

Relic neutrino absorption spectroscopy

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Resonant annihilation of extremely high-energy cosmic neutrinos on big-bang relic anti-neutrinos (and vice versa) into Z bosons leads to sizable absorption dips in the neutrino flux to be observed at Earth. The high-energy edges of these dips are fixed, via the resonance energies, by the neutrino masses alone. Their depths are determined by the cosmic neutrino background density, by the cosmological parameters determining the expansion rate of the Universe, and by the large redshift history of the cosmic neutrino sources. We investigate the possibility of determining the existence of the cosmic neutrino background within the next decade from a measurement of these absorption dips in the neutrino flux. As a by-product, we study the prospects to infer the absolute neutrino mass scale. We find that, with the presently planned neutrino detectors (ANITA, Auger, EUSO, OWL, RICE, and SalSA) operating in the relevant energy regime above 10^{21} eV, relic neutrino absorption spectroscopy becomes a realistic possibility. It requires, however, the existence of extremely powerful neutrino sources, which should be opaque to nucleons and high-energy photons to evade present constraints. Furthermore, the neutrino mass spectrum must be quasidegenerate to optimize the dip, which implies $m_\nu \gtrsim 0.1$ eV for the lightest neutrino. With a second generation of neutrino detectors, these demanding requirements can be relaxed considerably.

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I. INTRODUCTION

Neutrinos are the elementary particles with the weakest known interactions. Correspondingly, they can propagate to us through the cosmic microwave and neutrino background (CMB and $C\nu B$, respectively) without significant energy loss even from cosmological distances. A possible exception to this transparency is resonant annihilation of extremely high energy cosmic neutrinos (EHEC ν 's) on big-bang relic antineutrinos (and vice versa) into Z bosons [1–4,85]. This occurs near the respective resonance energies,

$$E_{\nu_i}^{\text{res}} = \frac{m_Z^2}{2m_{\nu_i}} = 4.2 \times 10^{22} \text{ eV} \left(\frac{0.1 \text{ eV}}{m_{\nu_i}} \right), \quad (1)$$

with $m_Z = 91.2$ GeV denoting the mass of the Z boson [5] and m_{ν_i} ($i = 1, 2, 3$) the nonzero neutrino masses—for which there is rather convincing evidence inferred from the apparent observation of neutrino oscillations [5]. On resonance, the corresponding cross sections are enhanced by several orders of magnitude. This leads to a few percent probability of annihilation within the Hubble radius of the Universe, even if one neglects further enhancing effects due to cosmic evolution. Indeed, it appears that—apart from the indirect evidence to be gained from cosmology, e.g., big-bang nucleosynthesis and large-scale structure formation—this annihilation mechanism is the unique process [6] having sen-

sitivity to the $C\nu B$ [1]. Moreover, observation of the absorption dips would present one of the few opportunities to determine absolute neutrino masses [7]. However, the mechanism requires that there exists a sufficiently large EHEC ν flux at the resonant energies of Eq. (1). One of the purposes of this work is to quantify how large this EHEC ν flux must be.

Since the original proposal in 1982 [1], significant advances have been made in theoretical and observational cosmology, experimental neutrino physics, and EHEC ν physics. Each of these three areas impacts immediately on the $C\nu B$ measurement. Thus, it is timely to investigate again the possibility of determining the existence of the $C\nu B$ and to study the prospects for determining the neutrino masses via resonant neutrino absorption.

What are the new findings that affect the EHEC ν absorption probability?

(a) The original calculation of neutrino absorption was done for a matter-dominated flat universe without a cosmological constant. Recent observations of large-scale gravity, deep-field galaxy counts, and type Ia supernovae favor a universe with energy fractions $\Omega_\Lambda \approx 0.7$ in the cosmological constant and $\Omega_M \approx 0.3$ in (mainly cold and dark) matter [8,9]. The position of the first Doppler peak in recent CMB measurements strongly suggests that the Universe is flat, i.e. the fractional contribution of curvature energy Ω_k is negligibly small. These cosmological parameters, together with the Hubble parameter H_0 , determine the expansion rate of the

Universe as a function of look-back distance. This expansion history, in turn, crucially impacts the EHEC ν fluxes at Earth, since their sources are almost certainly located at cosmological distances [10].

(b) The oscillation interpretation of the atmospheric neutrino data offers a lower limit on the heaviest of the three mass eigenstates of

$$m_{\nu_3} \geq \sqrt{\Delta m_{\text{atm}}^2} > 0.04 \text{ eV} \quad (2)$$

at 95% confidence level (CL), from the inferred mass splitting Δm_{atm}^2 [11]. On the other hand, studies of the cosmic evolution of the large-scale structure, as observed today, from primordial density perturbations, measured in the CMB, yield an upper bound on the sum of neutrino masses [8,9,12]

$$\sum_i m_{\nu_i} \leq 1.2 \text{ eV}. \quad (3)$$

Since oscillation studies also reveal that neutrino mass splittings are small compared to the eV scale, the cosmic bound per mass state is conservatively [13] ~ 0.4 eV—about a factor of five better than the laboratory bounds inferred from tritium beta decay [14,15], $m_{\nu_3} < 2.2$ eV (95% CL) [16], and neutrinoless double beta decay [17], $m_{\nu_3} \leq (0.66\text{--}2.70)$ eV [18]. Thus we have

$$0.04 \text{ eV} < m_{\nu_3} \leq 0.4 \text{ eV}. \quad (4)$$

It is remarkable that the neutrino mass, whose value was compatible with zero at the time of the original proposal for Z dips [1], is now known to be not only nonzero, but to lie within a one-order-of-magnitude range! Accordingly, the resonant annihilation energy for the heaviest mass state is also known to within an order of magnitude:

$$1 \times 10^{22} \text{ eV} \leq E_{\nu_3}^{\text{res}} < 1 \times 10^{23} \text{ eV}. \quad (5)$$

This resonant energy, when divided by the Z-decay multiplicity of ~ 40 , predicts secondary cosmic ray particles with energies spanning a decade or more above the Greisen-Zatsepin-Kuzmin (GZK) energy, $E_{\text{GZK}} = 4 \times 10^{19}$ eV. This is the energy beyond which the CMB is absorbing to nucleons, due to resonant photopion production [19]. The association of Z bursts with the mysterious cosmic rays observed above E_{GZK} is a controversial possibility [20–25]. Intriguingly, the neutrino mass window required in such a scenario coincides quantitatively with Eq. (4) [24].

(c) Recent proposals for progressively larger EHEC ν detectors, such as the Pierre Auger Observatory [26], IceCube [27], ANITA [28], EUSO [29], OWL [30], and SaLSA [31] offer credible hope that the collection of an event sample above 10^{21} eV may be realized within this decade [32] (cf. Fig. 1). Another encouraging sign for the progress in experimental sensitivity is that existing EHEC ν “observatories,” such as RICE [34], GLUE [35], and FORTE [36], have recently put the first sensible upper limits on the EHEC ν flux

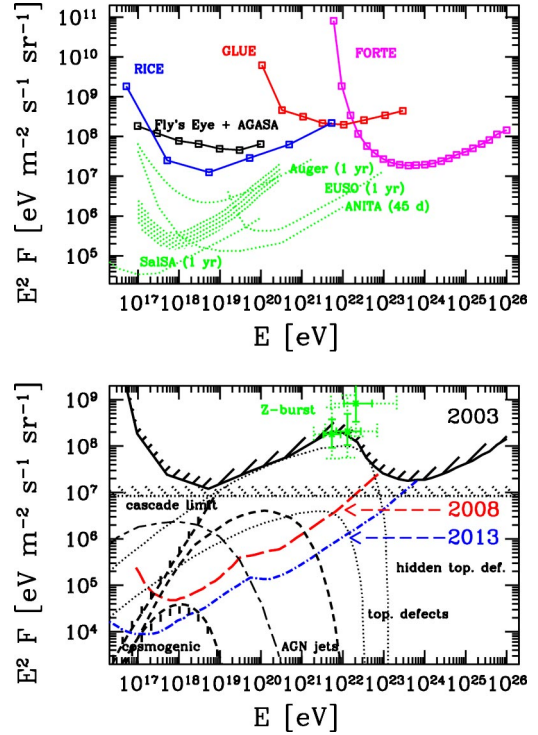


FIG. 1. Current status and next decade prospects for EHEC ν physics, expressed in terms of diffuse neutrino fluxes per flavor, $F_{\nu_\alpha} + F_{\bar{\nu}_\alpha}$, $\alpha = e, \mu, \tau$; full mixing among the flavors en route to Earth [33] is assumed. *Top*: Upper limits from RICE [37], GLUE [38], FORTE [39], and Fly’s Eye and AGASA [40]. Also shown are projected sensitivities of Auger in ν_e , ν_μ modes and in ν_τ mode (bottom swath) [41], ANITA [42], EUSO [43], and SaLSA [44], corresponding to one event per energy decade and indicated duration. *Bottom*: Roadmap for improvement in the next decade (2008 and 2013; adapted from Ref. [32]), corresponding to one event per energy decade, as well as the current (2003) observational upper bound (solid-shaded) obtained from Fig. 1 (top). For the year 2008 (long-dashed), we assume 3 yr of Auger data and one 15 d ANITA flight. For 2013 (dashed-dotted), we assume 8/3/4 yr Auger/EUSO/SaLSA, and 3 ANITA flights. The sensitivity will improve if further projects such as Auger North and OWL [30] are realized, or if the EUSO flight time is extended. Also shown is a wide sample of predictions for EHEC ν fluxes (discussed in Sec. II B).

in the region (5) relevant for neutrino absorption [37–39]. We display these limits in Fig. 1 (top).

The organization of this paper is as follows: In Sec. II, we determine the EHEC ν spectra, with their all-important absorption features, to be observed at Earth. In particular, the depths, widths, and locations of the relic neutrino absorption dips are calculated, for various proposed sources of a diffuse EHEC ν flux. The experimental prospects to detect the absorption dips within the next decade or beyond are discussed in Sec. III. Finally, in Sec. IV, we summarize our results and conclude.

II. EHEC ν SPECTRAL DIPS

Given an EHEC ν source emissivity distribution, one can determine the resulting neutrino spectrum to be observed at

Earth, by taking into account resonant annihilation with the $C\nu B$, and energy losses due to redshift. We formulate this problem, in Sec. II A, in terms of propagation functions [45,46]. We calculate these functions by means of the procedure outlined in the original papers on relic neutrino absorption spectroscopy [1,2], updated to modern cosmological parameters. The observable neutrino flux arriving at Earth is then obtained by folding the propagation function with the EHEC ν source emissivity distribution. The latter basic input is not known in the energy region of interest. Therefore, we introduce in Sec. II B various parameterizations to model neutrino emission from the most relevant classes of possible sources—astrophysical accelerators (bottom-up “Zevatrions”) or exotic massive particles and topological defects (top-down). In Sec. II C, we study the location, depths, and widths of relic neutrino absorption dips in the context of these model classes of sources, and for various neutrino mass scenarios.

A. Neutrino propagation functions

Let $\mathcal{L}_{\nu_\beta}(r, E_i)$ be the EHEC ν source emissivity distribution, i.e. the number of neutrinos ν_β of flavor $\beta=e, \mu, \tau$, per co-moving volume, per unit of time and per unit of energy as measured at the source, injected in the $C\nu B$ at a “propagation distance” $r=ct$ from Earth [47] with an energy E_i . The propagation through the $C\nu B$ can be described [45,46] by the functions $P_{b|a}(E; E_i, r)$, which are defined as the expected number of particles of type b above the threshold energy E if one particle of type a started at a distance r with energy E_i . With the help of these propagation functions, the differential flux of neutrinos of flavor α ($b=\nu_\alpha$) at Earth, i.e., their number N_{ν_α} arriving at Earth with energy E per units of energy, area (A), time (t) and solid angle (Ω), can be expressed as

$$F_{\nu_\alpha}(E) \equiv \frac{d^4 N_{\nu_\alpha}}{dE dA dt d\Omega} = \frac{1}{4\pi} \int_0^\infty dE_i \int_0^\infty dr \sum_\beta \frac{-\partial P_{\nu_\alpha|\nu_\beta}(E; E_i, r)}{\partial E} \mathcal{L}_{\nu_\beta}(r, E_i), \quad (6)$$

where the “propagation distance” $r=ct$ is related to the redshift z by

$$dz = (1+z)H(z)dr, \quad (7)$$

with the evolving Hubble parameter given by

$$H^2(z) = H_0^2 [\Omega_M(1+z)^3 + \Omega_k(1+z)^2 + \Omega_\Lambda]. \quad (8)$$

Here, $H_0 = h100$ (km/s)/Mpc, with $h = (0.71 \pm 0.07) \times_{0.95}^{1.15}$ [5], denotes the present value of the Hubble parameter. In Eq. (8), Ω_M , Ω_k , and Ω_Λ are the present matter, curvature, and vacuum energy densities, respectively, in terms of the critical density. From the Friedmann equation comes the constraint that fractional energies must sum to 100%, i.e. $\Omega_M + \Omega_k + \Omega_\Lambda = 1$. As default values we choose $\Omega_M = 0.3$, Ω_k

$= 0$, and $\Omega_\Lambda = 0.7$. These values are collectively known as the “concordance” model, for they fit a wide assortment of data [8,9,48].

Justified approximations are that (i) the only type of energy loss is due to the redshift, and (ii) the only relevant interaction is absorption on the relic neutrino background, dominated by resonant Z production [49]. With these approximations, the differential propagation function is given by (see the Appendix for a thorough derivation, properly taking into account neutrino mixing effects)

$$\frac{-\partial P_{\nu_\alpha|\nu_\beta}(E; E_i, z)}{\partial E} = \delta\left(E - \frac{E_i}{1+z}\right) \sum_j |U_{\alpha j}|^2 P_{\nu_j}(E_i, z) |U_{\beta j}|^2, \quad (9)$$

where $U_{\alpha j}$ is the leptonic mixing matrix and $P_{\nu_j}(E_i, z)$ is the survival probability of a cosmic neutrino ν_j injected at a redshift z with energy $E_i = E(1+z)$,

$$P_{\nu_j}(E(1+z), z) = \exp\left[-\int_0^z \frac{d\tilde{z}}{H(\tilde{z})(1+\tilde{z})} n_{\tilde{\nu}_j}(\tilde{z}) \sigma_{\nu_j \tilde{\nu}_j}^{\text{ann}}(s)\right], \quad (10)$$

with

$$s = 2m_{\nu_j} E(1+\tilde{z}). \quad (11)$$

We remind the careful reader that z is the emission redshift of the cosmic source, while \tilde{z} is the redshift at the time of neutrino annihilation, the latter being integrated from present time back to the emission time. The exponent in the brackets is (minus) the optical depth (also called the “opacity”) for a neutrino emitted at redshift z .

The momentum of a massless neutrino today is of order of the CMB energy, $\sim 3T_\gamma \sim 0.7$ meV. The neutrino momentum at an earlier epoch is then $\sim 0.7(1+z)$ meV. Thus, relic neutrinos will be nonrelativistic as long as $m_{\nu_j} \gg \langle p_{\nu_j} \rangle \sim 0.7(1+z)$ meV holds. In most of the region in redshift space which we consider, the relic neutrinos are nonrelativistic. This means that their number densities are given by [5]

$$n_{\nu_j}(\tilde{z}) = n_{\tilde{\nu}_j}(\tilde{z}) = \langle n_\nu \rangle_0 (1+\tilde{z})^3 = \frac{3}{22} \langle n_\gamma \rangle_0 (1+\tilde{z})^3 = 56 \text{ cm}^{-3} (1+\tilde{z})^3, \quad (12)$$

with $\langle n_\nu \rangle_0$ and $\langle n_\gamma \rangle_0$ denoting today’s (subscript “0”) average number density in relic neutrinos and relic photons, respectively [50].

The annihilation cross section $\sigma_{\nu_j \tilde{\nu}_j}^{\text{ann}}(s)$, which is dominated by the Z production peak with a rather narrow width, can be approximated by [1]

$$\sigma_{\nu_j \tilde{\nu}_j}^{\text{ann}}(s) = \langle \sigma_{\nu \tilde{\nu}}^{\text{ann}} \rangle \delta\left(\frac{s}{m_Z^2} - 1\right), \quad (13)$$

where

$$\langle \sigma_{\nu\bar{\nu}}^{\text{ann}} \rangle = \int \frac{ds}{m_z^2} \sigma_{\nu_j \bar{\nu}_j}^{\text{ann}} = 2\pi\sqrt{2}G_F = 40.4 \text{ nb} \quad (14)$$

is the energy averaged cross section [51] $G_F = 1.2 \times 10^{-5} \text{ GeV}^{-2}$ being the Fermi coupling constant [5]. Therefore, the survival probability (10), at the injection energy $E_i = E(1+z)$, can be approximated by [1]

$$P_{\nu_j}(E(1+z), z) \approx \exp \left[- \left(\frac{\langle n_{\nu} \rangle_0 \langle \sigma_{\nu\bar{\nu}}^{\text{ann}} \rangle}{H_0} \right) \times \frac{(E_{\nu_j}^{\text{res}}/E)^3}{[\Omega_M (E_{\nu_j}^{\text{res}}/E)^3 + \Omega_k (E_{\nu_j}^{\text{res}}/E)^2 + \Omega_\Lambda]^{1/2}} \right],$$

for $\frac{1}{1+z} < \frac{E}{E_{\nu_j}^{\text{res}}} < 1$, (15)

within the region of support indicated, and identically one (no absorption) outside the region of support. Here, $E_{\nu_j}^{\text{res}}$ is the resonant energy in the rest system of the relic neutrinos, given in Eq. (1).

Numerically, the annihilation probability on the CνB, neglecting cosmological expansion, is on the few percent level,

$$\frac{\langle n_{\nu} \rangle_0 \langle \sigma_{\nu\bar{\nu}}^{\text{ann}} \rangle}{H_0} = (0.71/h) \times 3.0\%. \quad (16)$$

Taking cosmological expansion into account, the annihilation probability is enhanced by the redshift-dependent ratio in the exponent appearing in the expression (15) for the survival probability. The enhancement is easy to understand: In the numerator of the ratio, the factor $(E_{\nu_j}^{\text{res}}/E)^3 = (1+\bar{z})^3$ accounts for the higher target density of the CνB in physical volume. The denominator accounts for the time or path length available per unit redshift for annihilation. What is noteworthy here is that, at early times, the Ω_M term dominates, and so the neutrino optical depth scales as $1/\sqrt{\Omega_M}$. Thus, the smaller the Ω_M , or equivalently, the larger the Ω_Λ , the greater the neutrino absorption probability from sources with $1+z \gtrsim (\Omega_\Lambda/\Omega_M)^{1/3}$. For example, the Λ CDM universe, with $\Omega_M = 0.3$ and $\Omega_\Lambda = 0.7$, produces absorption dips nearly twice as deep as a pure CDM universe with $\Omega_M = 1$. In turn, this alleviates the statistics requirement (discussed later) by a factor of ~ 3 .

In Fig. 2, we display the survival probability (15) for the modern concordance cosmological parameters, i.e. $\Omega_M = 0.3$, $\Omega_k = 0$, $\Omega_\Lambda = 0.7$, and $h = 0.71$ (top), and their allowed variations (bottom), respectively. It seems that variations of Ω_M , Ω_Λ , and h , within their uncertainties [8,9], amount to a $\sim 5\%$ effect. As is apparent from Fig. 2 (top), using the narrow-width approximation (13) for the annihilation

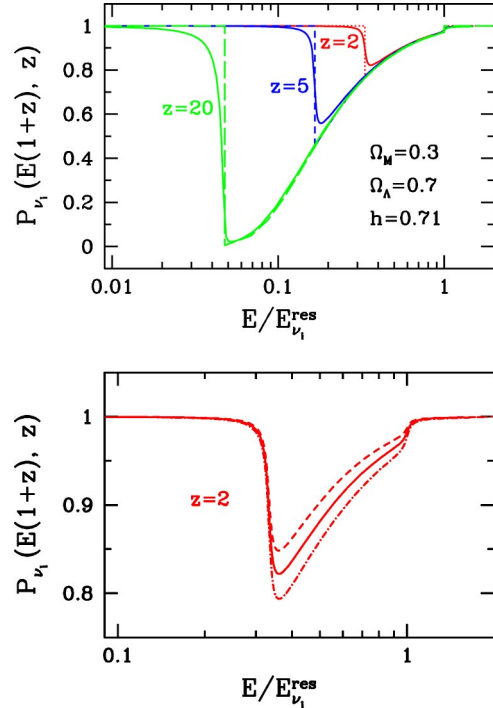


FIG. 2. The survival probability $P_{\nu_i}(E_i, z)$ of a cosmic neutrino ν_i , injected at redshift z with energy E_i , as a function of the energy at Earth, $E = E_i/(1+z)$, in units of the resonance energy $E_{\nu_i}^{\text{res}} = m_z^2/2m_{\nu_i}$. *Top*: The narrow-width approximation (15), for $z=2$ (dotted), $z=5$ (short-dashed), $z=20$ (long-dashed), and standard cosmological parameters ($\Omega_M = 0.3$, $\Omega_\Lambda = 0.7$, $h = 0.71$), compared with the complete energy dependence from the annihilation cross section of Ref. [3] (solid). *Bottom*: The survival probability for $z=2$ and standard cosmological parameters (solid) compared with the most extreme variations allowed by up-to-date global fits: $\Omega_M = 0.20$, $\Omega_\Lambda = 0.78$, $h = 0.81$ (dashed) and $\Omega_M = 0.40$, $\Omega_\Lambda = 0.61$, $h = 0.62$ (dashed-dotted) [9].

cross section rather than taking into account the complete energy dependence of the cross section is justified within the overall 5% error.

B. Neutrino source emissivity distributions

In the previous subsection, we have found that, for a given source emissivity distribution \mathcal{L}_{ν_β} of neutrinos of flavor β , the neutrino flux of flavor α to be observed at Earth is predicted to be [cf. Eqs. (6)–(9)]

$$F_{\nu_\alpha}(E) = \frac{1}{4\pi} \int_0^\infty \frac{dz}{H(z)} \sum_{\beta, j} |U_{\alpha j}|^2 P_{\nu_j}(E(1+z), z) \times |U_{\beta j}|^2 \mathcal{L}_{\nu_\beta}(z, E(1+z)), \quad (17)$$

where $\mathcal{L}_{\nu_\beta}(z, E(1+z))$ is the number of neutrinos of flavor β and energy $E_i = E(1+z)$ emitted per co-moving volume, per unit time and unit energy, at a redshift “distance” z [cf. Eq. (7)]. In this subsection and the one that follows, we evaluate this expression further for some benchmark source emissivities.

We will focus mainly on sources that produce pions (“hadronic” sources), be they astrophysical Zevatrons (bottom-up) or nonaccelerator (top-down) ones. From the decay sequence $\pi^\pm \rightarrow \mu^\pm + \nu_\mu \rightarrow e^\pm + 2\nu_\mu + \nu_e$ (neither we, nor experiments, will distinguish neutrinos from antineutrinos), the flavor ratios of the source emissivities are predicted to be

$$\mathcal{L}_{\nu_e} : \mathcal{L}_{\nu_\mu} : \mathcal{L}_{\nu_\tau} = 1 : 2 : 0. \quad (18)$$

In this case, one finds—exploiting the fact that phenomenologically $|U_{e3}|^2 \ll 1$ and $|U_{\mu 3}| \approx |U_{\tau 3}|$ —that the fluxes at Earth are well approximated by (see the Appendix and also Ref. [33])

$$F_{\nu_\alpha}(E) \approx \frac{1}{4\pi} \int_0^\infty \frac{dz}{H(z)} \frac{1}{3} \mathcal{L}_\nu^{\text{tot}}(z, E(1+z)) \sum_{j=1}^3 |U_{\alpha j}|^2 \times P_{\nu_j}(E(1+z), z), \quad (19)$$

where $\mathcal{L}_\nu^{\text{tot}}$ is the total, flavor-summed neutrino emissivity at the source. In fact, as discussed in the Appendix, Eq. (19) holds whenever the source emissivities satisfy $\mathcal{L}_{\nu_\mu} + \mathcal{L}_{\nu_\tau} = 2\mathcal{L}_{\nu_e}$. Thus, it holds also for a “democratic” flavor ratio of source emissivities,

$$\mathcal{L}_{\nu_e} : \mathcal{L}_{\nu_\mu} : \mathcal{L}_{\nu_\tau} = 1 : 1 : 1, \quad (20)$$

as might arise from the decay of topological defects not coupled directly to ordinary matter such as, for example, mirror-matter “necklaces” [52].

It is unlikely that neutrino observatories will be able to distinguish neutrino flavors at extreme high energy [53]. Phenomenologically then, we are mainly interested in the sum over all neutrino flavors α . With such a sum, unitarity completely removes the dependence on the leptonic mixing matrix from Eq. (19), leaving just

$$\sum_\alpha F_{\nu_\alpha}(E) \approx \frac{1}{4\pi} \int_0^\infty \frac{dz}{H(z)} \frac{1}{3} \mathcal{L}_\nu^{\text{tot}}(z, E(1+z)) \times \sum_{j=1}^3 P_{\nu_j}(E(1+z), z). \quad (21)$$

Another simplification of Eq. (19) occurs if the neutrino masses are quasidegenerate, which in fact is realized in Nature if the lightest neutrino has a mass $m_{\nu_1} \gg \Delta m_{\text{atm}}^2$, say, $m_{\nu_1} \gtrsim 0.1$ eV (cf. Fig. 3):

$$m_{\nu_1} : m_{\nu_2} : m_{\nu_3} \approx 1 : 1 : 1 \Rightarrow P_{\nu_1} \approx P_{\nu_2} \approx P_{\nu_3}. \quad (22)$$

With quasi-mass-degeneracy, Eq. (19) simplifies to

$$F_{\nu_\alpha}(E) \approx \frac{1}{4\pi} \int_0^\infty \frac{dz}{H(z)} P_{\nu_1}(E(1+z), z) \frac{1}{3} \mathcal{L}_\nu^{\text{tot}}(z, E(1+z)), \quad (23)$$

for each $\alpha = e, \mu, \tau$.

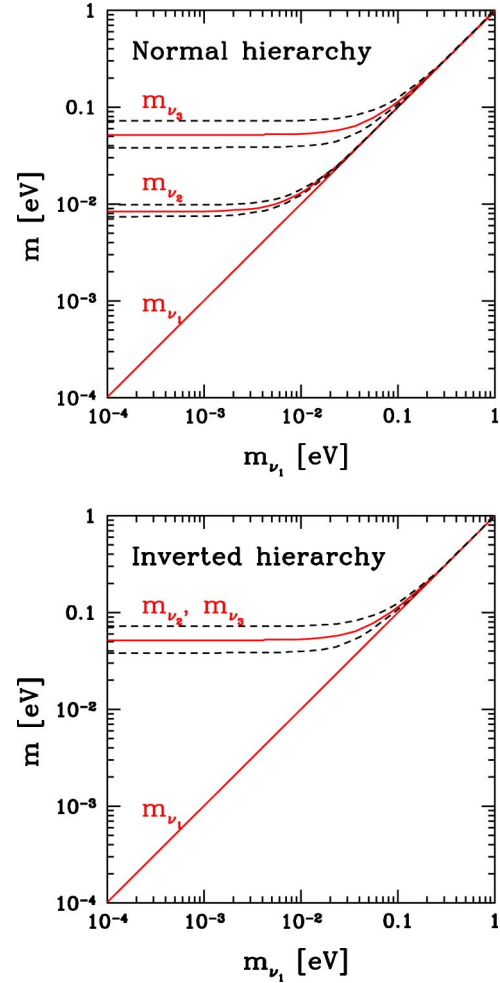


FIG. 3. Allowed ranges for the neutrino masses as a function of the lightest neutrino mass m_{ν_1} , in the normal (top) and inverted (bottom) three-neutrino scheme (adapted from Refs. [54,55]).

In the following, we assume that the EHEC ν sources, which build up the diffuse source emissivity, have identical injection spectra J_{ν_β} . Consequently, the z and E_i dependences of the source emissivity distribution factorize,

$$\mathcal{L}_{\nu_\beta}(z, E_i) = \eta(z) J_{\nu_\beta}(E_i). \quad (24)$$

The “activity” $\eta(z)$ is the number of sources per co-moving volume and per unit of time, at the redshift “time” z , while the injection spectra $J_{\nu_\beta}(E_i)$ are the number of neutrinos ν_β emitted by a single source per unit of energy. The z dependence of the activity accounts for any evolution of the co-moving number density and/or of the common luminosity of the sources.

We employ several parameterizations of $\eta(z)$, which allow us to study a variety of possible EHEC ν origins—ranging from astrophysical accelerator sources such as gamma ray bursts (GRB’s) and active galactic nuclei (AGN’s) which turned on at $z \sim \text{few}$, to non-accelerator sources such as topological defects which have been decaying all the way back to z very large.

We start with a parametrization of the activity motivated by astrophysics [56],

$$\eta_{\text{SFR}}(z) = \eta_0 \frac{1+a}{(1+z)^{-n_1+a(1+z)^{n_2}},} \quad (25)$$

with $\eta_0 = \eta(z=0)$ being the activity in the today's epoch. With the values $a=0.005$ (0.0001), $n_1=3.3$ (4.0), and $n_2=3.0$ (3.0), the parametrization fits the star formation rate (SFR) history derived from the blue and UV luminosity density, in line with the extreme ranges of optical/UV measurements without [57] (with [58]) dust extinction corrections. We will refer to these two cases as conservative and generous SFR, respectively. The star formation rate is believed to map out the earliest structures sufficiently bound by gravity to undergo fusion. As such, they may map out the history of AGN and GRB evolution, two potential sources of the EHEC rays (EHECR's) [56].

The parametrization (25) provides a peak activity at

$$1+z_{\text{peak}} = \left(\frac{n_1}{an_2} \right)^{1/(n_1+n_2)}, \quad (26)$$

which occurs at $z_{\text{peak}}=1.4$ for the conservative SFR and at $z_{\text{peak}}=2.9$ for the generous one. For $n_2=0$, $a=0$ it reduces to a simple power law, $\eta(z) = \eta_0 (1+z)^{n_1}$. Such a choice, with $n_1 \sim 4$ to fit the low- z SFR, is often used in the literature. Without a cutoff in z , however, the power law provides an extreme evolution history, as the activity increases indefinitely.

This brings us to a further parametrization, namely a simple power-law ansatz, but with cutoffs z_{min} and z_{max} to exclude the existence of nearby and early-time sources, respectively:

$$\eta_{\text{pow}}(z) = \eta_0 (1+z)^n \theta(z-z_{\text{min}}) \theta(z_{\text{max}}-z). \quad (27)$$

This ansatz has the advantage of leading to easily tractable analytic expressions, while still reproducing, for $n=4$, the generous SFR case, as long as $z_{\text{max}} < z_{\text{peak}}$. In addition to describing the evolution of GRB's and AGN's up to $z \sim 2$, the power law with $n \approx 1-2$ and $z_{\text{max}} \gg 1$ also characterizes the activity expected from nonaccelerator sources. For example, topological defect sources [59] are characterized by $n=3/2$ [60] and z_{max} arbitrarily large [61]. Throughout, we will take $z_{\text{min}}=0$ as a default value [62].

Turning to the source injection spectra $J_{\nu_\beta}(E_i)$, we only need to specify them in the energy decade around the respective resonance energies. For this energy decade, we again assume a power-law behavior [63],

$$J_{\nu_\beta}(E_i) = j_{\nu_\beta} E_i^{-\alpha} \theta(E_i - E_{\text{min}}) \theta(E_{\text{max}} - E_i). \quad (28)$$

As with the activity power law, this spectral power-law parametrization facilitates the analytic study of different scenarios for the origin of EHEC ν 's.

In most of the source models, the neutrinos originate from pion decays, the latter either being produced in inelastic pp or $p\gamma$ interactions (astrophysical sources), or, alternatively,

arising in the fragmentation of QCD jets in the decays of superheavy particles (top-down sources). Although the resulting neutrino spectra J_{ν_β} can be fairly well calculated for given injection spectra of protons or pions, the details are tedious and not worth the effort for our present purposes, and so we retain the simple parametrization (28) instead. The power law should mimic the various neutrino injection spectra reasonably well, with the spectral index in the range $\alpha \approx 1-2$. Throughout, we will take $E_{\text{min}}=0$ as a default value.

For astrophysical accelerators, one expects that E_{max} is a fraction ($\sim 5\%$) of the maximum proton energy, $E_{p \text{ max}}$. In the case of shock acceleration, $E_{p \text{ max}}$ is determined by the requirement that the gyromagnetic radius of the protons in the ambient magnetic field B be less than the spatial extension R of the accelerating source. The result is

$$E_{p \text{ max}}^{\text{shock}} \approx 10^{21} \text{ eV } (R/\text{kpc})(B/\text{mG}). \quad (29)$$

Even higher energies are possible in proposed nonshock mechanisms, such as unipolar induction, acceleration in strong electromagnetic waves in plasmas (wake fields) [64], or by magnetic recombination in the vicinity of massive black holes [65].

For top-down scenarios, E_{max} can be much larger, basically bounded only by the huge mass of the decaying particle or defect. These huge masses are thought to reflect the energy scale of an underlying phase transition. Popular examples include $M_{\text{GUT}} \sim 10^{16}$ GeV from grand unification, and $M_{\text{wimpzilla}} \sim 10^{11-13}$ GeV from the end of inflation. Neutrinos from the decay of these supermassive objects carry an energy about one order of magnitude less than the mass.

In general, there may be several classes of sources—such as GRB's, AGN's, and EHEC protons scattering inelastically off the CMB ("cosmogenic" neutrinos), and topological defects—which build up the total emissivity distribution. Predictions from a variety of such source classes are shown in Fig. 1 (bottom): from jets of AGN's [66], from ordinary topological defects ($M_X=10^{14}$ GeV) [67], from hidden-sector topological defects ($M_X=4 \times 10^{14}$ GeV) [68], and from the Z-burst scenario normalized to fit the highest energy cosmic ray anomaly [24] with different assumptions about the universal radio background. Also shown are the upper and lower bounds [46] (short-dashed-shaded) and one example [67] (short-dashed) of the cosmogenic neutrino flux. Finally, the cascade limit [69] from Ref. [70] (dotted-shaded) on transparent neutrino sources (discussed later) is shown. It applies to all scenarios where neutrinos originate from pion decays or even from electroweak jets [71]. These neutrinos are accompanied by photons and electrons that cascade down in energy during their propagation through the Universe. The cascade limit arises from the requirement that the associated diffuse gamma-ray fluxes should not exceed measurements [72].

It is straightforward to generalize Eq. (24) to a sum of source classes. In reality, it is more prudent to wish for even one class of sources at the energies of interest here, 10^{22-23} eV, and so we continue to work within the single source-class hypothesis.

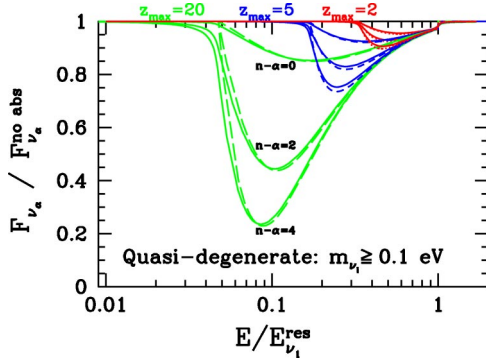


FIG. 4. Predicted flux (23) of neutrinos ν_α of flavor $\alpha=e, \mu, \tau$ at Earth, for a source emissivity characterized by a power-law activity (27) and a power-law injection spectrum (28), for the case of quasi-degenerate neutrino masses (22), normalized to the predicted flux for no absorption. E/E_ν^{res} scales as the degenerate mass m_ν . Three values of z_{max} are shown: 2 (dotted), 5 (short-dashed), and 20 (long-dashed). For each choice of z_{max} , three choices of $n-\alpha$ are shown: 0 (upper), 2 (middle), and 4 (lower). The corresponding solid lines show the same quantity evaluated with the complete energy dependence of the annihilation cross section from Ref. [3] arising from the finite Z width, instead of exploiting the zero-width approximation (13). For all curves, $E_{\text{max}} > E_{\nu_1}^{\text{res}}(1+z_{\text{max}})$ is assumed.

Another possibility, in principle, is to collect events from a small number of point sources, possibly just one. This presents the advantage that the spectra are not smeared by the integration over the redshift distribution. Since neutrinos are not deflected in flight, point-source selection is possible. However, we will find that the event numbers required for statistical significance is on the order of 100 or more. This large number presents a luminosity challenge for a small number of very bright sources [2]. In the end, it does not matter much for our purposes whether the EEC ν flux is diffuse or granular, since the isotropy of the relic background ensures nearly the same absorption shapes in the spectrum of either source type.

C. Case studies

Let us start our analyses with the quasidegenerate neutrino mass-spectrum case, whose implications are summarized in Eq. (22). In view of the expectation that EHEC ν neutrino fluxes are rapidly falling with energy (cf. Fig. 1), this case has the best prospects of observability, since it corresponds to the largest neutrino masses (cf. Fig. 3), and it brings all three resonance energies to a common lowest value (1). As discussed above [cf. Eq. (23)], we expect in this case an identical absorption dip in the neutrino flux of each flavor to be observed at Earth, if the flavor ratios at the sources were hadronlike (18) or democratic (20).

In Fig. 4, we show the predicted flux F_{ν_α} for a hadronic or flavor democratic source emissivity characterized by a power-law activity (27) with redshift evolution index n , and power-law injection spectrum (28) with index α . This flux is normalized to the predicted flux for no absorption, $F_{\nu_\alpha}^{\text{no abs}}$

[obtained by replacing the survival probability in Eq. (23) by unity]. A nice feature following from the two power laws, source activity and source emissivity, is that the normalized spectrum depends on n and α only through the combination $n-\alpha$; the source evolution and the energy fall-off compensate each other in a simple way. Three particular $n-\alpha$ combinations are shown in Fig. 4 for each fixed z_{max} : $n-\alpha=0$ (upper curves), $n-\alpha=2$ (middle curves), and $n-\alpha=4$ (lower curves).

Figure 4 illustrates some general features. Viewed as a function of decreasing energy, the absorption dip starts abruptly at the resonance energy $E_{\nu_i}^{\text{res}}$. This initial depth is determined by today's annihilation probability, $\approx 3\%$ according to (16), and reflects the absorptions occurring in nearby ($z\sim 0$) space. The dip extends in energy down to $E_{\nu_i}^{\text{res}}/(1+z_{\text{max}})$, with this minimum energy being the redshifted value of the resonant energy for annihilations occurring at cosmic distance ($z\sim z_{\text{max}}$). The overall depth of the dip increases dramatically with $n-\alpha$, thereby showing a strong preference for source evolution and/or flat energy spectra. With increasing $z\sim z_{\text{max}}=2, 5, 20$, the absorption depths are roughly 5%, 8%, and 15% for $n-\alpha=0$; 7%, 18%, and 55% for $n-\alpha=2$; and 10%, 27%, and 77% for $n-\alpha=4$, respectively. We observe in Fig. 4, as we did in Fig. 2 (top), that the replacement of the finite Z width with the δ -function approximation (13) is well justified for our purposes. From now on, we will exploit this simplification.

Next, we consider some nondegenerate neutrino mass scenarios, and source activities other than power-law. In the four panels of Figs. 5–8, we show the prediction (21) for $\Sigma_\alpha F_{\nu_\alpha} / \Sigma_\alpha F_{\nu_\alpha}^{\text{no abs}}$ for a neutrino spectrum which is quasidegenerate, i.e. lightest mass $m_{\nu_1}=0.4$ eV (upper panel), normal hierarchical, with $m_{\nu_1}=0.01$ eV (2nd panel) and $m_{\nu_1}=0.002$ eV (3rd panel), and inverted hierarchical with $m_{\nu_1}=0.002$ eV (bottom panel). Figures 5 and 6 display results for $z_{\text{max}}=10$, $\alpha=2$ (solid) and 0 (dashed), and the conservative and generous star-formation activities, respectively. Figure 7 shows results for a power-law activity, $n-\alpha=0$, and $z_{\text{max}}=10$ (short-dashed) and 20 (solid), mimicking a topological defect source scenario. Figure 8 displays results for a power-law activity, $n-\alpha=2$, and $z_{\text{max}}=2, 5, 10$ from upper to lower curves, corresponding to bottom-up acceleration. As expected, the depths of the dips increase markedly with z_{max} , lower spectral index α (28), and increased source evolution (27), the latter represented either by the shift from “conservative” to “generous” SFR as in Figs. 5 and 6, or by the increase in $n-\alpha$ as in Figs. 7 and 8.

The second panels in the figures illustrate that in the case of a normal hierarchy and an intermediate lightest neutrino mass of $m_{\nu_1}=0.01$ eV, corresponding to central values $m_{\nu_2}=0.013$ eV and $m_{\nu_3}=0.053$ eV in Fig. 3 (top), the absorption dips from ν_1 and ν_2 cannot be resolved. The two dips appear as one dip, which is understandably about twice as deep as the single dip from ν_3 .

The two lower panels present results for a very small,

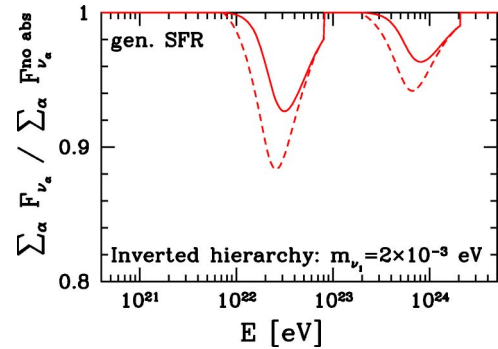
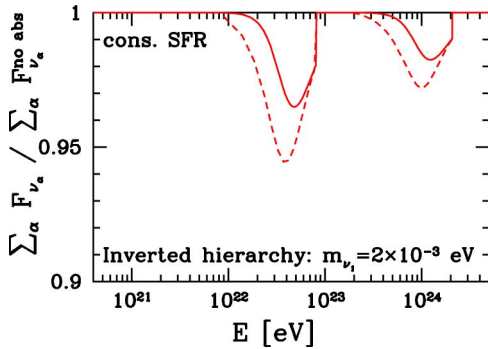
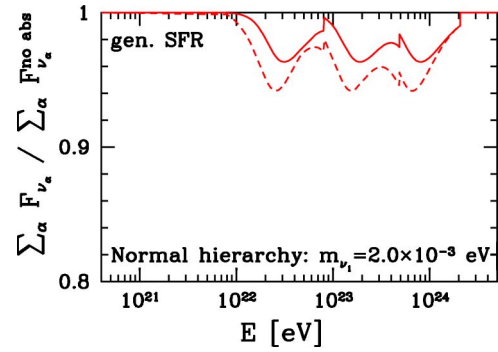
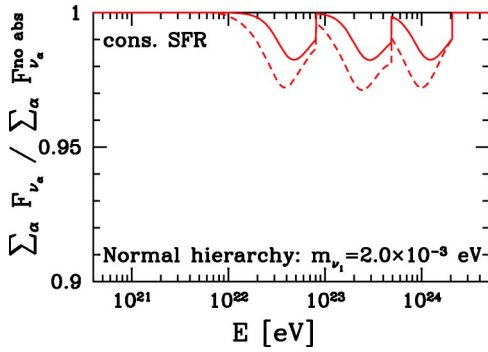
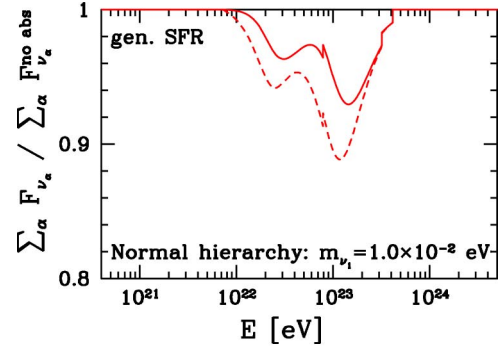
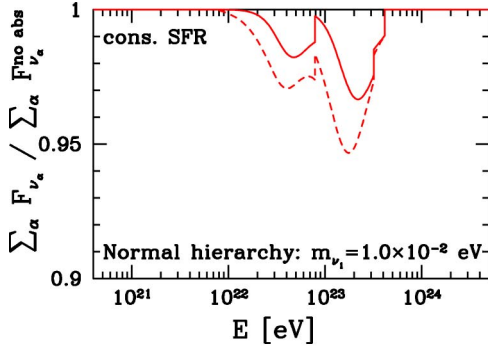
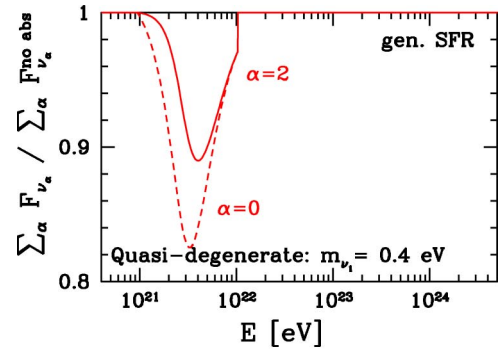
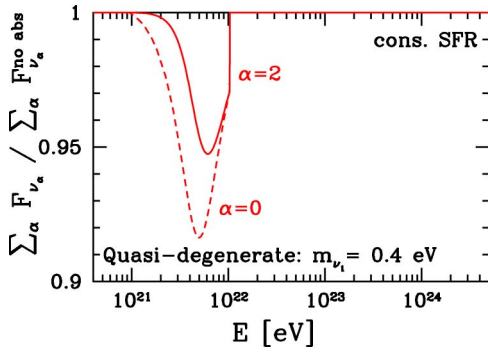


FIG. 5. Predicted flux of neutrinos summed over flavors at Earth (21), normalized to the predicted flux for no absorption, for a conservative SFR activity (25) to $z_{\text{max}} = 10$, injection-spectrum indices $\alpha=2$ (solid) and $\alpha=0$ (dashed) (28), and neutrino spectra that are quasidegenerate (top), normal-hierarchical (second and third panels), and inverted-hierarchical (bottom panel). For all curves it is assumed that $E_{\text{max}} > E_{\nu_1}^{\text{res}}(1 + z_{\text{max}})$.

FIG. 6. Same as Fig. 5, but with the generic SFR activity (25).

lightest neutrino mass of $m_{\nu_1} = 0.002 \text{ eV}$. In the normal hierarchical case (third panel), three absorption dips—associated with the masses $m_{\nu_1} = 0.002 \text{ eV}$, and, from the central values of Fig. 3 (top), $m_{\nu_2} = 0.0085 \text{ eV}$ and

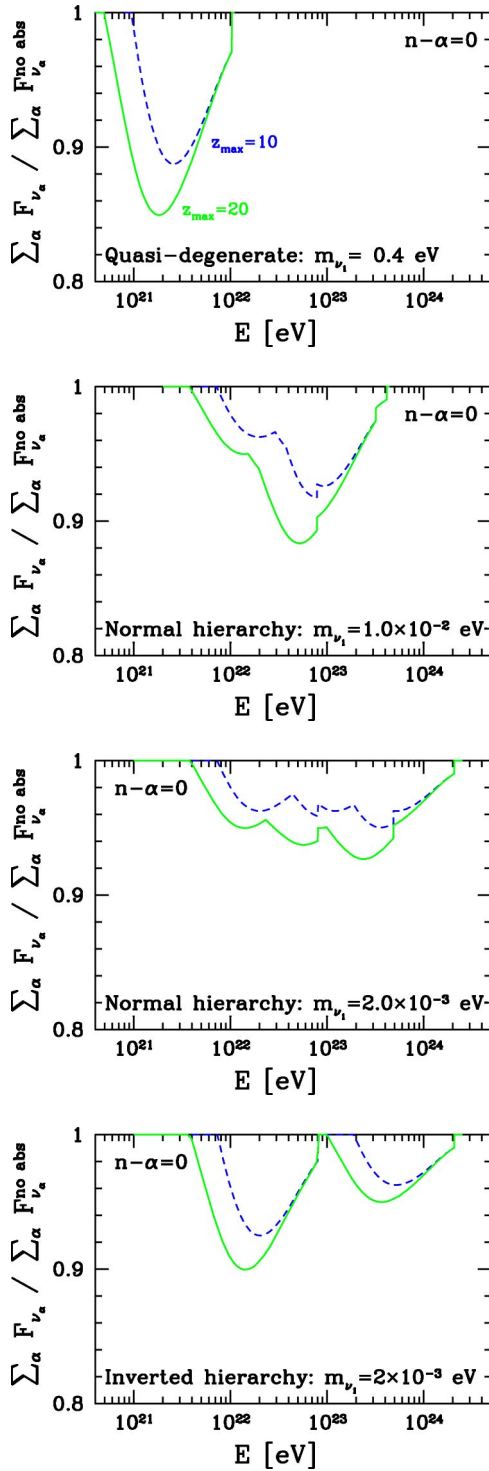


FIG. 7. As in Fig. 5, but with a power-law activity (27) and $n - \alpha = 0$, with $z_{\min} = 0$ and $z_{\max} = 10$ (short-dashed), 20 (solid), mimicking a topological defect source scenario. For all curves it is assumed that $E_{\max} > E_{\nu_1}^{\text{res}}(1 + z_{\max})$.

$m_{\nu_3} = 0.052$ eV—are clearly visible. These third panels present relic neutrino mass absorption spectroscopy at its best! For an inverted hierarchy, on the other hand, one finds, as illustrated in the bottom panels, that the absorption dips from ν_2 and ν_3 cannot be resolved individually. With m_{ν_1}

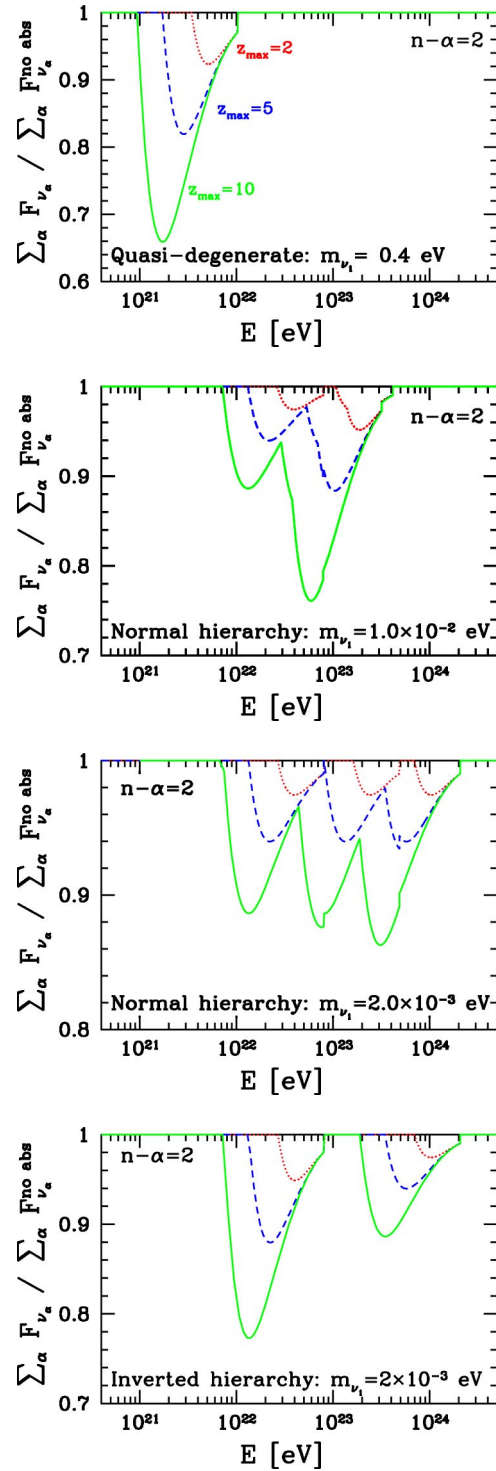


FIG. 8. Same as Fig. 7, but with $n - \alpha = 2$ fixed, and $z_{\max} = 2.5, 10$ (from upper to lower curves), corresponding to a bottom-up acceleration source scenario.

$= 0.002$ eV, the central values for the nearly degenerate heavy masses are $m_{\nu_2} \approx m_{\nu_3} = 0.052$ eV, from Fig. 3 (bottom).

The 2 meV mass is a critical one in the sense that the relic neutrino ν_1 is indeed nonrelativistic for z up to $z_{\max} = 2$, in which case Eq. (12) is applicable. For even smaller masses

TABLE I. Expected number of neutrinos (plus antineutrinos) to be detected in upcoming EHEC ν observatories with energies in the indicated intervals until the indicated year, for two different EHEC ν flux scenarios—one saturating the current observational upper bound and one saturating the cascade limit [cf. Fig. 1 (bottom)].

Energy decade	$\Sigma_{\alpha} \Delta (N_{\nu_{\alpha}} + N_{\bar{\nu}_{\alpha}})$					
	10 ^{21–22} eV		10 ^{22–23} eV		10 ^{23–24} eV	
Year	2008	2013	2008	2013	2008	2013
Observ. limit	240	700	30	90	2	5
Cascade limit	13	40	3	10	1	2

(or larger z_{\max}), the relic ν_1 density becomes relativistic and the absorption dip is known to wash out [1]. Thus, the two lower panels in Figs. 5–8 remain applicable for any $m_{\nu_1} < 0.002$ eV, if one simply removes the contribution of the highest-energy dip, at $\sim 10^{24}$ eV.

III. EXPERIMENTAL PROSPECTS

In the last paragraph we have seen that, depending on the source activity, injection spectrum, and neutrino masses, we may expect absorption dips of (10–20)% depth, located at $(0.1–0.5)E_{\nu_i}^{\text{res}}$, with a width of about an order of magnitude in energy. Is there any hope of discovery of these absorption dips in the next decade or beyond? The answer to this question depends critically on the magnitude of the EHEC ν flux at the resonance energies realized in Nature. In the following we will study a few benchmark flux scenarios (Sec. III A) and discuss further experimental (Sec. III B) issues concerning the clean reconstruction of absorption dips.

A. Benchmark flux scenarios

1. Most optimistic scenario: hidden sources

The most favorable case for relic neutrino absorption spectroscopy is the one in which the neutrino fluxes saturate the current observational upper bounds. It is known quantitatively [cf. the points marked “Z-burst” from Ref. [24] in Fig. 1 (bottom)] that in this most favorable flux scenario, the secondary fluxes of protons (and photons) from hadronic Z decay are remarkably of just the right order of magnitude to explain the highest energy cosmic rays above the GZK energy by Z bursts [20,22].

As can be seen from Fig. 1 (bottom) and as summarized in Table I, the upcoming EHEC ν observatories expect to see in this case in 2013 about 230 neutrinos (plus antineutrinos) per flavor $\alpha = e, \mu, \tau$ in the energy interval from 10²¹ eV to 10²² eV. The total number of neutrino events, then, is $N \approx 700$ neutrinos of all flavors, which implies a 1σ fluctuation of $\sqrt{N} \approx 26$. For a 3σ evidence for an absorption dip in this energy interval, an absorption depth of $3\sqrt{N}/N \approx 11\%$ is required; for a 5σ discovery, a depth of $5\sqrt{N}/N \approx 19\%$ is required. These statements have to be interpreted with caution, since they assume that the half-width of the dip roughly spends an energy decade.

As a comparison with our case studies in Sec. II C reveals, these depth requirements are achievable as long as the

dip is maximized via a quasidegenerate neutrino mass. The quasidegenerate condition is $m_{\nu_1} \gtrsim 0.1$ eV, which in turn implies that $E_{\nu_i}^{\text{res}} \lesssim 4 \times 10^{22}$ eV. For example, power-law source emissivities produce such depths for $n - \alpha \gtrsim 0$ and/or $z_{\max} \gtrsim 10$ [cf. Figs. 4, 7 (top), and 8 (top)]. Therefore, the realization of the Z -burst mechanism for the already observed highest energy cosmic ray events implies not only a discovery of the corresponding EHEC ν flux by 2008 (cf. Table I), but also the discovery of solid evidence for the associated absorption dip as soon as 2013.

What kind of source could produce such an optimistic EHEC ν flux, of the order of the observational limit at 10²¹ eV to 10²² eV? Bottom-up astrophysical Zevatrons may in principle produce such a flux, with the problem that they must accelerate protons to energies $E_{p \max} \gtrsim 10^{23}$ eV [24,25]. Much more favorable in this respect are top-down sources, involving physics beyond the standard model (SM), which naturally produce extremely high energetic particles. With either bottom-up or top-down, the sources must be “hidden” [73], in order to avoid the cascade limit [cf. Fig. 1 (bottom)]. Hidden sources are, by definition, sources from which neither nucleons nor photons (with energies $\gtrsim 100$ MeV) escape.

One popular hidden top-down example is a topological defect, here taken with mass $M_X \sim 4 \times 10^{14}$ GeV, which couples to SM particles only indirectly through a non-SM sector [52]. In this case of topological defect origin, we have $n - \alpha \approx 0$ and $z_{\max} \gtrsim 1$, and so Fig. 7 applies; the dip presented there is $\sim 15\%$ for the most interesting case of quasidegenerate neutrinos, $m_{\nu_1} \gtrsim 0.1$ eV. In Fig. 9, we show the significant wiggle from quasidegenerate neutrinos in the otherwise power-law spectrum expected for the hidden topological defect neutrino flux as in Fig. 1 (bottom). By design, this flux scratches the observational upper bound. Such a flux is also sufficient to explain the EHEC ν ’s via the Z -burst mechanism. The indicated error bars show the statistical significance that is expected with planned and proposed experiments by the year 2013 [cf. Fig. 1 (bottom)]. This example illustrates concretely our previous claim that substantial evidence for a relic neutrino absorption dip and the existence of the $C\nu B$ seems achievable within the foreseeable future, only if Z bursts are the dominant mechanism for the EHEC ν ’s. This requires speculative hidden sources.

For a hierarchical neutrino spectrum, on the other hand, with $m_{\nu_1} \lesssim 0.04$ eV, the lowest resonant energy is $\sim 10^{23}$ eV.

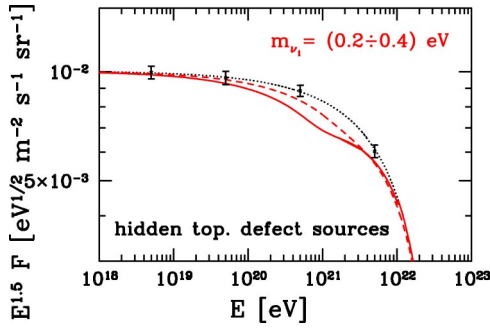


FIG. 9. Predicted neutrino flux at Earth, summed over all flavors, from a powerlike source emissivity, with $n=1.5$, $z_{\max}=\infty$, $\alpha=1.5$, and $E_{\max}=4 \times 10^{22}$ eV. This flux mimics one from hidden-sector topological defects with $M_X=4 \times 10^{14}$ GeV [cf. Fig. 1 (bottom)] and is also sufficient to explain the EHECR's above E_{GZK} via the Z-burst mechanism. Curves are without (dotted) and with relic neutrino absorption. Assumed neutrino masses are degenerate at $m_{\nu_1}=0.2$ eV (dashed) and $m_{\nu_1}=0.4$ eV (solid). The error bars indicate the statistical accuracy achievable per energy decade by the year 2013, for a flux which saturates today's observational bound from Fig. 1 (bottom).

Even if the EHEC ν flux saturates the current observational limit, the expected event numbers in the 10^{22-23} eV energy interval are apparently too small (cf. Table I) to allow a discovery of absorption dips in the next decade: with 90 expected events in this energy range by 2013, one needs absorption depths of $\sim 53\%$ ($\sim 32\%$) for a 5σ discovery (3σ evidence). As can be gleaned from Fig. 7, such large dips are not expected.

2. Less optimistic scenario: transparent sources

Let us turn our attention now to the non-optimized EHEC ν fluxes. From a neutrino flux that saturates the cascade limit [cf. Fig. 1 (bottom)], we may expect, in the 10^{21-22} eV energy bin, just 40 events by the year 2013 (cf. Table I). This yields a 3σ evidence for an absorption dip only if it has a depth of $\sim 48\%$ [74], again assuming a half-width of roughly one energy decade. Dips of this depth we have seen in Sec. II C only for quasidegenerate neutrinos and extreme parameter choices for the source activity. For example, for power indices $n-\alpha \gtrsim 2$, as might happen with highly evolving bottom-up sources, one needs $z_{\max} \gtrsim 20$ in addition [cf. Figs. 4 and 8 (top)]. The latter is easily achieved in top-down sources ($z_{\max}=\infty$); however, their restriction to $n-\alpha \approx 0$ tends to decrease the depth of the dip [cf. Figs. 4 and 7 (top)]. This pessimistic outlook is, however, ameliorated when one realizes that an increase in statistics by a factor of 10 reduces the required absorption depth by a factor of ~ 3 , to $\sim 15\%$ for a 3σ evidence, and to $\sim 25\%$ for a 5σ discovery. Such an increase in statistics could be achieved, for example, by undertaking more ANITA flights, by extending the EUSO flight time, or by developing the OWL or SalSA experiments. If the ~ 13 neutrinos in the 10^{21-22} eV interval (cf. Table I), expected from a flux which saturates the cascade limit, are measured by the pre-2008 experiments, then such extensions of the upcoming EHEC ν observatories would be warranted.

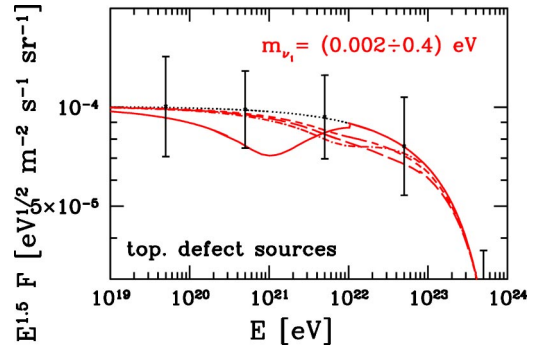


FIG. 10. As in Fig. 9, but with $E_{\max}=10^{24}$ eV, to mimic topological defects with $M_X=10^{16}$ GeV [cf. Fig. 1 (bottom)]. The assumed neutrino spectra are (i) quasidegenerate, $m_{\nu_1}=0.4$ eV (solid), (ii) normal hierarchical, $m_{\nu_1}=0.01$ eV (long-dashed) and $m_{\nu_1}=0.002$ eV (short-dashed), and (iii) inverted hierarchical, $m_{\nu_1}=0.002$ eV (dashed-dotted). Here the error bars indicate the statistical accuracy achievable per energy decade by the year 2013, for a flux which saturates today's cascade limit from Fig. 1 (bottom).

What classes of sources could deliver EHEC ν fluxes that saturate the cascade limit? It has recently been pointed out that cosmogenic neutrino fluxes from photopion production by cosmic protons on the CMB can reach this limit [46,67]. However, the limit is scratched typically just above E_{GZK} , at around 10^{20} eV and not higher [cf. Fig. 1 (bottom)]. Other possible sources include bottom-up Zevatrons with a large $E_{p \max} \gtrsim 10^{23}$ eV, and topological defects with $M_X \gtrsim 10^{14}$ GeV.

As examples, we show in Figs. 10 and 11, for various neutrino mass spectra, the predicted modulations in the otherwise power-law spectra expected for a topological defect and an astrophysical neutrino flux, respectively, both saturating the cascade limit. As expected, we see that a further increase in statistics by a factor of $\sim 10-100$, to reduce the error bars by a factor of $\sim 3-10$, will be mandatory to establish absorption dips. Moreover, a quasidegenerate neutrino mass spectrum is required.

B. Further experimental issues

To temper optimism, it is fair to mention three further issues that mitigate an experiment's ability to cleanly ob-

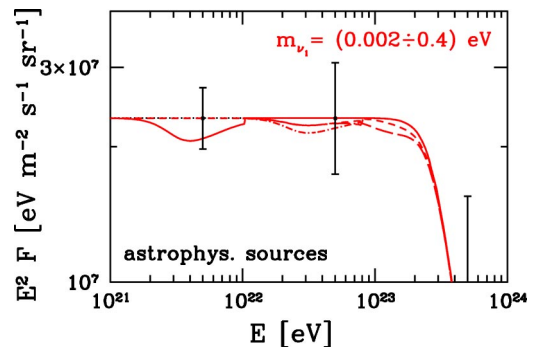


FIG. 11. As in Fig. 10, but with a generous SFR activity (25), $z_{\max}=20$, injection spectrum index $\alpha=2$, and $E_{\max}=10^{24}$ eV, to mimic astrophysical sources with $E_{p \max}=2 \times 10^{16}$ GeV.

serve an absorption dip. These are (i) the resolution with which an experiment can reconstruct the initial neutrino energy from the visible event energy, (ii) the flavor-dependent nature of energy reconstruction, and (iii) the possible confusion of a deviation of the (assumed) power-law spectrum due to absorption with other possible origins of deviation. We discuss each of these in turn.

Proposed EHEC ν detectors will measure shower energy from ground-based scintillator or water, or atmospheric nitrogen fluorescence, or radio signals in ice. From EHECR physics, energy reconstruction for hadron-initiated showers is known to be about 25%. Neutrino observatories expect to reconstruct shower energies with about the same error. This will smear the dip somewhat, but the smearing is tolerable because it is \leq the dip's half-width.

However, there is a subtlety connected with the various showers that result from different neutrino flavors [75]. Consider first neutral current (NC) events, identical for all neutrino flavors. On average, the final state neutrino carries away 80% of the incident energy, leaving 20% in the detected jet. The NC cross section is $\sim 45\%$ that of the charged-current (CC) cross section. Next consider CC interactions. The ν_μ and ν_τ CC events produce charged leptons, generally unobservable, which carry away 80% of the energy. Thus, the ν_μ and ν_τ CC events leave the same 20% energy deposition in the shower as do the NC interactions. For all these events, the absorption dip in the shower events will appear lower in energy than the true dip by a factor ~ 0.2 . This is easily corrected. However, there remains some smearing due to the event-by-event variance of the energy transfer about the mean (y distribution). We do not include the effect of this variance in the present work.

In contrast to the above event classes, a ν_e CC event will produce a hadronic jet plus an electron which creates an electromagnetic jet, and so the interaction deposits 100% of the incident energy into the combined shower.

In summary of this flavor discussion, NC and ν_μ and ν_τ CC scattering constitute about 77% of the events. The showers from these events contain about 20% of the incident neutrino energy. The other 23% of events are due to ν_e CC, with all of the incident energy observed. Thus, the observed events will look like a superposition of two fluxes with relative weights 23% and 77%, the latter displaced downward in energy by a factor of five in the mean, but with fluctuations.

Finally we come to the confidence issue, whether an observed feature near the end of the EHEC ν spectrum can be claimed to be an absorption dip. For example, the spectral end point may have structure simply due to differences among the individual contributing sources. To better ensure that a dip feature is observed, the continuation of the spectrum *above* the dip region should be observed. This requires more events, at higher energies still. Looking again at Table I, one sees even more reason to doubt the observation of an absorption feature, except for relatively low resonance energies. In turn, this means that quasidegenerate neutrinos, with mass not much below present cosmological bounds, are required if the relic neutrinos are to be detected via an absorption dip.

IV. SUMMARY AND CONCLUSIONS

Relic neutrino absorption spectroscopy via the observation of absorption dips in the EHEC ν flux may be feasible. The Z dips, if observed, are rich in particle and astrophysical information. The position of the high-energy edge of each dip is fixed by the neutrino mass. The depth and shape are determined by the density of the C ν B, by the usual cosmological parameters, by the activity and injection spectrum of the EHEC ν sources, and by the neutrino mass pattern. Moreover, as we have seen, the width of the dip region reflects the spectrum and evolutionary history of the neutrino source(s). Thus, were we so fortunate as to resolve a dip in some detail, the information to be mined truly belongs to interdisciplinary particle-astrophysics.

However, the statistics at the dip is entirely determined by the magnitude of the EHEC ν flux. Moreover, the energy positions of the dips—and, to some extent, their depths—critically depend on the neutrino mass spectrum.

Large event samples, $N \gg 100$, beyond 10^{21} eV are needed to reveal Z dips with statistical significance. To get these event numbers within the next decade or so, an EHEC ν flux at least as large as the present cascade limit is required. Almost certainly, even higher fluxes must be invoked. They could be generated by hidden sources, which are opaque to nucleons and high-energy photons, thereby evading the cascade limit. Such sources are likely indicative of new physics. Moreover, the high flux requirement implies a Z-burst contribution to the EHECR events beyond E_{GZK} within an order of magnitude of the present AGASA rate. The Auger project should detect these large EHEC ν and EHECR fluxes within a few years, and EUSO should easily measure them within the decade.

As inferred from our numerous figures, a quasidegenerate neutrino mass spectrum with $m_{\nu_1} \gtrsim 0.1$ eV seems to be required to produce a detectably deep absorption dip at a possibly accessible resonant energy. Such a spectrum is testable in several ways [9]. Neutrinoless double beta decay experiments (assuming neutrinos are Majorana), such as CUORE [76] and NEMO-3 [77], and the KATRIN tritium decay experiment [78] can be expected to show positive results already in the upcoming decade. The transfer functions relating CMB fluctuations to today's large-scale matter distributions can be expected to show positive contributions from massive neutrinos [79]. It seems sensible to say that, if pre-2008 experiments do not see any EHEC ν flux in the 10^{21-23} eV region, then, in the context of the concordance cosmological model, absorption dips will not be observed within the next decade or two. If such is the case, then Nature will have overlooked a wonderful opportunity to produce direct evidence for the C ν B [80].

In summary, the presently planned neutrino detectors open up a window of opportunity for relic neutrino absorption spectroscopy. The next decade will be really exciting and decisive in this respect.

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APPENDIX: IMPLEMENTING NEUTRINO FLAVOR PHYSICS

At production in the cosmic sources, let the ratios of the neutrino flavors be written as $\nu_e : \nu_\mu : \nu_\tau = w_e : w_\mu : w_\tau$, with $\sum w_\beta = 1$. A convenient description of the flavor mixture is given by the density matrix

$$\rho(t=0) = \sum_{\beta} w_{\beta} |\nu_{\beta}\rangle \langle \nu_{\beta}|. \quad (\text{A1})$$

The density matrix is properly normalized to $\text{Tr}(\rho(0)) = 1$, and so describes the ensemble-averaged, single neutrino.

The relic neutrinos, with thermal energies $\sim 3kT_\nu \sim 0.5$ meV in the today's epoch, long ago decohered into their mass-eigenstates [21]. Therefore, rewriting the density matrix in the mass basis allows us to simply include the losses due to resonant absorption. Making use of the mixing matrix notation $|\nu_{\beta}\rangle = U_{\beta j} |\nu_j\rangle$, or equivalently, $U_{\beta j} = \langle \nu_j | \nu_{\beta} \rangle = \langle \nu_{\beta} | \nu_j \rangle^*$, the forward-propagated density matrix in the mass basis, with resonant absorption, is

$$\begin{aligned} \rho(t) &= \sum_{\beta} w_{\beta} \sum_j [e^{-i(m_j^2/2E)t} e^{-(\Gamma_j/2)t} U_{\beta j} |\nu_j\rangle] \\ &\quad \times \sum_k [\langle \nu_k | U_{\beta k}^* e^{+i(m_k^2/2E)t} e^{-(\Gamma_k/2)t}] \\ &= \sum_{\beta} w_{\beta} \sum_{j,k} U_{\beta j} U_{\beta k}^* e^{+i\Delta m_{kj}^2 t} e^{-[(\Gamma_j + \Gamma_k)/2]t} |\nu_j\rangle \langle \nu_k|. \end{aligned} \quad (\text{A2})$$

Here, $\Gamma_j = cn_j \sigma_j$ is the annihilation rate, and the factor $e^{-(\Gamma/2)t}$ correctly accounts for absorption at the amplitude level [81]. As written, this formula does not include energy losses due to redshifting; however, it is straightforward to incorporate redshifting in the final expressions.

At Earth, the probability to detect flavor α is then,

$$\begin{aligned} P_{\nu_{\alpha} \text{ detected}} &= \langle \nu_{\alpha} | \rho(t) | \nu_{\alpha} \rangle = \sum_{\beta} w_{\beta} \sum_{j,k} U_{\beta j} U_{\beta k}^* U_{\alpha k} U_{\alpha j}^* \\ &\quad \times e^{-i\Delta m_{kj}^2 t} e^{-[(\Gamma_j + \Gamma_k)/2]t}. \end{aligned} \quad (\text{A3})$$

This probability expression naturally divides into a sum diagonal in the mass, and a sum of interfering mass terms from which oscillations arise. After a bit of algebra, the result is

$$\begin{aligned} P_{\nu_{\alpha} \text{ detected}} &= \sum_{\beta} w_{\beta} \sum_j |U_{\alpha j}|^2 |U_{\beta j}|^2 e^{-\Gamma_j t} \\ &\quad + 2 \sum_{\beta} w_{\beta} \sum_{j < k} e^{-[(\Gamma_j + \Gamma_k)/2]t} \\ &\quad \times \left[\Re(U_{\beta j} U_{\beta k}^* U_{\alpha k} U_{\alpha j}^*) \cos\left(\frac{\Delta m_{kj}^2 t}{2E}\right) \right. \\ &\quad \left. - \Im(U_{\beta j} U_{\beta k}^* U_{\alpha k} U_{\alpha j}^*) \sin\left(\frac{\Delta m_{kj}^2 t}{2E}\right) \right]. \end{aligned} \quad (\text{A4})$$

Setting the annihilation rates to zero returns the usual formula for neutrino oscillations (cf., e.g., Ref. [55]).

Although the neutrino wave function at extreme high energies may remain coherent over a cosmic distance, the oscillating terms are effectively averaged away in any realistic observation. Let us look first at the coherence of the neutrino wave function. The spread in the neutrino mass eigenstates after travel through a distance D results from the difference in the group velocities $\beta = \delta E / \delta p$:

$$\begin{aligned} D_{\text{spread}} &= ct \delta\beta = \frac{D \Delta m^2}{2E^2} \\ &= 6 \times 10^{-20} \left(\frac{0.7}{h}\right) \left(\frac{D}{D_H}\right) \Delta m_{-3}^2 E_{22}^{-2} \text{cm}, \end{aligned} \quad (\text{A5})$$

where $\Delta m_{-3}^2 \equiv \Delta m^2 / 10^{-3} \text{eV}^2$, $E_{22} \equiv E / 10^{22} \text{eV}$, and the Hubble distance $D_H \equiv cH_0^{-1} = 4.2 (0.7/h)$ Gpc characterizes the typical cosmic distance. Decoherence of the wave packet results when D_{spread} exceeds the natural length of the wavepacket, call it $c\tau_{\Psi}$, the latter being determined by conditions at the source [82]. The natural length may be c times the production time [83] (e.g., $c\tau_{\pi^\pm} \sim 3$ m), the mean free path (mfp) between interactions in a dense source [83], or the spatial uncertainty in the location of the production point in the source [84]. The decoherence length is obtained by setting D_{spread} in Eq. (A5) equal to $c\tau_{\Psi}$, and solving for D . The result, as a fraction of the Hubble size, is

$$\frac{D_{\text{decohere}}}{D_H} = 0.5 \left(\frac{h}{0.7}\right) \left(\frac{\tau_{\Psi}}{3\text{m}}\right) \frac{E_{22}^2}{\Delta m_{-3}^2} \times 10^{20}. \quad (\text{A6})$$

Clearly, decoherence does not occur in our Universe for a 10^{22}eV neutrino.

Now we turn to the averaging effects of measurement. To observe the oscillating terms requires measurement of the phase $\phi = \Delta m^2 t / 2E$ to better than 2π , or equivalently, knowledge of $\delta(D/E)$ to better than $4\pi / \Delta m^2$. This in turn requires knowledge of δD to better than $\lambda_{\text{osc}} \equiv 4\pi E / \Delta m^2$, and knowledge of δE to better than $E(\lambda_{\text{osc}} / D)$. With the oscillation length being so short $\sim E_{22} / \Delta m_{-3}^2$ kpc compared to cosmic scales, there is no hope to observe the oscillating term.

From here on, we may use for the detection probability just the averaged value of Eq. (A4), i.e. just the first term

$$P_{\nu_\alpha \text{ detected}} = \sum_j |U_{\alpha j}|^2 e^{-\Gamma_j t} \sum_\beta w_\beta |U_{\beta j}|^2. \quad (\text{A7})$$

Finally, taking into account energy losses due to redshift, Eq. (A7) is easily generalized to Eq. (9) in the main text for the differential propagation function of neutrinos, and to Eq. (17) for the neutrino flux to be observed at Earth.

As simple as Eq. (A7) is, further simplifications are possible. For example, if $w_e : w_\mu : w_\tau = 1:1:1$, as might arise from “democratic” neutrino emission from a topological defect, then by unitarity of the mixing matrix, $\sum_\beta w_\beta |U_{\beta j}|^2 = 1/3$ independently of j . It is also known [33] that the same result obtains for the case $w_e : w_\mu : w_\tau = 1:2:0$ as results from charged pion decay to neutrinos, in the limit of $\nu_\mu \leftrightarrow \nu_\tau$ interchange-symmetry. The latter is exact if the atmospheric mixing angle θ_{32} is maximal (45°) and $\Re(U_{e3}) = 0$. Oscillation data [5] show that these two conditions are satisfied (or nearly so). Since any linear combination of the two sets $\{w_\beta\}$ just discussed must also give the same conclusion, we arrive at a minitheorem: whenever $w_e = 1/3$ and $\nu_\mu \leftrightarrow \nu_\tau$ interchange symmetry is (nearly) valid, then $\sum_\beta w_\beta |U_{\beta j}|^2 = 1/3$ independently of j .

Because $w_e = 1/3$ includes two popular neutrino-production cases, charged pion decay and democratic emission, we will assume that $w_e = 1/3$ holds and work with the very compact result

$$P_{\nu_\alpha \text{ detected}} = \frac{1}{3} \sum_j |U_{\alpha j}|^2 e^{-\Gamma_j t}. \quad (\text{A8})$$

This expression leads immediately to the main text’s Eq. (19) for the neutrino flux at Earth, after again taking into account the energy loss due to redshifting.

Cosmic-neutrino detectors are unlikely to be capable of neutrino flavor identification (see Ref. [53] for a discussion). Accordingly, it is tempting to also sum over α , the flavor arriving at Earth, to arrive at the especially simple result

$$P_{\nu_{\text{any}}} = \sum_\alpha P_{\nu_\alpha \text{ detected}} = \frac{1}{3} \sum_j e^{-\Gamma_j t}, \quad (\text{A9})$$

leading to Eq. (21). Caution is warranted, however, as the interactions of the different flavors deposit different mean energies into the detector. For example, the short mfp of a high-energy electron and the long mfp’s of high-energy muons and taus imply that in charged-current interactions, ν_μ ’s and ν_τ ’s will leave much less visible energy in the detector than will ν_e ’s. This issue is discussed further in the main text in Sec. III.

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- [62] The contribution from small z in Eq. (21) is negligible, for any reasonable activity. We note that a “cosmological” distance of 50 Mpc corresponds to the small value $z=0.012$, for our default cosmological parameters.
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