

Generally, the conversion of neutrinos depends on E and E' , on the matter density profile along the neutrino trajectory—and thus on the arrival direction, for neutrinos crossing the Earth—and on the neutrino masses and mixings. In particular, p and \bar{p} are functions of the mixing angles θ_{12} and θ_{13} (assuming the standard parameterization of the

mixing matrix, e.g. [21]) and of the sign of the atmospheric mass squared splitting, Δm_{31}^2 . The θ_{12} angle is measured to be

$$\tan^2 \theta_{12} \simeq 0.3\text{--}0.73 \quad (5)$$

at 99.73% C.L. [22], while θ_{13} is still unknown, with the bound $\sin^2 \theta_{13} \leq 0.02$ [23]. For brevity, I will refer to the intervals $\sin^2 \theta_{13} \leq 2 \cdot 10^{-6}$ and $\sin^2 \theta_{13} \geq 3 \cdot 10^{-4}$ as “small” and “large” θ_{13} respectively. The sign of Δm_{31}^2 is also unknown; the two possibilities are called “normal hierarchy” ($\Delta m_{31}^2 > 0$) and “inverted hierarchy” ($\Delta m_{31}^2 < 0$) of the neutrino mass spectrum.

Three scenarios of conversion in the star are most relevant here (see Table I) [24–26]:

- (i) the “*adiabatic normal*”, characterized by the normal hierarchy and large θ_{13} . In this scenario neutrinos undergo resonant conversion [27] at matter density of $\sim 10^3$ g cm $^{-3}$. Their propagation is adiabatic, i.e. without transitions between the eigenstates of the Hamiltonian, resulting in $p = 0$. The propagation of antineutrinos is adiabatic too, with $\bar{p} = \cos^2 \theta_{12}$.
- (ii) the “*adiabatic inverted*”, with the inverted hierarchy and large θ_{13} . Both neutrinos and antineutrinos propagate adiabatically, but the resonant conversion affects antineutrinos, so that one has $p = \sin^2 \theta_{12}$ and $\bar{p} = 0$.
- (iii) the “*nonadiabatic*”, characterized by small θ_{13} . Here resonant conversion is made ineffective by the complete breaking of adiabaticity, and one has $p = \sin^2 \theta_{12}$ and $\bar{p} = \cos^2 \theta_{12}$ for any hierarchy.

In these three cases, p and \bar{p} have very little dependence on the neutrino energy, thus justifying the use of the expressions (3). For intermediate values of θ_{13} and normal (inverted) mass hierarchy, things are as in case (iii) with the difference that p (\bar{p}) is energy-dependent and $0 < p < \sin^2 \theta_{12}$ ($0 < \bar{p} < \cos^2 \theta_{12}$). The survival probabilities receive small corrections from oscillations in the Earth [24,26,28] and shockwave effects [29]. The former will be discussed briefly later; the latter are neglected as they are smaller than 1–2% in the energy interval where the

TABLE I. The survival probabilities p and \bar{p} , and the ratio $\Phi_e/\Phi_{\bar{e}}$ for different combinations of mass hierarchy and degree of adiabaticity of the neutrino propagation, under the conditions $\phi_e/\phi_{\bar{x}} \leq 1$ and $\phi_{\bar{e}}/\phi_x \leq 1$. The numbers in brackets give the results without those conditions. I used $\max(\phi_{\bar{x}}/\phi_x) = 1.5$. The small effects of the Earth matter and of shockwaves, as well as $O(\theta_{13}^2)$ terms are neglected.

Character	Hierarchy	p	\bar{p}	$\Phi_e/\Phi_{\bar{e}}$
adiabatic	normal	0	$\cos^2 \theta_{12}$	0.67–4.6 (< 4.6)
adiabatic	inverted	$\sin^2 \theta_{12}$	0	0.39–1 (> 0.39)
fully nonadiabatic	any	$\sin^2 \theta_{12}$	$\cos^2 \theta_{12}$	0.39–4.6 (any)

DSN ν F is largest, $E \sim 20\text{--}30$ MeV [30]. I also neglect terms of order θ_{13}^2 and higher in p and \bar{p} .

Now I derive the upper bound on the ν_e flux, Φ_e . I consider high energy thresholds, $E_1 \gtrsim 20$ MeV, motivated by the cuts necessary for background reduction at existing ν_e and $\bar{\nu}_e$ detectors [1,13]. First, I find an expression that is always greater or equal than the ratio $\Phi_e/\Phi_{\bar{e}}$:

$$\frac{\Phi_e}{\Phi_{\bar{e}}} \leq \frac{p\phi_e + (1-p)\phi_x}{(1-\bar{p})\phi_{\bar{x}}}. \quad (6)$$

This follows from Eq. (3) by dropping the $\phi_{\bar{e}}$ term there. Next, I use the fact that $\phi_x/\phi_{\bar{x}} \leq 1$, because the ν_x spectrum is slightly softer than the $\bar{\nu}_x$ one, and has average energy at or below the threshold E_1 (see e.g. fig. 4 in Ref. [17]). I also consider that the ratio (6) is largest when \bar{p} has its maximum value \bar{p}_{\max} . Thus:

$$\frac{\Phi_e}{\Phi_{\bar{e}}} \leq \frac{1}{1-\bar{p}_{\max}} \left[1 + p \left(\frac{\phi_e}{\phi_{\bar{x}}} - 1 \right) \right]. \quad (7)$$

It appears that, in general, an upper limit on $\Phi_e/\Phi_{\bar{e}}$ depends on a variety of unknowns through the factor $\phi_e/\phi_{\bar{x}}$. However, the bound simplifies to

$$\frac{\Phi_e}{\Phi_{\bar{e}}} \leq \frac{1}{1-\bar{p}_{\max}}, \quad \text{if } p = 0 \quad \text{or} \quad \phi_e/\phi_{\bar{x}} \leq 1. \quad (8)$$

From Table I I have $\bar{p}_{\max} = \cos^2\theta_{12} + \mathcal{O}(10^{-2})$, with the second term representing the (positive) correction due to oscillations in the Earth. From Eqs. (1), (5), and (8) I get:

$$\frac{\Phi_e}{\Phi_{\bar{e}}} \leq \frac{1}{\sin^2\theta_{12}} + \mathcal{O}(10^{-2}) \simeq 4.6, \quad (9)$$

$$\Phi_e(E > 19.3 \text{ MeV}) \lesssim 5.5 \text{ cm}^{-2} \text{ s}^{-1}, \quad (10)$$

if $p = 0$ or $\phi_e/\phi_{\bar{x}} \leq 1$. Explicitly, Eqs. (8)–(10) hold under one or the other of the following conditions:

- (a) the adiabatic normal scenario is realized (large θ_{13} , $p = 0$). *No other assumptions*—on the neutrino fluxes or on the supernova rate, etc.—are necessary, and thus in this case the constraint (10) is robust.
- (b) any other oscillations scenario is realized, or the conversion pattern remains unknown, and we know that $\phi_e/\phi_{\bar{x}} \leq 1$. The latter condition will be called “ ν_x dominance”. It is generally valid at $E \gtrsim 20$ MeV for neutrino fluxes motivated by theory; I will return to this point in a moment.

Equation (10) was obtained using the smallest θ_{12} in Eq. (5) and includes the correction due to Earth effects. This was estimated numerically following refs. [9,26] for a large variety of neutrino spectra and supernova rate functions; it contributes to the final result (10) by at most 5%. About the significance of the bound (10), using the likelihood information from Refs. [1,22] I find that the inequality $\Phi_e/\sin^2\theta_{12} < 5.5 \text{ cm}^{-2} \text{ s}^{-1}$ holds at 98% confidence level.

A stronger bound follows from Eq. (8) for the adiabatic inverted case ($\bar{p}_{\max} = 0$), with ν_x dominance:

$$\frac{\Phi_e}{\Phi_{\bar{e}}} \leq 1, \quad \Phi_e(E > 19.3 \text{ MeV}) \lesssim 1.2 \text{ cm}^{-2} \text{ s}^{-1}. \quad (11)$$

The more general limit, Eq. (10), is the main result of this work. It is highly motivating for ν_e detectors, as it approaches their sensitivities. These were estimated to be $\sim 6 \text{ cm}^{-2} \text{ s}^{-1}$ in the energy interval 22.5–32.5 MeV for SNO (using charged current ν_e scattering on deuterium) [13], and $\sim 1.6 \text{ cm}^{-2} \text{ s}^{-1}$ in the interval 14–40 MeV for ICARUS (using absorption in liquid Argon) [14].

One may wonder how Eq. (10) changes when one includes the dependence of neutrino conversion on the energy, on the redshift and on the arrival direction of the neutrinos. The derivation, and thus the results (7) and (8), still apply, provided that for p , \bar{p} and \bar{p}_{\max} one uses their values averaged over energy, redshift and zenith angle, with the product of neutrino spectrum and supernova rate as weight, see Eq. (4). This procedure gives the correction due to Earth effects, explaining the $\mathcal{O}(10^{-2})$ term in Eq. (9). The limits (9) and (11) remain valid, as they refer to limiting cases, where \bar{p} is truly a constant.

A second question is the validity of the condition of ν_x dominance above the threshold $E_1 \sim 20$ MeV. While there are no direct tests of this, the condition is supported by current supernova simulations, which predict the ν_e 's to have comparable luminosity (within a factor of 2) and a softer spectrum than $\bar{\nu}_e$ and ν_x . Figure 1 gives the ratio $\phi_e/\phi_{\bar{x}}$ calculated using Eq. (4) with typical parameters (see caption). One sees that the ν_x dominance is violated by at most $\sim 20\%$ when the ν_x component is very soft and suppressed in luminosity.

What can be said on the ν_e flux if evidence of the DSN ν F is not seen in the ν_e channel but is found in the $\bar{\nu}_e$ one? Similarly to what shown so far, a positive detection

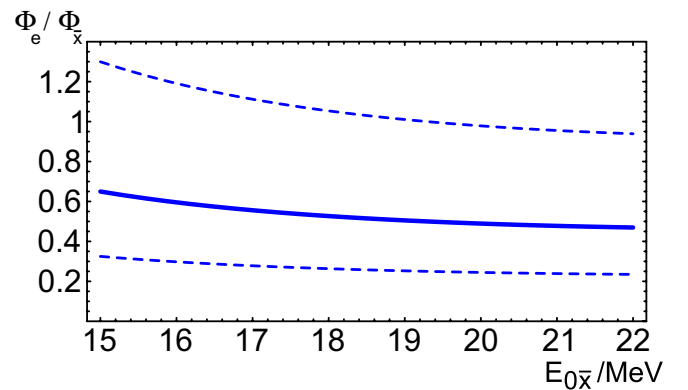


FIG. 1 (color online). The ratio $\phi_e/\phi_{\bar{x}}$ as a function of $E_{0\bar{x}}$, for $E_{0e} = 12$ MeV and ratio of ν_e and $\bar{\nu}_x$ luminosities of 2, 1, 0.5 (upper to lower curves). I parametrized the neutrino spectra and the supernova rate function as in [9], with $\alpha_e = \alpha_{\bar{x}} = 2.3$, $\alpha = 0$, $\beta = 3.4$. The SuperKamiokande energy threshold $E_1 = 19.3$ MeV was used.

of $\bar{\nu}_e$ can be translated into a *lower limit* on the ν_e component. Indeed, one can constrain the ratio $\Phi_{\bar{\nu}_e}/\Phi_e$ from above, using nearly the same steps as in Eqs. (6)–(11), with the exchanges $e \leftrightarrow \bar{e}$, $x \leftrightarrow \bar{x}$, $p \leftrightarrow \bar{p}$. The analogous of Eqs. (9) and (10) are:

$$\frac{\Phi_{\bar{\nu}_e}}{\Phi_e} \leq \frac{\max(\phi_{\bar{x}}/\phi_x)}{\cos^2\theta_{12}} + \mathcal{O}(10^{-2}), \quad (12)$$

$$\Phi_e \geq \Phi_{\bar{\nu}_e} \left[1.73 \max\left(\frac{\phi_{\bar{x}}}{\phi_x}\right) + \mathcal{O}(10^{-2}) \right]^{-1}, \quad (13)$$

where the largest value of θ_{12} from Eq. (5) was used. These results are valid for the adiabatic inverted case ($\bar{p} = 0$) or for ν_x dominance, $\phi_{\bar{\nu}_e}/\phi_x \leq 1$. They depend on the maximum value allowed by theory of the ratio $\phi_{\bar{x}}/\phi_x$, $\max(\phi_{\bar{x}}/\phi_x)$, just like the derivation of Eqs. (9) and (10) contains the maximum of $\phi_x/\phi_{\bar{x}}$. The latter was set to 1 there.

To estimate $\max(\phi_{\bar{x}}/\phi_x)$ requires a dedicated work, with the evaluation of the dependence of $\phi_{\bar{x}}/\phi_x$ on the details of the ν_x and $\bar{\nu}_x$ original spectra. For illustration, here I take $\max(\phi_{\bar{x}}/\phi_x) = 1.5$, which is the largest ratio of the $\bar{\nu}_x$ and ν_x spectra found in Ref. [17] in the energy interval $E = 20\text{--}80$ MeV. With this value Eq. (13) gives $\Phi_e \geq 0.39\Phi_{\bar{\nu}_e} - \mathcal{O}(10^{-2})$ (the minus sign indicates a negative Earth effect).

For the adiabatic normal scenario ($p = 0$) with the ν_x dominance one finds the analogous of Eq. (11):

$$\frac{\Phi_{\bar{\nu}_e}}{\Phi_e} \leq \max(\phi_{\bar{x}}/\phi_x). \quad (14)$$

This gives $\Phi_e \geq 0.67\Phi_{\bar{\nu}_e}$ for $\max(\phi_{\bar{x}}/\phi_x) = 1.5$.

The results of this paper are summarized in Table I. The Table helps the discussion that follows.

What will we learn from a detection of the DSN ν F in the ν_e channel, together with the limits (10) and (11)? If evidence is found above the value in (10), the physics

and/or the data analyses involved would have to be revised, e.g. by relaxing possible priors or by reconsidering background subtraction. If both the ν_e and the $\bar{\nu}_e$ data analyses are proved to be robust, one would have to admit the possibility of a very large ν_e component in the flux before neutrino conversion (i.e. $\phi_e > \phi_x$ above the detection threshold), or invoke exotic physics that could give a larger $\bar{\nu}_e$ survival probability, \bar{p}_{\max} (see Eq. (7)). A measurement above the limit (11) would exclude the adiabatic inverted scenario with ν_x dominance. A ν_e detection would also constrain the $\bar{\nu}_e$ flux indirectly. For example, Eq. (13) (with $\max(\phi_{\bar{x}}/\phi_x) = 1.5$) tells that a measured ν_e flux smaller than $\sim 0.48 \text{ cm}^{-2} \text{ s}^{-1}$ above 19.3 MeV would be incompatible with a $\bar{\nu}_e$ flux as large as $1.2 \text{ cm}^{-2} \text{ s}^{-1}$ in the same energy bin, thus improving the SK bound for $\bar{\nu}_e$, Eq. (1).

In summary, I found an upper limit of $\Phi_e \leq 5.5 \text{ cm}^{-2} \text{ s}^{-1}$ above 19.3 MeV for the ν_e component of the diffuse supernova neutrino flux, using the SK bound on the $\bar{\nu}_e$ flux, the properties of supernova neutrino spectra and neutrino oscillations. This new limit is the strongest available, improving the LSD bound by 3 orders of magnitude. It defines the necessary sensitivity for future searches of the ν_e flux. Noticeably, the result proves that the ν_e flux needs not be equal or smaller than the $\bar{\nu}_e$ one—as it would be the case without oscillations—but instead it can be larger by a factor of several above $E \sim 20$ MeV, depending on the oscillation scenario. It follows that, in spite of their smaller volumes, ν_e telescopes can be competitive with SK in the search of the diffuse supernova neutrino flux. A dedicated analysis of the existing SNO data is compelling, in this perspective.

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