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Minimal cosmogenic neutrinos

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The observed flux of ultrahigh energy (UHE) cosmic rays (CRs) guarantees the presence of high-energy cosmogenic neutrinos that are produced via photohadronic interactions of CRs propagating through intergalactic space. This flux of neutrinos doesn't share the many uncertainties associated with the environment of the yet unknown CR sources. Cosmogenic neutrinos have nevertheless a strong model dependence associated with the chemical composition, source distribution or evolution and maximal injection energy of UHE CRs. We discuss a lower limit on the cosmogenic neutrino spectrum which depends on the observed UHE CR spectrum and composition and relates directly to experimentally observable and model-independent quantities. We show explicit limits for conservative assumptions about the source evolution.

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I. INTRODUCTION

Cosmogenic neutrinos are produced when ultrahigh energy (UHE) cosmic rays (CRs) interact with the cosmic radiation background while propagating between their sources and Earth. The frequent interactions with the cosmic microwave background (CMB) limits the propagation of nucleons with energies greater than $E_{\rm GZK} \simeq 40$ EeV to within a few 100 Mpc and is responsible for the so-called Greisen-Zatsepin-Kuzmin (GZK) cutoff of extragalactic protons [1,2]. Mesons produced in these interactions quickly decay and produce an observable flux of cosmogenic (or GZK) neutrinos [3]. In fact, the observed spectrum of CRs extending up to energies of a few 100 EeV shows a suppression above $\sim E_{\rm GZK}$ with high statistical significance [4,5]. This could be an indication that protons are dominating the flux at these energies. In this case the flux of cosmogenic neutrinos is typically large.

However, the experimental situation is less clear. Measurements of the elongation rate distribution of UHE CR showers indicate a transition of their arrival composition from light to heavy within 4–40 EeV [6,7]. If a heavy component dominates also at higher energies the prospect for cosmogenic neutrino production is "disappointing" [8] or at least less favorable than for the proton scenario [9,10]. A crucial uncertainty of this scenario is the maximal injection energy of the nucleus with mass number A. Rigidity scaling of the energy cutoff in CR sources allows for larger maximal energies for the case of nuclei [11]; as long as $E_{\rm max} \gg AE_{\rm GZK}$, even this scenario will produce an appreciable amount of cosmogenic neutrinos [12]. If this condition is not met interactions with the subdominant cosmic photon background from the optical/infrared will still contribute to the cosmogenic neutrino flux. We will use the estimate of Ref. [13] for our calculation.

The IceCube neutrino observatory has reached the sensitivity for the detection of optimistic cosmogenic neutrino fluxes [14]. In the case of a non-observation it is of interest to know a lower limit on the various source emission possibilities for their definite exclusion. Lower cosmogenic neutrino flux limits have already been discussed in the context of proton-dominated scenarios via a deconvolution of early Auger data [15]. We will discuss in this article updates of these lower limits and extensions to more general assumptions for the source distribution and chemical composition. Similar to Ref. [15] we will not attempt to construct a specific source emission model that fits the Auger spectrum and elongation rate distribution but we will derive the limits directly from the observed composition measurement and spectrum. From this we can derive a strict lower limit on the cosmogenic flux.

II. COSMIC RAY PROPAGATION

The propagation of UHE CR nuclei is affected by photodisintegration [16–18], photohadronic interactions [19], Bethe-Heitler pair production [20] and redshift losses due to the expansion of the Universe. It is convenient to consider a homogenous and isotropic distribution of CR sources and derive the observed CR from the comoving number density $Y_i \equiv n_i/(1+z)^3$ as a solution to a set of Boltzmann equations [21],

$$\dot{Y}_i = \partial_E(HEY_i) + \partial_E(b_iY_i) - \Gamma_i^{\rm tot}Y_i + \sum_j \int {\rm d}E_j\gamma_{ji}Y_j + \mathcal{L}_i. \quad (1)$$

The cosmic expansion rate $H(z)$ follows the usual "concordance model" dominated by a cosmological constant with $\Omega_\Lambda \sim 0.73$ and a (cold) matter component, $\Omega_{\rm m} \sim 0.27$ with $H^2(z) = H_0^2[\Omega_{\rm m}(1+z)^3 + \Omega_\Lambda]$, normalized to its value today of $H_0 \sim 72$ km s^{-1} Mpc^{-1} [22]. The first and second terms on the r.h.s. of Eq. (1) describe, respectively, redshift and other continuous energy losses with rate $b \equiv -{\rm d}E/{\rm d}t$. In the following we will treat

Bethe-Heitler pair production as a continuous energy losses process [20]. The third and fourth terms describe more general interactions involving particle losses ($i \to$ anything) with total interaction rate $\Gamma_i^{\rm tot}$, and particle generation of the form $j \to i$ with differential interaction rate γ_{ij}. The last term on the r.h.s., $\mathcal{L}_i$, corresponds to the emission rate density of CRs of type i per comoving volume. The detailed description of the interaction rates and their scaling with redshift has been discussed in our previous publications [21,23].

We first discuss the case of proton sources. The flux of cosmogenic neutrinos today ($z = 0$) depends on the comoving number density of protons at all redshifts and can be approximated as [21]

$$J_\nu(E_\nu) \simeq \frac{1}{4\pi}\int_0^\infty \frac{dz'}{H(z')} \times \int d\mathcal{E}_p \gamma_{p\nu}(z', \mathcal{E}_p, (1+z')E_\nu) Y_p(z', \mathcal{E}_p), \quad (2)$$

where $\mathcal{E}_p$ is the solution to the differential equation $\dot{\mathcal{E}}_p = -H\mathcal{E}_p - b_{\rm BH}(z, \mathcal{E}_p)$ with initial condition $\mathcal{E}_p(0, E_p) = E_p$. The comoving number density of protons can be written as

$$Y_p(z, \mathcal{E}_p(z)) \simeq \frac{1}{1+z}\int_z^\infty \frac{dz'}{H(z')}\mathcal{L}_{p,\rm eff}(z', \mathcal{E}_p(z')) \times \exp\left[\int_z^{z'} dz'' \frac{\partial_E b_{\rm BH}(z'', \mathcal{E}_p(z'')) - \Gamma(z'', \mathcal{E}_p(z''))}{(1+z'')H(z'')}\right], \quad (3)$$

where the effective source term is defined as

$$\mathcal{L}_{p,\rm eff}(z, E_p) = \mathcal{L}_p(z, E_p) + \int d\mathcal{E}_p \gamma_{pp}(z, \mathcal{E}_p, E_p) Y_p(z, \mathcal{E}_p). \quad (4)$$

III. MINIMAL NEUTRINOS FROM PROTONS

A minimal contribution to the flux of cosmogenic neutrinos can be estimated as follows. As a first step we approximate the UHE CR spectrum measured by Auger via the phenomenological fit given in Ref. [24]. This fit is shown in Fig. 1 as a dashed-dotted line together with recent data of Auger, HiRes [4] and the Telescope Array [25] (TA). We only consider CRs above the ankle feature, i.e., above an energy of 4 EeV. Extragalactic contributions that extend below the ankle would contribute on top of the lower neutrino limits presented here. Note, that the normalization of the Auger data is lower by a about a factor two than HiRes and TA and hence cosmogenic neutrinos derived from this data are the lowest.

Whereas the spectrum of UHE CRs is dominated by close-by sources, the neutrino flux receives contributions up to the Hubble scale. The overall flux will hence increase for an increasing number of sources with redshift. We assume that redshift evolution decouples from the source emission spectrum, i.e., $\mathcal{L}_p(z, E) = \mathcal{H}(z)Q_p(E)$ and we consider two scenarios for the source evolution $\mathcal{H}(z)$. In the most conservative case we assume source contributions within redshift $z_{\rm max} = 2$ with no source evolution, i.e., $\mathcal{H}_0 = \Theta(z_{\rm max} - z)$. This corresponds to the case that CR emission rate starts rapidly at high redshift and remains constant (per comoving volume) throughout the rest of the Universe's history. However, most of the UHE CR candidate source are associated with a stellar or quasistellar origin. Hence a more realistic scenario assumes a source evolution following the star formation rate (SFR)[26,27]. We will use the estimate [28,29]

$$\mathcal{H}_{\rm SFR}(z) = \begin{cases} (1+z)^{3.4} & z < 1, \\ N_1(1+z)^{-0.3} & 1 < z < 4, \\ N_1N_4(1+z)^{-3.5} & z > 4, \end{cases} \quad (5)$$

with normalization factors, $N_1 = 2^{3.7}$ and $N_4 = 5^{3.2}$. Since we assume conservative choices of the source evolution the associated cosmogenic neutrino flux can be regarded as *lower limits* on the expected cosmogenic neutrino flux.

In the following we will derive approximate solutions to Eqs. (2) and (3) using an iterative scheme. For the iteration start we choose $Q_p^{(0)}(E_p) = (H_0 + \partial_E b_0 + \Gamma_0)4\pi J_{\rm CR}^{\rm obs}(E_p)$, where b_0 and Γ_0 are the energy loss and interaction rate, respectively, at redshift $z = 0$. The iteration step is then given by

$$Q_p^{(n+1)}(E_p) = 4\pi J_{\rm CR}^{\rm obs}(E_p)/\eta^{(n)}(E_p), \quad (6)$$

where we use the phenomenological fit of Ref. [24] for $J_{\rm CR}^{\rm obs}(E)$ and introduce the effective survival distance

$$\eta^{(n)}(E_p) = \int_0^\infty \frac{dz'}{H(z')} \frac{\mathcal{L}_{p,\rm eff}^{(n)}(z', \mathcal{E}_p(z'))}{Q_p^{(n)}(E_p)} \times \exp\left[\int_0^{z'} dz'' \frac{\partial_E b_{\rm BH}(z'', \mathcal{E}_p(z'')) - \Gamma(z'', \mathcal{E}_p(z''))}{(1+z'')H(z'')}\right]. \quad (7)$$

We continue this iteration until the relative correction $\sum_i (Q_{{\rm p},i}^{(n+1)}/Q_{{\rm p},i}^{(n)} - 1)^2$ stops to decrease or a maximal (sufficiently large) iteration step is achieved. This compensates for numerical instabilities.

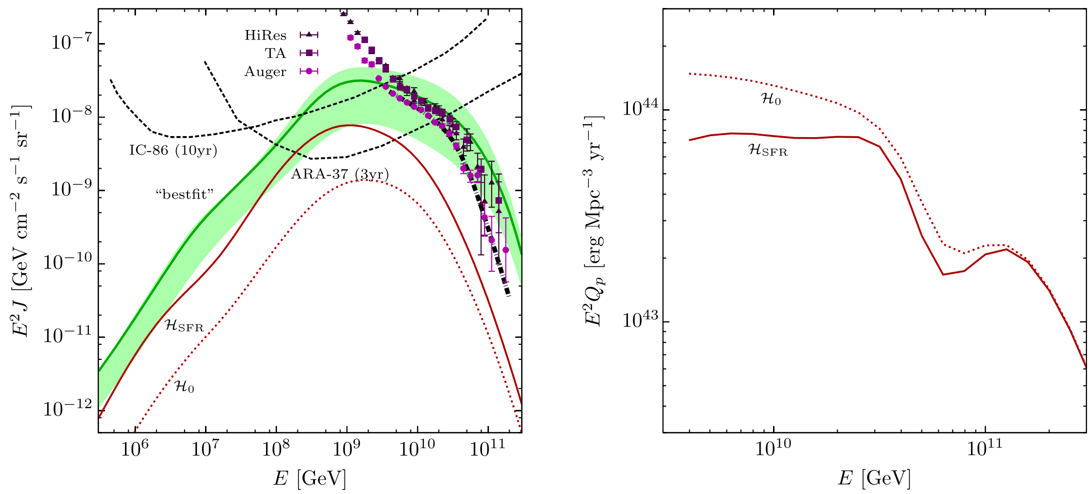

FIG. 1 (color online). Left-hand panel: minimal flux of cosmogenic neutrinos assuming dominance of protons above 4 EeV. We show the results without source evolution (dotted line) and assuming source evolution according to the star formation rate (solid line). Also shown are the projected sensitivities of IceCube (10 years) and the ARA-37 (3 years) as dashed lines. The thick dashed-dotted line shows the approximation of the Auger spectrum above the ankle. For comparison, we also show the best-fit cosmogenic neutrino flux (green solid line) from Ref. [40] ($E_{\rm min} = 10^{18.5}$ eV) including the 99% C.L. (green shaded area) obtained by a fit to the HiRes spectrum. Right-hand panel: the minimal proton emission rate density derived from the iteration method explained in the main text.

In Fig. 1 we show the resulting cosmogenic neutrino flux (left-hand plot) and the corresponding proton emission rate density (right-hand plot) for this procedure for the two evolution scenarios. The limit for the SFR evolution agrees well with that derived from a deconvolution analysis in Ref. [15]. We also indicate in this plot the sensitivity of IceCube [14] and the proposed Askaryan Radio Array (ARA) [30]. Three years of observation with the 37 station configuration of ARA ("ARA-37") is sufficient to reach the proton emission model for the SFR case. In the case of no source evolution this scenario is reached after ten years. As we already emphasized, this result depends on the absolute normalization and/or energy calibration of the observed UHE CR spectrum. For a normalization to HiRes and TA data we expect our limits to scale up by about a factor 2.

IV. GENERALIZATION TO HEAVY NUCLEI

The case of a more general scenario including UHE CR sources of heavy nuclei is more complicated. The chemical composition observed at Earth is the result of rapid photodisintegration in the radiation background and there is no simple connection to the source composition. However, since photodisintegration conserves the energy per nucleon we can derive a lower neutrino limit by tracking the leading (heaviest) nucleus back to its source starting from a composition A_o and Z_o inferred from UHE CR observations.

The parent nuclei during this back-tracking are at least as heavy as the observed mass composition. For instance, a single helium nucleus in the observed spectrum might be produced via the production chain $^{10}\mathrm{B} \rightarrow {}^{9}\mathrm{Be}(+p) \rightarrow {}^{4}\mathrm{He}(+{}^{4}\mathrm{He} + p)$ from the source. The parent nuclei in each step of this chain determine the interaction and energy loss rates during propagation. For a lower limit on the cosmogenic neutrino flux we have to *minimize* the emission rate density of the UHE CR nuclei associated with their cascades in the CMB. This corresponds to a *maximal* survival probability of nucleons. Hence, we can derive a strict lower limit with the assumption that the back-tracking of the nuclei is indefinite, i.e., we assume no upper limit on the atomic mass number in the nuclei cascades.

It is important to realize how this apparently *unphysical* approximation enables us to set a lower bound on the cosmogenic neutrino flux. In contrast to the pure-proton model discussed in the previous section the lower bounds on cosmogenic neutrinos from heavy nuclei do *not* correspond to a flux prediction. Clearly–on astrophysical grounds–we are not expecting that elements heavier than iron should play a significant role in the sources of CRs or at any point in the nuclei cascades developing during propagation. However, as we are only interested in lower bounds on the flux we are free to relax any astrophysical requirement on the CR sources *as long as* the corresponding neutrino yield *decreases*. Limiting the maximal nucleon mass in the back-tracking method would require that the sources have to become more luminous to compensate for photodisintegration losses and would *increase* the lower limit on cosmogenic neutrinos.

Photodisintegration that drives the cascades competes with photohadronic interactions and Bethe-Heitler energy loss. To first order, a photohadronic interaction of the nucleon with energy E, charge Z and mass number A can be approximated via the interaction rate of the free proton as $\Gamma_{A\gamma}(E) \simeq A\Gamma_{p\gamma}(E/A)$ [18]. Hence, the interaction rate

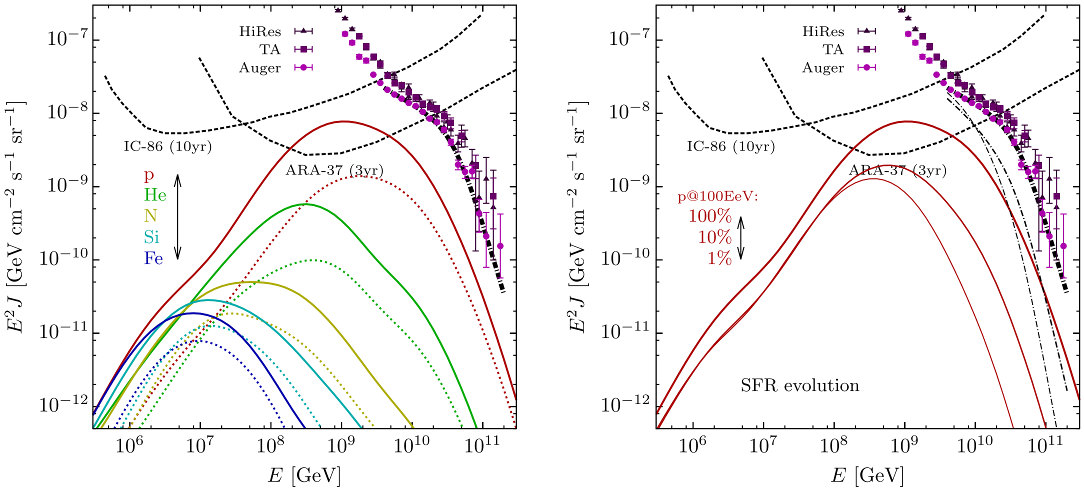

FIG. 2 (color online). Minimal flux of cosmogenic neutrinos for a mixed composition. Left-hand panel: minimal flux of cosmogenic neutrinos assuming dominance of protons, helium, nitrogen, silicon or iron (upper to lower lines) in UHE CRs above 4 EeV. We show the results without source evolution (dotted line) and assuming source evolution according to the star formation rate (solid line). Right-hand panel: the contribution of protons (red lines) in a mixed composition scenario assuming 100% (upper line), 10% (middle line) and 1% (lower line) proton contribution (black dashed-dotted lines) at 100 EeV.

per nucleon of the parent nucleus is approximately the same.[1] Energy loss via Bethe-Heitler pair production, however, scales as $b_{A\gamma}(E) \simeq Z^2 b_{p\gamma}(E/A)$ and the effective energy loss per nucleon scales as Z^2/A. Again, for a maximal survival probability of the nucleons and hence a minimal emission rate density of the sources, we assume a minimal Bethe-Heitler energy loss of the nucleons. This corresponds to the energy loss of a nucleus with charge Z_o and atomic mass number A_o associated with the observed composition.

In summary, a lower limit on the cosmogenic neutrino flux can hence be derived by the same Eqs. (2) and (3) where we now replace the continuous energy loss by its minimal contribution $b_{\rm min}(z, E) \simeq (Z_o^2/A_o) b_{\rm BH}(z, E)$, where $b_{\rm BH}$ correspond to the energy loss of a free proton. The photohadronic interaction of the nucleons is given by the average interaction of protons and neutrons. The total number of nucleons per nucleon energy depends on the observed (or inferred) mass composition of UHE CRs. Assuming a single component we have the relation $E_N J_N(E_N) = A_o E_{\rm CR} J_{\rm CR}(E_{\rm CR})$ with $E_N = E_{\rm CR}/A_o$ or $J_N(E_N) = A_o^2 J_{\rm CR}(E_{\rm CR})$.

In the left-hand panel of Fig. 2 we show the minimal cosmogenic neutrino fluxes for the case of helium, nitrogen, silicon and iron dominance of the Auger spectrum. The level of these fluxes is not in reach of present or future neutrino observatories. However, cosmogenic neutrino fluxes strongly depend on the maximal injection energy of the sources. We conservatively assume for our method that the maximal energy does not exceed the observed energy of UHE CRs. However, it is in principle possible that these models produce detectable fluxes of cosmogenic neutrinos [12] if the maximal energy significantly exceeds $A \times E_{\rm GZK}$. We will briefly discuss this in the following section.

We can also generalize our method to the case of a mixed compositions, which is indicated by the Auger CR elongation rate distribution. For instance, if $f_i(E_{\rm CR})$ denotes the fraction of nuclei with mass A_i at CR energies $E_{\rm CR}$ the mean mass number is given by

$$J_N(E_N) \simeq \sum A_i^2 f_i(A_i E_N) J_{\rm CR}(A_i E_{\rm N}). \quad (8)$$

Hence the minimal cosmogenic neutrino flux in this case is $J_\nu^{\rm min}(E_\nu) = \sum_i J_i^{\rm min}(E_\nu)$, where the individual $J_i^{\rm min}$ are derived in the same way as before but using $f_i(E_{\rm CR}) J_{\rm CR}(E_{\rm CR})$ as the input spectrum. As an example we show in the right-hand panel of Fig. 2 the lower limit associated with protons in a multi-component model, where we decrease the proton contribution at 100 EeV to 10% ($\alpha = 1$) and 1% ($\alpha = 2$) using $f_p = 1 - (1 + (E/10^{19}\ {\rm eV})^{-\alpha})^{-1}$ with $f_{\rm A} = 1 - f_p$.

V. OPTIMISTIC COSMOGENIC NEUTRINOS

Before we conclude we would like to take a more "optimistic" point of view and consider the opposite question how large cosmogenic neutrino fluxes of heavy CR nuclei can become. Predictions of the cosmogenic neutrino spectra are very sensitive to the maximal energy of UHE CR nuclei. Hence, optimistic neutrino predictions assume that the maximal energy of CR nucleons is much larger than the GZK cutoff, i.e., $E_{\rm CR}/A \gg E_{\rm GZK}$.

[1] Coherent photonucleus interactions that can dominate the interaction rate for higher photon energies are known to shadow this linear mass scaling [31]. However, their contribution to the interaction rate is subdominant after folding with a smooth background photon spectrum.

For the discussion it is convenient to introduce the energy density ($\mathrm{eV\,cm^{-3}}$) of the GZK neutrino background at redshift z defined as

$$\omega_{\rm GZK} \equiv \int \mathrm{d}E_\nu E_\nu Y_\nu(E_\nu). \tag{9}$$

From the Boltzmann equations (1) we can derive the evolution of the energy density as

$$\dot{\omega}_{\rm GZK} + H\omega_{\rm GZK} = \sum_i \int \mathrm{d}E b_{i,\rm GZK}(z, E) Y_i(z, E), \tag{10}$$

where $b_{i,\rm GZK}(E) \simeq 0.2E\Gamma_{\gamma\pi}(E/A_i)$ is an approximation of the energy loss of the nuclei into GZK neutrinos [12].

The UHE CR interactions with background photons are rapid compared to cosmic time scales. The energy threshold of these processes scale with redshift z as $A_iE_{\rm th}/(1+z)$ where $E_{\rm th} \gtrsim E_{\rm GZK}$ is the (effective) threshold today. We can therefore approximate the evolution of the energy density as

$$\dot{\omega}_{\rm GZK} + H\omega_{\rm GZK} \sim \frac{3K_\pi \mathcal{H}(z)}{4(1+K_\pi)} \sum_i \int_{A_iE_{\rm th}/(1+z)} \mathrm{d}E E Q_i(E), \tag{11}$$

where K_π is the ratio of charged to neutral pions produced in $p\gamma$ interactions. Assuming a power-law emission rate density $Q_i(E) \propto E^{-\gamma_i}$ with sufficiently large cutoff $E_{\rm max} \gg E_{\rm th}$ we see that cosmic evolution enhances the GZK flux as

$$\omega_{\rm GZK} \sim \frac{3}{8}\sum_i \eta_i \frac{(A_iE_{\rm th})^2 Q_i(A_iE_{\rm th})}{\gamma_i - 2}, \tag{12}$$

where the last term assumes $\gamma_i > 2$ and the effective survival distance of the nucleons is defined as

$$\eta_i = \int_0^\infty \frac{\mathrm{d}z}{H(z)} \mathcal{H}(z)(1+z)^{\gamma_i - 4}. \tag{13}$$

For $\gamma_i \simeq 2$ and for those evolution scenarios $\mathcal{H}$ that we have considered so far in this article, the effective survival distances range from $0.48/H_0$ (no evolution) to $2.4/H_0$ (SFR). This agrees well with the relative ratio ~ 5 of the energy densities associated with lower neutrino limits in the proton-dominated scenario shown in Fig. 1.

The relation (12) shows that as long as the maximal energy per nucleon is much larger than the pion production threshold in the CMB (i.e., $E_{\rm max} \gg AE_{\rm GZK}$) and the injection index is $\gamma_i \simeq 2$ the main difference in the energy density of GZK neutrinos comes from the underlying evolution model, not by the inclusion of heavy elements. In principle, this factor can be large even for heavy nuclei if the sources have a strong evolution. The fact that typical CR models including heavy nuclei produce significantly less GZK neutrinos can be traced back to a low maximal energy per nucleon and/or a weak evolution of CR sources [9,10,12]. Note that the latter is an important ingredient of proton-dominated low-crossover models [32], whereas CR models of heavy nuclei including more model degrees of freedom are less predictable with respect to the source evolution.

Note that, ultimately, the inferred energy density ω_γ of the extragalactic diffuse γ-ray background in the GeV–TeV region constitutes an upper limit for the total electromagnetic energy from pion-production of UHE CR nuclei [33–35]. An upper limit is given via the relation

$$\omega_\gamma \gtrsim \left(\frac{1}{3} + \frac{4}{3K_\pi}\right)\omega_{\rm GZK}. \tag{14}$$

Recent result from Fermi-LAT [36] translates into an energy density of $\omega_\gamma \simeq 6 \times 10^{-7}\ \mathrm{eV/cm^3}$ [37]. Assuming an E^{-2} neutrino spectrum between energies E_- and E_+ a numerical simulation gives a *cascade limit* of [38]

$$E^2 J^{\rm cas}_{\rm all\,\nu}(E) \simeq \frac{3\times 10^{-7}}{\log_{10}(E_+/E_-)}\ \mathrm{GeV\,cm^{-2}\,s^{-1}\,sr^{-1}}. \tag{15}$$

This is only slightly lower than the estimate (14) for $K_\pi = 1$. Cosmogenic neutrino fluxes that saturate this bound in the EeV region are already ruled out by IceCube upper limits [14]. In fact, unidentified γ-ray source like BL Lacs could significantly contribute to the extragalactic γ-ray background [39]. For this case the cascade bound (15) can be considered conservative.

VI. DISCUSSION

We have discussed in this article a simple procedure to derive lower limits on the cosmogenic neutrino flux. The limits are based on the observed spectrum and composition of UHE CRs and depend on the unknown evolution of sources. For the case of a proton-dominance in the UHE CR data we show that ARA-37 should identify the flux of cosmogenic neutrinos after 3 years of observation if UHE CR sources follow the star formation rate. For the less optimistic (and less realistic) case of no source evolution it would require 10 years of observation.

In the case of heavy nucleus dominance of the CR flux cosmogenic neutrino predictions are less optimistic. We can derive a lower limit in this scenario by tracking the leading nucleus back to its source. Since photodisintegration conserves the energy per nucleon of the interaction we can base our analysis on the observed number of nucleons in UHE CRs, which depends on the observed mass composition.

The dominant contribution to the cosmogenic neutrino flux is expected from the proton content in the UHE CR spectrum. We show in Fig. 2 two cases where we decrease the contribution of protons to 10 and 1% at 100 EeV and assume source evolution with the star-formation rate. Even this less optimistic case is in reach of ARA-37 after 5 years of observation.

The prediction of cosmogenic neutrinos is very sensitive to the maximal CR injection energy per nucleon. If this is significantly larger than the GZK cutoff, even UHE CR scenarios dominated by heavy nuclei can produce large fluxes of cosmogenic neutrinos. For flat spectra that are sufficiently close to E^{-2} the energy density of these optimistic GZK neutrino predictions depends on the cosmic evolution of the sources.

All cosmogenic neutrino fluxes shown in this analysis are normalized to Auger data. The spectra observed with HiRes and the Telescope are in general larger, which could be a result of an overall systematic energy shift by 20–30%. This corresponds to an upward shift of up to a factor 2 of the energy density $E_{CR}^2 J_{CR}(E_{CR})$. Hence the lower limits shown in Figs. 1 and 2 should be similarly scaled upward.

Finally, we would like to stress that the present analysis does not take into account statistical uncertainties of the CR data. However, the method can be easily extended to this case. In Refs. [40,41] it was shown that an actual fit to HiRes data assuming a proton power-law injection in the sources is statistically consistent with cosmogenic neutrino fluxes that exceed the minimal bound by up to an order of magnitude and are in reach of the IceCube detector.

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