Cosmic ray bound for models of extragalactic neutrino production

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We obtain the maximum diffuse neutrino intensity predicted by hadronic photoproduction models of the type which have been applied to the jets of active galactic nuclei or gamma ray bursts. For this, we compare the proton and gamma ray fluxes associated with hadronic photoproduction in extragalactic neutrino sources with the present experimental upper limit on cosmic ray protons and the extragalactic gamma ray background, employing a transport calculation of energetic protons traversing cosmic photon backgrounds. We take into account the effects of the photon spectral shape in the sources on the photoproduction process, cosmological source evolution, the optical depth for cosmic ray ejection, and discuss the possible effects of magnetic fields in the vicinity of the sources. For photohadronic neutrino sources which are optically thin to the emission of neutrons we find that the cosmic ray flux imposes a stronger bound than the extragalactic gamma ray background in the energy range between 10^5 GeV and 10^{11} GeV, as previously noted by Waxman and Bahcall [Phys. Rev. D **59**, 023002 (1999)]. We also determine the maximum contribution from the jets of active galactic nuclei, using constraints set to their neutron opacity by gamma ray observations. This present upper limit is consistent with the jets of active galactic nuclei producing the extragalactic gamma ray background hadronically, but we point out future observations in the GeV-to-TeV regime could lower this limit. We also briefly discuss the contribution of gamma ray bursts to ultrahigh-energy cosmic rays as it can be inferred from possible observations or limits on their correlated neutrino fluxes.

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I. INTRODUCTION

The connection between the emission of cosmic rays, gamma rays, and neutrinos from astrophysical accelerators is of considerable interest for the solution of the problem of the origin of cosmic rays [1,2]. The reason why a fundamental relation between these components must exist can be understood as follows. Particle acceleration mechanisms in cosmic plasmas generally require the presence of a magnetic field which is able to confine the accelerated charged particles, i.e., electrons and protons (or ions). The accelerated electrons lose their energy quickly in synchrotron radiation in the magnetic field. These synchrotron photons provide a target for accelerated protons to undergo photohadronic interactions, resulting in the production of mesons, which decay. The particles which eventually emerge from this process are high-energy photons, electrons (pairs), neutrons, and neutrinos. Neutrinos are directly ejected due to their low interaction cross section. Gamma rays and secondary electrons initiate electromagnetic cascades, shifting the power from ultrahigh energies to energies below which the absorption of gamma rays by pair production is unimportant [3]. Finally, the neutrons, which, unlike the protons, are not confined in the magnetic field, can escape and convert into cosmic ray protons after β decay, but their flux may be diminished by photoinduced $n \rightarrow p$ conversions. The branching ratios which distribute the available energy into the different channels are thereby generally of order unity. This leads to the conclusion that *cosmic proton accelerators produce cosmic rays, gamma rays, and neutrinos with comparable luminosities* [4].

The fundamental relation between cosmic ray and gamma ray production has the obvious consequence that active galactic nuclei (AGN), which are known to produce a large fraction of the gamma rays in the Universe, are a prime candidate for the sources of ultrahigh-energy cosmic rays (UHECR) [4]. The spectra of the emitted GeV-TeV gamma radiation of AGN also agree with the predictions of a hadronic production of these gamma rays [5]. Other prominent gamma ray sources, in particular the violent events connected with gamma ray bursts (GRB), have also been suggested as UHECR source candidates. Moreover, most of the extragalactic gamma ray energy is found in a diffuse background rather than in point sources, which allows for the possibility that the UHECR sources could be relatively large objects which would have a low gamma ray surface brightness, such as radio galaxies [6], galaxy clusters [7,8], or even larger structures [9]. Whatever the sources are, the fundamental relation between gamma ray and neutrino fluxes implies that, if in fact the extragalactic gamma ray emission is due to hadronic processes, a neutrino flux of a similar bolo-

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metric luminosity must exist. This prediction is the major motivation for high-energy neutrino experiments, which are currently operated, under construction, or planned. In fact, most model predictions for extragalactic high-energy neutrino fluxes have been made by using the source model to determine the spectral shape, and then by normalizing the total flux to some fraction of the diffuse extragalactic gamma ray background (EGRB) [10,11].

In contrast to the limits set by gamma ray observations, the limits which could arise from the corresponding cosmic ray emission of the neutrino sources have been given little attention. A detailed treatment of this problem regarding predictions for neutrino fluxes from the decay of topological defects (TD) has been given by Protheroe and Stanev [12], and a brief discussion of the possible relevance for diffuse neutrino fluxes from AGN by Mannheim [10]. Recently, it was proposed by Waxman and Bahcall [1] that indeed the measured flux of ultrahigh-energy cosmic rays provides the most restrictive limit on extragalactic diffuse neutrino fluxes for a broad class of sources. They claim that this cosmic ray bound is for neutrinos of all energies two orders of magnitude lower than the bound previously used which was based on the EGRB. In addition to the obvious restriction of their result to neutrinos from proton accelerators (i.e., excluding TD models), their claim is mainly based on three assumptions: (i) neutrons produced in photohadronic interactions can escape freely from the source, (ii) magnetic fields in the Universe do not affect the observed flux of extragalactic cosmic rays, and (iii) the overall injection spectrum of extragalactic cosmic rays is $\propto E^{-2}$. A key role is played by their assumption (iii): By assuming a specific cosmic ray input spectrum, they can normalize their bound at the ultrahigh energies, where they argue that (ii) also applies. Assumption (i) is justified by showing that some particular sources of specific interest, like the TeV-blazar Mrk 501, or also GRB, are transparent to the emission of neutrons. The authors claim that this new bound set by cosmic ray data essentially rules out the hypothesis that hadronic processes in AGN jets can produce the EGRB, and consequently that their neutrino fluxes are overestimated.

The purpose of the present paper is to reexamine the role cosmic ray observations can play to constrain models of neutrino production. The paper is organized as follows. In Sec. II, we give a brief review on the properties of photohadronic interactions, and derive the production spectra for cosmic rays and neutrinos for power-law photon target spectra. In Sec. III we briefly describe the effect of the propagation of extragalactic cosmic rays. We then follow Waxman and Bahcall in deriving a cosmic ray bound on neutrino fluxes, adopting their assumptions (i) and (ii), but instead of assuming a specific cosmic ray injection spectrum, we assume a specific spectrum for the *observable* extragalactic cosmic ray flux, which is constructed such that it complies with all existing observational limits on the cosmic ray proton intensity. In Sec. IV we turn to AGN, and discuss in particular the photohadronic opacity of blazar jets as can be estimated from observations. We shall show that most AGN in fact have large photohadronic opacities at ultrahigh energies, and we derive an upper bound for the neutrino contribution from AGN considering the effect of neutron opacity. In Sec. V, we discuss the possible effect on cosmic ray protons from neutron decay of the magnetic fields known to exist in clusters of galaxies, and radio galaxies which are considered the hosts of gamma ray emitting AGN. We derive critical energies below which they could increase the bound. We conclude by discussing the combined effect of our results, and to what extent the cosmic ray data can indeed constrain models for expected neutrino fluxes, and vice versa.

II. COSMIC RAY, GAMMA RAY, AND NEUTRINO EMISSION FROM EXTRAGALACTIC PROTON ACCELERATORS

In this section, we obtain the form of the spectra of cosmic rays, gamma rays, and neutrinos escaping from cosmic proton accelerators. Here, we shall assume that protons are confined within the acceleration region whereas photoproduced neutrons may escape to become cosmic rays. This is justified in particular for cosmic ray sources connected to relativistic outflows, like AGN jets or GRB, since the lifetime of protons is here limited by adiabatic energy losses and generally is much shorter than the diffusive escape time. For example, if protons are accelerated near the beginning of an AGN jet, say where the jet width is $\sim 10^{16}$ cm and where the gamma rays are probably produced, and are released near the end of the jet after its width has expanded to at least a few parsec, then their energies on release will be down by a least two orders of magnitude. As a consequence, only protons resulting from the decay of neutrons which have escaped from the acceleration region contribute significantly to the cosmic ray spectrum. Escaping neutrons are produced by accelerated protons which interact with ambient soft photons, together with pions which decay into neutrinos.

We shall first discuss the properties of photohadronic interactions. We shall then obtain the ambient proton spectrum in the emission region resulting from shock acceleration followed by radiative cooling, adiabatic losses, and advection away from the shocked region where the photon density is sufficient for efficient photoproduction of neutrons. From this we obtain the form of the spectra of neutrons and neutrinos on production, and the escaping neutron spectrum which may be modified by neutron absorption in photohadronic $n \rightarrow p$ conversions.

A. Photohadronic interactions

Photohadronic interactions can be divided into two processes: photoproduction of pions (and other mesons), and Bethe-Heitler production of $e^{\pm}$ pairs. Charged pions decay as $\pi^{\pm} \rightarrow \mu^{\pm} \nu_{\mu}$, $\mu^{\pm} \rightarrow e^{\pm} \nu_{\mu} \nu_e$ (here and in the following we disregard the difference between neutrinos and antineutrinos), neutral pions decay into gamma rays as $\pi^0 \rightarrow \gamma\gamma$. Electrons and positrons (from pion decay and Bethe-Heitler production) cascade in the magnetic field and radiation field, and so can be assumed to convert all their energy into synchrotron radiation in the magnetic field required for the acceleration of protons. The production of charged pions allows the production of secondary neutrons through isospin exchange.

The physics of photohadronic interactions in ambient photon spectra has been extensively studied in Monte Carlo simulations [13,14]. The properties of the production of secondary particles can be expressed in the fraction ξ of the proton energy given to a specific particle component per interaction. For neutrinos, gamma rays, and neutrons, the values

$$\xi_\nu \approx \xi_\gamma \approx 0.1, \tag{1}$$

$$\xi_n \approx 0.5, \tag{2}$$

respectively, have been found for power-law target spectra typical in AGN jets, while for GRB target spectra the values are $\xi_\nu \approx \xi_\gamma \approx \xi_n \approx 0.2$ [15]. The energy per particle in units of the proton energy have been found for neutrinos and neutrons as $\langle E_\nu \rangle / E_p \approx 0.033$ and $\langle E_n \rangle / E_p \approx 0.83$, respectively, while for GRB they are $\langle E_\nu \rangle / E_p \sim 0.02$ and $\langle E_n \rangle / E_p \sim 0.5$, respectively [15]. From this we can immediately define the relative energy of escaping neutrinos and neutrons as

$$\eta_{\nu n} = \langle E_\nu \rangle / \langle E_n \rangle \approx 0.04 \tag{3}$$

for both AGN and GRB target spectra. The fractional energy loss of the proton per interaction is $\kappa_p \approx 0.2$ in the AGN case, and $\kappa_p \approx 0.5$ for GRB [15]. Note that the quantities given here as typical for GRB apply only at ultrahigh proton energies, at lower energies they approach the values found for AGN [15].

Electromagnetic radiation initiated by the Bethe-Heitler pair production, and by photons and electrons from neutral and charged pion decay, are reprocessed in synchrotron pair cascades. This energy will emerge as a component of the gamma ray background radiation, for AGN mainly in an energy range 10 MeV–1 TeV. The contribution of the Bethe-Heitler process to the production of gamma rays depends on the target energy spectrum index, α, since its cross section peaks at energies about two orders of magnitude lower than that of photopion production. Assuming that the power law extends over this range without change, one can find the relation

$$L_\gamma = [1 + \exp(5\alpha - 5)] L_\nu, \tag{4}$$

where L_γ and L_ν are the bolometric photohadronic luminosities in gamma rays and neutrinos [16], and we note that for $\alpha = 1$, $L_\gamma = 2L_\nu$. We also note that in general, $L_\gamma \geqslant L_\nu$ holds as a direct consequence of the isospin symmetry of charged and neutral pions—hence, for any kind of neutrino production involving pion decay, the bolometric flux in correlated photons sets a robust upper limit on the possible bolometric neutrino flux.

B. Ambient proton spectrum

We assume a spectrum of protons on acceleration of the form, $Q(E_p) \propto E_p^{-2} \exp(-E_p/E_{\max})$ [s^{-1} GeV^{-1}]. In order to calculate the spectra of cosmic ray protons and neutrinos escaping from AGN jets we first need to obtain the ambient spectrum of protons, $N_p(E_p)$ [GeV^{-1}], in the shocked regions where the soft photon target density is sufficiently high for pion photoproduction to take place. From the ambient proton spectrum we can then obtain the spectra of neutrinos and neutrons produced, $Q_n(E_n)$ and $Q_\nu(E_\nu)$ [s^{-1} GeV^{-1}], respectively. Then taking account of the optical depth of the emission region for neutrons escaping photoinduced $n \rightarrow p$ conversions, one can obtain from the neutron production spectrum, the spectrum of cosmic ray protons resulting from the β decay of the neutrons escaping from the jet, $Q_{\rm cr}(E_p)$.

The spectrum of protons on acceleration and the ambient proton spectrum are related by the proton loss time scale $t_p(E_p)$,

$$N_p(E_p) \sim Q_p(E_p) t_p(E_p). \tag{5}$$

The processes mainly contributing to losses of protons are interactions with radiation, advection away from the shock region of dimension R, and adiabatic energy losses if the emission region expands.

The advection time scale is expected to be $t_{\rm adv} \sim R/(\beta_{\rm sh} c)$ where $\beta_{\rm sh}$ is the shock velocity and R is the dimension of the jet in the shocked region. In relativistic flows streaming away from a central source, this energy independent time scale is usually in competition with adiabatic energy losses of the protons due to the expansion of the flow [16]. In a relativistic outflow, characterized by a bulk Lorentz factor Γ and an opening angle Θ, the expansion velocity in the comoving frame (in units of c) is $\beta_{\rm ex} \approx \Gamma\Theta$ for $\Theta < \Gamma^{-1}$, and $\beta_{\rm ex} \approx 1$ otherwise. This leads to adiabatic cooling on a time scale $t_{\rm ad} \approx R/(\beta_{\rm ex} c)$. For example, in GRB one can generally assume that $\beta_{\rm ex} \simeq 1$, and also observations of superluminal motions in AGN jets are consistent with $\Theta \sim \Gamma^{-1}$ [17], thus $\beta_{\rm ex} \sim 1$. In the following, we shall assume that adiabatic losses are relevant with $0.3 \lesssim \beta_{\rm ex} \lesssim 1$. This has the important consequence that the lifetime of protons in the jet (or outflow) is limited to about one crossing time. The time scale for diffusive escape of protons is usually much longer (except, maybe, near the maximum proton energy), thus protons with $E_p \ll E_{\max}$ can be assumed to be confined in the emitter and do not contribute to the cosmic ray emission.

The photon target spectrum will be assumed to have a power-law shape $n(\epsilon) \propto \epsilon^{-\alpha-1}$ extending to energies sufficiently above the threshold for photopion production by protons of energy E_p. Then, for $\alpha > 0$ the time scale for energy loss by photohadronic interactions is asymptotically of the form

$$t_{p\gamma}(E_p) \propto E_p^{-\alpha}, \tag{6}$$

where $t_{p\gamma}$ is understood as including Bethe-Heitler and pion production losses, the cooling time for pion production will be called $t_{p,\pi} > t_{p\gamma}$. For very flat target spectra, as, for example, in GRB at ultrahigh proton energies, the photoproduction time scale is approximately constant [16], thus, Eq. (6) applies with $\alpha = 0$. We shall confine the discussion to the values $\alpha = 1$, relevant in AGN, and $\alpha = 0$ hereafter. Hence, for $\alpha = 1$ we obtain

$$t_{p\gamma}(E_p) = (E_1/E_p)(R/c), \tag{7}$$

with $t_{p,\pi}(E_p)\approx 1.5 t_{p\gamma}(E_p)$. E_1 is the energy corresponding to unit optical depth for photoproduction losses, $\tau_{p\gamma}(E_1) = R/[c t_{p\gamma}(E_1)] = 1$. Setting $t_p^{-1} = t_{p,ad}^{-1} + t_{p\gamma}^{-1}$ we obtain

$$N_p(E_p) = Q_p(E_p)(R/c)[(E_p/E_1) + a]^{-1} \tag{8}$$

with $a = \max(\beta_{\rm ex}, \beta_{\rm sh})$, and for $Q_p(E_p) \propto E^{-2}$ we have

$$N_p(E_p) \propto \begin{cases} a^{-1}E_p^{-2} & (E_p < aE_1), \\ E_1 E_p^{-3} & (E_p > aE_1). \end{cases} \tag{9}$$

(Note that for clarity, here and in the next section we omit the exponential cutoff in the spectrum at $E_{\rm max}$.) Obviously, for $\alpha = 0$ the optical depth for photoproduction is constant, and $N_p(E_p) \propto Q_p(E_p)$. One can show that for typical photon densities in GRB fireballs $\tau_{p\gamma} < 1$, and that Bethe-Heitler losses are unimportant, viz., $t_{p,\pi} \approx t_{p\gamma}$ [18,16].

C. Generic cosmic ray proton and neutrino production spectra

The time scale for photohadronic production of neutrons is $t_{p\gamma\to n} \approx t_{p,\pi}\kappa_p/\langle N_n\rangle$. For $\alpha = 1$, this is $t_{p\gamma\to n} \approx 0.5 t_{p\gamma} \propto E_p^{-1}$, while for $\alpha = 0$ we have $t_{p\gamma\to n} \approx 2.5 t_{p\gamma}$. This immediately gives the production spectrum of neutrons,

$$Q_n(E_n) \approx N_p(E_n)/t_{p\gamma\to n}(E_n) \propto \begin{cases} a^{-1}E_n^{-1} & (E_n < aE_1), \\ E_1 E_n^{-2} & (E_n > aE_1). \end{cases} \tag{10}$$

Neutrons may escape to become cosmic ray protons. However, because neutrons themselves suffer pion photoproduction losses, the cosmic ray production spectrum will differ from $Q_n(E_p)$ above the energy bE_1 at which the optical depth for neutron escape, $\tau_{n\gamma}$, is one. Neutrons can be considered as ''absorbed'' after they are converted into a proton, or after they have lost most of their energy in $n\gamma$ interactions, whichever time scale is shorter. For $\alpha = 1$ this means $\tau_{n\gamma} \approx 2\tau_{p\gamma}$, giving $b \approx 0.5$ for AGN jets, while for $\alpha = 0$ and typical GRB photon densities, $\tau_{n\gamma} \approx \tau_{p\gamma} < 1$, which means that neutron absorption is unimportant in GRB.

We note that in a homogeneous spherical medium of radius R the optical depth decreases radially $\propto (1 - r/R)$ from its central value τ to zero at $r = R$ giving rise to the geometrical escape probability of an interacting particle propagating in a straight line

$$\mathcal{P}_{\rm esc}(\tau) \approx (1 - e^{-\tau})/\tau \approx \begin{cases} 1 & \tau < 1, \\ \tau^{-1} & \tau > 1, \end{cases} \tag{11}$$

resulting in the cosmic ray proton production spectrum being steepened above bE_1, compared with $Q_n(E_p)$. For $\alpha = 1$, the cosmic ray proton production spectrum is therefore

$$Q_{\rm cr}(E_p) \propto \begin{cases} a^{-1}E_p^{-1}E_1^{-1} & (E_p < aE_1), \\ E_p^{-2} & (aE_1 < E_p < bE_1), \\ bE_1 E_p^{-3} & (bE_1 < E_p). \end{cases} \tag{12}$$

However, because the two break energies aE_1 and bE_1 are probably very close, $a \sim b$, we shall adopt a single break energy $E_b = bE_1$, and use the following approximations:

$$Q_n(E_n, L_p) \propto L_p \exp\left[\frac{-E_n}{E_{\rm max}}\right] \begin{cases} E_n^{-1}E_b^{-1} & (E_n < E_b), \\ E_n^{-2} & (E_b < E_n), \end{cases} \tag{13}$$

$$Q_{\rm cr}(E_p, L_p) \propto L_p \exp\left[\frac{-E_p}{E_{\rm max}}\right] \begin{cases} E_p^{-1}E_b^{-1} & (E_p < E_b), \\ E_p^{-3}E_b & (E_b < E_p). \end{cases} \tag{14}$$

Note that we have now put in explicitly the cutoff in the accelerated proton spectrum, and the proportionality with the proton luminosity of the source, L_p. In Sec. IV, we shall relate L_p to the observed photon luminosities for specific models of neutrino emission by AGN jets. Obviously, in GRB we simply have $Q_{\rm cr}(E_n, L_p) \approx Q_n(E_n, L_p) \propto L_p E_p^{-2} \exp(-E_p/E_{\rm max})$.

The production spectrum of muon neutrinos will have the same broken power-law form as the neutron production spectrum, and is related to it by

$$Q_{\nu_\mu}(E) = \frac{2}{3}\frac{\langle \xi_\nu \rangle}{\langle \xi_n \rangle \eta_{\nu n}^2} Q_n(E/\eta_{\nu n}), \tag{15}$$

where we count ν_μ and $\overline{\nu_\mu}$ together, and the corresponding spectrum of electron neutrinos at the source would be $Q_{\nu_e}(E) \approx \frac{1}{2} Q_{\nu_\mu}(E)$. Putting in the numbers given in Sec. II A, we find

$$Q_{\nu_\mu}(E) \approx 83.3 Q_n(25E) \quad \text{for } \alpha = 1, \tag{16}$$

$$Q_{\nu_\mu}(E) \approx 416. Q_n(25E) \quad \text{for } \alpha = 0. \tag{17}$$

We shall refer to Eqs. (14) and (15) as the *generic cosmic ray and neutrino production spectra*. We emphasize the strong dependence of the number of produced neutrinos per produced neutron on the assumed target photon spectral index: at ultrahigh energies, GRB produce about 5 times more neutrinos per neutron than AGN. We shall return to the implications of this result at the end of the paper.

III. PROPAGATION OF NEUTRINOS, PHOTONS, AND PROTONS OVER COSMOLOGICAL DISTANCES

In this section, we discuss propagation of cosmic rays and neutrinos in an expanding Universe filled with the cosmic microwave background radiation. To illustrate the problem, we compare the energy-loss horizons of protons and neutrinos. We shall then briefly discuss the physical problems connected to several approaches to cosmic ray propagation calculations. Using the numerical propagation code described by Protheroe and Johnson [19], we then calculate the observable neutrino and cosmic ray spectra from a cosmological distribution of generic photohadronic sources, as described in the last section. Here we assume that the sources are transparent to neutrons, while protons are confined, and that gamma rays are reprocessed in synchrotron-pair cascades until emitted in the energy range of 10 MeV–30 GeV. Using an extrapolated cosmic ray spectrum which is consistent with present observational limits on the light component of cosmic rays (i.e., protons) as an upper limit on the possible extragalactic proton contribution, and the diffuse extragalactic gamma ray background (EGRB) observed by the Energetic Gamma Ray Experiment Telescope (EGRET) as an upper limit on the hadronic extragalactic gamma ray flux, we determine an energy dependent upper bound on the neutrino flux from cosmic ray sources with the assumed properties. Our result is compared with the energy independent bound on extragalactic neutrino fluxes recently proposed by Waxman and Bahcall [1].

A. Comparison of energy-loss horizons

We wish to compare the distances that neutrinos, photons with energies below threshold for cascading in background radiation fields, and protons will travel through the Universe without significant energy losses. We define the energy loss horizon by

$$\lambda = cE/|dE/dt|, \tag{18}$$

such that for linear processes the energy is reduced to $1/e$ of its initial value on traversing a distance λ.

For gamma rays below $\sim$30 GeV and neutrinos, the energy-loss process is due to expansion of the Universe [20]. For simplicity, we adopt an Einstein–de Sitter cosmology (i.e., $\Lambda=0$ and $\Omega=1$), so that the horizon λ can be related to a redshift z by the redshift distance relation

$$\lambda(z) = \frac{2}{3}\frac{c}{H_o}[1-(1+z)^{-3/2}] \tag{19}$$

where $H_o = 50h_{50}\,\mathrm{km\,s^{-1}\,Mpc^{-1}}$ is the Hubble constant. Since we require the distance for which the energy is reduced by a factor e during propagation, we have $(1+z)=e$ giving the horizon for redshift losses:

$$\lambda_z = \frac{2c}{3H_o}(1-e^{-3/2}). \tag{20}$$

This is also the horizon for neutrinos and gamma rays below $\sim$30 GeV. Normalizing the neutrino horizon to the radius of the Einstein–de Sitter universe,

$$\hat{\lambda}_z \equiv \frac{3H_o}{2c}\lambda_z, \tag{21}$$

we obtain $\hat{\lambda}_\gamma = \hat{\lambda}_\nu = \hat{\lambda}_z \approx 0.78$.

In addition to redshift losses, extragalactic cosmic rays suffer energy losses from photohadronic interactions with cosmic backgrounds, mainly the microwave background, and this is the reason for the Greisen-Zatsepin-Kuzmin (GZK) cutoff expected for a cosmic ray spectrum originating from a cosmologically homogeneous source distribution [21,22]. Photopion production and Bethe-Heitler pair production govern the energy loss in different energy regimes due to their very different threshold energies. The Bethe-Heitler process limits the propagation of protons with energies $E_p > 2m_p m_e c^4/kT_{\mathrm{mbr}} \approx 4\times10^9$ GeV to $\lambda_p \sim 1$ Gpc, while pion production reduces the horizon for protons with $E_p \gtrsim m_p m_\pi c^4/kT_{\mathrm{mbr}} \approx 5\times10^{11}$ GeV to $\lambda_p \sim 10$ Mpc. Again, in units of the radius of the Einstein–de Sitter Universe, the energy-dependent horizon for protons can be written

$$\hat{\lambda}_p(E) = [1/\hat{\lambda}_z + 1/\hat{\lambda}_{p,\mathrm{BH}}(E) + 1/\hat{\lambda}_{p,\pi}(E)]^{-1}, \tag{22}$$

where the components expressing redshift, Bethe-Heitler, and pion production losses can be written as

$$\hat{\lambda}_z \simeq 0.78,$$
$$\hat{\lambda}_{p,\mathrm{BH}}(E) \approx 0.27h_{50}\exp(0.31/E_{10}), \tag{23}$$
$$\hat{\lambda}_{p,\pi}(E) \approx 5\times10^{-4}h_{50}\exp(26.7/E_{10}),$$

with $E_{10}=E_p/10^{10}$ GeV. The approximations for $\hat{\lambda}_{p,\mathrm{BH}}$ and $\hat{\lambda}_{p,\pi}$ fit the exact functions determined numerically in [19] and the exact interaction kinematics within $\sim$10% up to $E_p \sim 10^{12}$ GeV.

The different energy-loss horizons for gamma rays and neutrinos, and protons strongly affect the relative intensities of their diffuse isotropic background fluxes. This is true in particular for evolving source populations such as quasars, galaxies, or GRBs (if they trace star formation activity), since here most of the energy is released at large redshifts. Cosmic rays above the ankle ($E_{\mathrm{cr}} \sim 3\times10^9$ GeV) originate only from sources with redshifts $z \lesssim 0.27$, while neutrinos and gamma rays originate from sources within $z_\nu = z_\gamma \sim 1.7$. This will give rise to the neutrino intensity being enhanced relative to the protons because of their larger horizon, and because of the evolution of the sources (e.g., quasars) with cosmic time (redshift).

We may illustrate the problem as follows. The basic method of calculating the approximate present-day diffuse fluxes of neutrinos, gamma rays below $\sim$30 GeV, and cosmic rays of photoproduction origin, would be to integrate the contributions from sources at redshifts up to those corresponding to the respective energy-loss horizons. Assuming a constant source number per comoving volume element for simplicity, the resulting fluxes are proportional to $V_c(\hat{\lambda})/d_L^2(\hat{\lambda}) \sim \hat{\lambda}$, where $V_c(\hat{\lambda})$ and $d_L(\hat{\lambda})$ are the cosmological comoving volume and the luminosity distance, respectively, corresponding to the horizon $\hat{\lambda}$. Assuming photohadronic production of neutrinos and cosmic rays in, e.g., a cosmological (nonevolving) distribution of AGN, the relative flux of neutrons (assuming no absorption) with energy E_n, and corresponding neutrinos with energy $E_\nu = 0.04E_n$, is at the source given by $[E_\nu^2 N_\nu(E_\nu)]/[E_n^2 N_n(E_n)] = \xi_\nu/\xi_n \approx 0.2$. The same ratio will be observed in the integrated fluxes, as long $E_n < 4\times10^9$ GeV. For higher cosmic ray energies, the flux ratio must be multiplied by a factor $\hat{\lambda}_\nu/\hat{\lambda}_p$, which yields a flux ratio of $\sim$0.6 for $E_n \sim 10^{10}$ GeV, and $\sim$30 for $E_n \sim 10^{11}$ GeV. Obviously, the differences would be much larger if we had assumed strong source evolution which enhances the contribution from large distances. If we want to determine a neutrino spectrum from an observed, correlated cosmic ray spectrum from the same sources, the result must therefore approximately reflect the changes in the ratio of the energy-loss horizons.

B. Exact calculation of present-day neutrino and cosmic ray spectra

The example above, of simply integrating the cosmic ray and neutrino contribution of cosmologically distributed sources up to the energy-loss horizon, disregards several important aspects of particle propagation. First, the particle number is not conserved in this method. Second, the energy evolution of the particles is neglected, which removes the dependence on spectral properties of the source. Both caveats are removed in a method known as the continuous-loss approximation, which follows the particle energy along fixed trajectories as a function of cosmological distance (or redshift) [23,6]. For example, the trajectories for neutrino energies would be simply $E(z)=E_0(1+z)$. This method is exact for adiabatic losses due to the expansion of the Universe (i.e., particle redshift), and is still a very good approximation for Bethe-Heitler losses, but it gives only poor results for photopion losses. The reason for the latter is the large mean free path, and the large inelasticity of this process, which results in strong fluctuations of the particle energy around its mean trajectory [24,25]. Cosmic ray transport in the regime where photopion losses are relevant is therefore best described by numerical approaches, either solving the exact transport equation [24], or by Monte Carlo simulations [26]. Another important aspect concerning relative fluxes of cosmic rays, gamma rays, and neutrinos is the fact that the interaction of the cosmic rays with the cosmic background radiation themselves produces secondary particles, which have to be considered as an additional contribution to the primary neutrinos and gamma rays. The full problem is treated by the cascade propagation code which has been described in detail by Protheroe and Johnson [19], and which we shall use also here. The code is based on the matrix-doubling technique for cascade propagation developed by Protheroe and Stanev [27].

For an input spectrum $(dP_{\rm gal}/dV_c)\langle Q(E,z)\rangle$ per unit comoving volume per unit energy per unit time, the intensity at Earth at energy E is given by

$$I(E)\propto\frac{1}{4\pi}\int_{z_{\rm min}}^{z_{\rm max}}M(E,z)\,\frac{(1+z)^2}{4\pi d_L^2}\,\frac{dV_c}{dz}\,\frac{dP_{\rm gal}}{dV_c}\times\langle Q[(1+z)E,z]\rangle dz, \tag{24}$$

where d_L and V_c are luminosity distance and comoving volume, and $M(E,z)$ are ''modification factors'' for injection of protons at redshift z as defined by Rachen and Biermann [6]; for neutrinos, $M(E,z)=1$. The modification factors for protons depend on the input spectra, and are calculated numerically using the matrix method [19].

C. An abstract bound on extragalactic neutrino fluxes from neutron-transparent sources

From the above considerations, it is obvious that one can use the observed cosmic ray spectrum to construct a correlated neutrino spectrum, under the assumption that all observed extragalactic cosmic rays are due to neutrons ejected from the same sources as the neutrinos. A difficulty here is that the contribution of *extragalactic* cosmic rays to the total observed cosmic ray flux is unknown. Since neutrons convert into cosmic ray protons, we can clearly consider only the proton component at all energies. Above the ankle in the cosmic ray spectrum at $E_{\rm cr}\sim3\times10^9$ GeV, observations are generally consistent with a ''light'' chemical composition, i.e., a 100% proton composition is possible. Since protons at these energies cannot be confined in the magnetic field of the Galaxy, they are also likely to be of extragalactic origin. At extremely high energies, however, there is the problem that the event statistics is very low, and different experiments disagree on the mean cosmic ray flux at $\sim10^{11}$ GeV by one order of magnitude (see Bird *et al.* [28] for a comparison of the results until 1994 of the four major experiments Akeno, Havarah Park, Fly's Eye, and Yakutsk). This energy region is very important, since we expect here the existence of the GZK cutoff due to photoproduction losses in the microwave background. Currently, no clear evidence for the existence of this cutoff has been found, and the results of at least two major experiments (Havarah Park [29], and Akeno Giant Air Shower Array (AGASA) 1998 [30]) are consistent with the result obtained from a superposition of all experiments using a maximum likelihood technique [31,32], that is, a continuation of the cosmic ray spectrum as a power law $\propto E^{2.75}$ up to $\gtrsim3\times10^{11}$ GeV (see also Fig. 1).

The situation is even more difficult at lower energies: cosmic rays are here assumed to be mainly of Galactic origin, and there is evidence that a considerable, maybe dominant fraction consists of heavy nuclei rather than protons. Around the knee or the cosmic ray spectrum at $E_{\rm cr}\sim10^6-10^7$ GeV, recent results from the KASCADE air shower experiment suggest that the fraction of heavy nuclei in the cosmic ray flux is at least $\sim$30%, and further increasing with energy [33]. Also below the ankle, in the energy range 10^8-10^9 GeV, the analysis of air shower data has produced tentative evidence of a composition change from heavy to light (with increasing energy), supporting a dominantly heavy composition of cosmic rays between the knee and the ankle [34]. (Note that this result is under dispute, and it has been shown that it depends on the Monte Carlo simulation codes used to construct the air shower properties in dependence of the primary particle mass [35]. These simulation codes involve particle interaction models based on extrapolations many orders of magnitude above the energy range currently accessible with particle accelerators [36].)

Using all the available data, we find that an extragalactic cosmic ray spectrum of the form

$$N_{p,\rm obs}(E)=0.8\times(E/1\ \ {\rm GeV})^{-2.75}\ \ {\rm cm^{-2}\,s^{-1}\,sr^{-1}\,GeV^{-1}}\quad(3\times10^6\ \ {\rm GeV}<E<10^{12}\ \ {\rm GeV}) \tag{25}$$

is consistent with all data and limits on the cosmic ray proton flux (Fig. 1). It represents the current *experimental upper limit* on the extragalactic cosmic ray proton flux, which we shall use to construct an *upper limit* on the possible, diffuse extragalactic neutrino flux. If it can be shown that the inten-

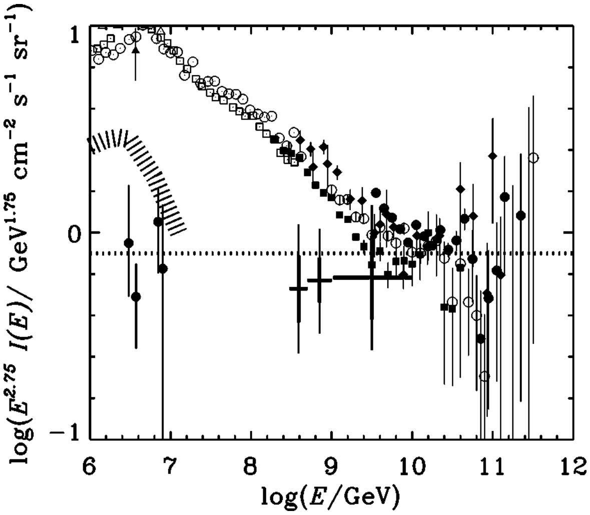

FIG. 1. The observed all-particle cosmic ray spectrum taken from the article by T. K. Gaisser and T. Stanev in the 1998 Review of Particle Properties [81], and supplemented by Fly's Eye monocular data [28] (open circles at high energy), and recent AGASA data [30] (filled circles at high energy). Also shown are estimates of the cosmic ray proton component: based on the proton fraction estimates by [34] (thick lines with thick error bars, extended by thin lines which indicate the systematic error due to normalization to the all-particle data); Norikura data [82] (filled circles with large error bars at $3\times10^6-10^7$ GeV); proton fraction estimated from KASCADE data [33] normalized to all-particle data (hatched band from 10^6-10^7 GeV). The spectrum we adopt for the proton component which forms the upper bound to any extragalactic cosmic ray proton spectrum [cf. Eq. (25)] is shown by the dotted line.

sity of protons at 10^9 GeV and lower energies is below that assumed in this paper (dotted line in Fig. 1), then the neutrino bound we have constructed below 10^7 GeV would need to be reduced.

To construct the neutrino bound, we assume test spectra of the form $Q_{\rm cr}(E)=Q_n(E)\propto E^{-1}\exp(-E/E_{\rm max})$ with $10^6\,{\rm GeV}<E_{\rm max}<10^{12}$ GeV. The corresponding neutrino spectra are determined using Eq. (16). We assume a source distribution following the cosmological evolution function found for galaxies and AGN ([37], see next section). For a given $E_{\rm max}$ the total contribution of cosmic rays and neutrinos is calculated using Eq. (24). The resulting spectrum is then normalized so that its maximum reaches the cosmic ray flux given by Eq. (25). By varying $E_{\rm max}$ between 10^6 and 10^{12} GeV, we then obtain the desired maximum flux of neutrinos consistent with the present cosmic ray data (see Fig. 2). We also consider the correlated gamma ray output, assumed to be twice the neutrino energy flux, and check it does not exceed the observed power-law component of the diffuse gamma ray background (we estimate the background between 3 MeV and 30 GeV to be $\approx1.5\times10^{-5}$ GeV cm^{-2} s^{-1} sr^{-1} [38]). Note that the cosmic ray curve for $E_{\rm max}=10^6$ GeV does not reach the estimated cosmic ray proton spectrum in Fig. 2(a) in order to avoid over-producing diffuse gamma rays. Note also that the propagated cosmic ray spectra cut off at or below 10^{11} GeV for all $E_{\rm max}$ values, and so our bound is insensitive to the assumed extragalactic cosmic ray spectrum above 10^{11} GeV.

The upper bound on the neutrino flux according to this model is given by the minimum of the cosmic ray and the gamma ray bound, shown in Fig. 3 by the curve "$\tau_{n\gamma}<1$." We see that the cosmic ray limit starts to dominate the bound for $E_\nu\gtrsim100$ TeV, which then decreases to a minimum at $E_\nu\approx10^9$ GeV, after which it rises again. This rise is a consequence of the strongly increasing ratio of neutrino and proton horizons above this energy, in conjunction with the assumptions that the distribution of sources is homogeneous in space. A more realistic scenario might be that the lack of evidence for the GZK cutoff in the cosmic ray data is due to the dominant contribution of a local cosmic ray source. In this case, our estimate that $[E_\nu^2N_\nu(E_\nu)]/[E_n^2N_n(E_n)]\propto\hat\lambda_\nu/\hat\lambda_p$ would not apply. Obviously, if an extended data set would confirm the existence of the GZK cutoff in the UHECR spectrum, the neutrino bound would remain at the level of the Waxman-Bahcall result (also shown in Fig. 3, see discussion below) for $E_\nu>10^9$ GeV.

By assumption, the neutrino bound constructed this way applies only to sources which are transparent to neutrons. For the opposite extreme, i.e., sources with a very high neutron opacity, $\tau_{n\gamma}\gg1$, it is still possible to set an upper limit using the observed EGRB, assuming that the dominant part of the emitted gamma radiation is in the EGRET range. This is shown by the line labeled "$\tau_{n\gamma}\gg1$" in Fig. 3. The range in between can be regarded as the "allowed range" for the neutrino emission from sources with $\tau_{n\gamma}>1$. In the next section, we will estimate specific neutrino spectra for AGN models which imply that $\tau_{n\gamma}(E_n)>1$ for high neutron energies.

Our result may be compared with the cosmic ray bound to extragalactic neutrino fluxes recently proposed by Waxman and Bahcall [1]. Their bound was constructed using a different approach: Waxman and Bahcall assume an E^{-2} input spectrum of extragalactic cosmic rays, and normalize the propagated spectrum to the observed flux at $E_{\rm cr}=10^{10}$ GeV. Consequently, their bound [determined for a source evolution $\propto(1+z)^{3.5}$] agrees with the one derived in this work at $E_\nu\sim5\times10^8$ GeV, where the cosmic ray limit is most restrictive.

We have chosen E^{-1} trial spectra with variable exponential cutoffs in order to be able to mimic the effect of the superposition of spectra from various source classes. The pronounced peak in the energy flux associated with the trial spectra allows one to normalize the neutrino flux consistent with the experimental upper limit on extragalactic protons ($\propto E^{-2.75}$) at any chosen energy. Other hard spectra or delta function distributions as trial spectra would have yielded practically the same result. Although canonical AGN jets, which we discuss in the next section, are an example for sources with E^{-1} (or similar) spectra, our result does not imply that we assume AGN jets to actually saturate the upper limit in general. As a matter of fact, at neutrino energies $\gtrsim10^9$ GeV, a class of photohadronic sources saturating our bound has neither been suggested nor does its existence seem likely on the basis of current knowledge. By restricting their source spectra to E^{-2}, Waxman and Bahcall have con-

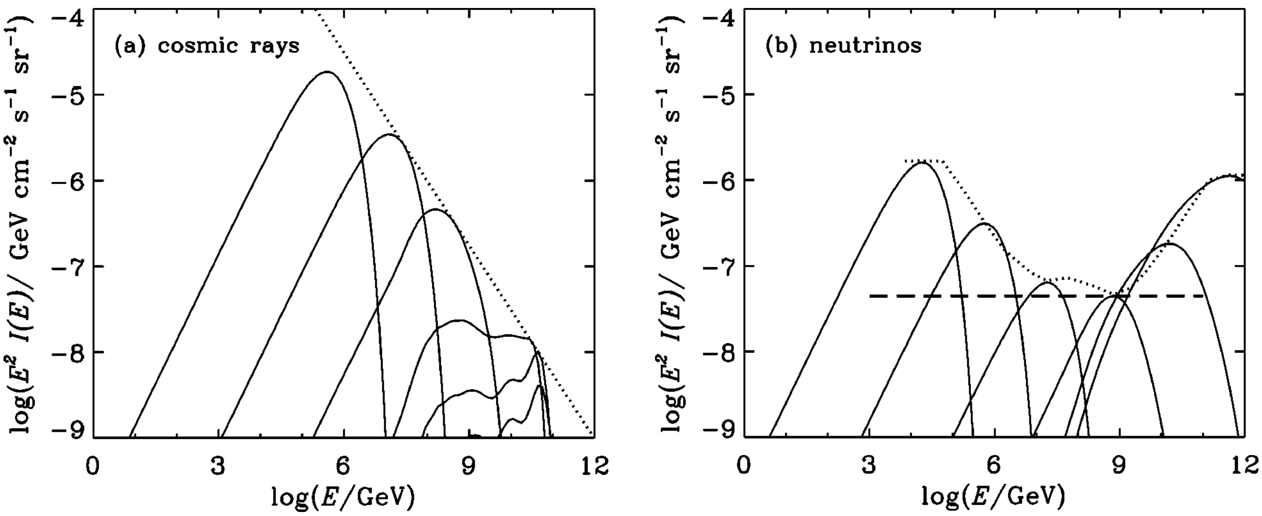

FIG. 2. Spectra of (a) cosmic rays and (b) neutrinos after propagation through the Universe of input spectra for optically thin pion photoproduction sources for $E_{max}=10^6$, 3×10^7, 10^9, 3×10^{10}, 10^{12}, and 3×10^{13} GeV, assuming galaxy evolution as described in the text. Spectra are normalized such that the cosmic ray intensity does not exceed the cosmic ray proton spectrum estimated from observations [dotted line in part (a)] and such that the neutrino energy flux does not exceed 0.5 of the observed photon energy density above 3 MeV. The dotted curve in part (b) joins the peaks in the neutrino spectra and forms our neutrino upper bound for optically thin pion photoproduction sources. The dashed line is the bound obtained by Waxman and Bahcall [1].

structed a bound for neutron-transparent sources which is probably closer to current models for cosmic ray and neutrino production than our general upper limit above 10^{19} eV. Examples for such models are the model for diffuse neutrino fluxes from GRBs proposed by the same authors [18,1] (note that this prediction assumes no cosmological evolution for GRBs), and model A in Mannheim [10] which was constructed using the cosmic ray limit with the assumption that the emerging neutrons at $\sim10^{10}$ GeV contribute to the extragalactic cosmic ray spectrum (both shown in Fig. 3).

IV. DIFFUSE NEUTRINO SPECTRA FROM AGN JETS

In this section, we shall consider the spectra of cosmic ray protons and neutrinos emerging from jets of two classes of gamma ray emitting AGN, i.e., BL Lac objects and radio quasars, which are usually combined as the class of *blazars*. We shall use the estimated optical depths of gamma rays to photon-photon pair production in a typical AGN of each type, to infer the corresponding neutron-photon optical depths in these objects. We shall show that high luminosity AGN (like 3C279) can be expected to be opaque to neutrons at energies above about 10^8-10^9 GeV, while low luminosity BL Lacs (like Mrk501) must be transparent to neutrons at all energies. Then assuming a model for the luminosity dependence of the optical depths, and the local luminosity functions of BL Lacs and quasars, we shall estimate the form of the production spectra of cosmic ray protons and neutrinos per unit volume of the local Universe. Applying the source evolution functions found for BL Lacs and radio quasars, respectively, we shall derive model estimates for the diffuse neutrino contribution from these sources which are compatible with cosmic ray limits. We shall also construct an upper bound for the contribution of AGN jets, using the same method as in the previous section, but for the appropriate generic source spectra, Eqs. (14) and (15), varying the break energy E_b over the range allowed by the models.

A. Cosmic ray proton and neutrino production spectra from blazars

As our starting point, we assume a target photon spectrum with index $\alpha=1$ which we have already seen leads to

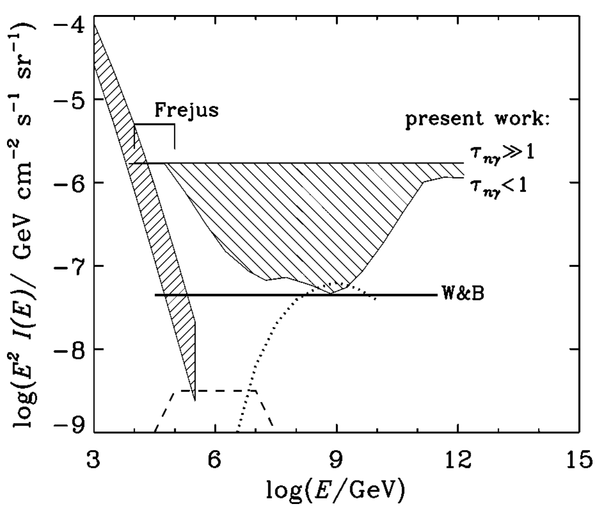

FIG. 3. Muon neutrino upper bounds for optically thin pion photoproduction sources (curve labeled $\tau_{n\gamma}<1$) and optically thick pion photoproduction sources (curve labeled $\tau_{n\gamma}\gg1$); the hatched range between the two curves can be considered the allowed region for upper bounds for sources with $\tau_{n\gamma}>1$. For comparison we show the bound obtained by Waxman and Bahcall [1] (for an evolving source distribution). Predictions for optically thin photoproduction sources are also shown: proton-blazar (Mannheim 1995 [10], model A)—dotted curve, and GRB sources (Waxman and Bahcall 1997, [18])—dashed curve. Also shown is an observational upper limit from Fréjus [83] and the atmospheric background [84].

$\tau_{p\gamma}(E_p) \propto E_p$. A similar energy dependence applies to gamma rays interacting with the same photons by photon-photon pair production ($\gamma\gamma \to e^+e^-$), and so for $\alpha=1$ we have

$$\tau_{p\gamma}(E_p) = \frac{\langle \kappa_p \sigma_{p\gamma} \rangle}{\langle \sigma_{\gamma\gamma} \rangle} \tau_{\gamma\gamma}\left(\frac{2m_e c^2 E_p}{m_\pi [m_p + m_\pi/2]} \right)$$
$$\approx 5 \times 10^{-4} \tau_{\gamma\gamma}([4 \times 10^{-6}]E_p). \tag{26}$$

Here we have used $\langle \kappa_p \sigma_{p\gamma} \rangle / \langle \sigma_{\gamma\gamma} \rangle \approx 300\ \mu\text{barn}/\sigma_T$ from averaging over a photon spectrum with $\alpha=1$, where $\langle \kappa_p \sigma_{p\gamma} \rangle$ includes Bethe-Heitler pair production, $\sigma_{\gamma\gamma}$ is the total cross section for the process $\gamma\gamma \to e^+e^-$ [39], and σ_T is the Thomson cross section. Using $\tau_{n\gamma}(E_n) \approx 2\tau_{p\gamma}(E_n)$, and assuming that $\tau_{n\gamma}(E) \propto E$ holds for a range $10^5 E_\gamma \lesssim E_n \leqslant E_{\max}$, we obtain the relation

$$\tau_{n\gamma}(E_n)/\tau_{\gamma\gamma}(E_\gamma) \approx 4 \times 10^{-9} E_n/E_\gamma. \tag{27}$$

We apply this relation to two reference AGN: the BL Lac object Mrk501, and the quasar 3C279. The combination of the observed TeV spectrum and the EGRET flux limits for Mrk501 [40,41] gives rise to our assumption of a break energy at about 0.3–1 TeV in the context of photohadronic models, while for 3C279 the EGRET spectrum allows for a break at $\sim$3–10 GeV [42]. Note that internal opacity due to the presence of low-energy synchrotron photons with spectral index α, which plays a crucial role in photohadronic models for the gamma ray emission, gives rise to a spectral steepening by $E^{-\alpha}$, and *not* to an exponential cutoff, above the energy where $\tau_{\gamma\gamma}=1$, see Eq. (11). A TeV power law being steeper than the gamma ray spectrum at EGRET energies could thus imply optical thickness in spite of reaching some 25 TeV as seen by the HEGRA telescopes [43]. Clearly, lower values for the optical depth are obtained if the data are interpreted with models which assume that the observed steepening is due to the radiation process itself (e.g., synchrotron–self-Compton emission) [40,41]. Applying Eq. (27) we find that the break energy in the cosmic ray production spectra should be $E_b \sim 10^{11}$ GeV in Mrk501, and $E_b \sim 10^9$ GeV in 3C279. Thus, UHE cosmic rays can escape from Mrk501 under optically thin conditions, and our upper bound is not affected by the actual value of the gamma ray optical depth in this source.

We note that the break energies derived above depend on the assumption of an undistorted power-law target spectrum of photons. This assumption might be reasonable for BL Lacs, which show power-law photon spectra extending from the infrared (relevant for $p\gamma$ interactions) into the x-ray regime (relevant for $\gamma\gamma$ absorption of $\sim$10–100 GeV photons, assuming a Doppler boosting of the emission by a factor $\delta \sim 10$). In high luminosity radio quasars, however, this is not the case (see, e.g., [44]), and relation (27) does not necessarily hold. Moreover, in these sources the dominant target spectrum for $p\gamma$ and $\gamma\gamma$ interactions could be given by the external accretion disk photons, forming roughly an (unboosted) power law spectrum with $\alpha=1$ up to $\sim$10 eV [45], where it drops by about one order of magnitude (e.g., [46]). This still implies a relation like Eq. (27), but with $\tau_{n\gamma}$ increased relative to $\tau_{\gamma\gamma}$ by a factor of 10—consequently, the possible break in the gamma ray spectrum of 3C279 would correspond to $E_b \sim 10^8$ GeV. A principal lower limit to the break energy in EGRET sources is set at $E_b \sim 10^7$ GeV, since otherwise EGRET photons ($\gtrsim$1 GeV) could not be emitted.

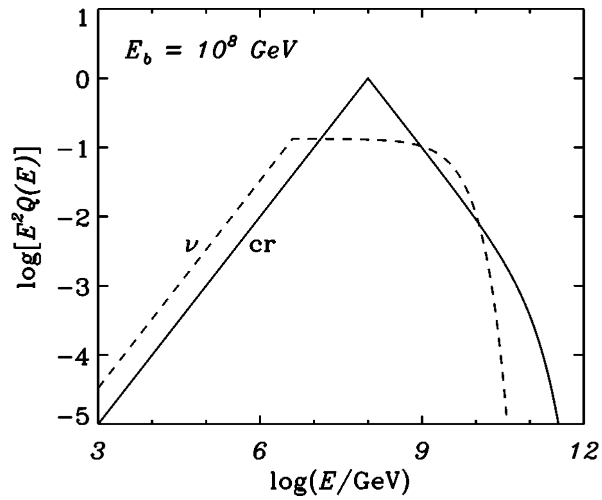

FIG. 4. From of the spectra of neutrinos and cosmic rays escaping from optically thick photoproduction sources, Eqs. (14) and (15). The example shown here has a break energy of 10^8 GeV, and an exponential cutoff above 10^{11} GeV.

Using cosmic ray proton production spectra with a break at $\sim 10^8$ GeV (see Fig. 4) will, of course, have a strong effect on the neutrino bound implied by cosmic ray data. Since the cosmic ray proton spectrum of the source above the break drops faster than the upper limit spectrum on cosmic ray protons, Eq. (25), the bound at a neutrino energy $E_\nu > E_b/25$ is essentially set by the cosmic ray flux at the break energy E_b rather than by the more restrictive flux at $25E_\nu$. If all sources had the same E_b, the bound would be increased roughly by a factor $(25E_\nu/E_b)^{0.75}$. Clearly, the assumption of optical depths allowing break energies of 10^8 GeV or below in luminous quasars is not directly supported, but only *consistent with* current observations. The maximal neutrino intensity for AGN jets, which we shall derive below on the basis of this assumption, is therefore at the currently allowed maximum for the adopted source model, and likely to be lowered when more detailed gamma ray data become available.

As a general limitation of this approach it must be noted that most AGN jets have not been detected by EGRET. If their lack of high-energy emission is due to a large gamma ray opacity, this would also imply very low break energies for their ejected cosmic ray spectrum. One such possible class of sources are the gigahertz-peaked sources (GPS) and compact steep-spectrum (CSS) quasars [47], which make up about 40% of the bright radio source population. Objects such as these could thus produce a diffuse neutrino intensity at the level of the EGRB or maybe even above, without violating any constraint. Also low-opacity sources for which the maximum proton energy remains much less than its theoretically allowed value and, in particular, below the value of the break energy, could lead to a higher neutrino intensity at low E_ν (limited only by the EGRB). We did not consider this possibility by keeping $E_{p,\max}$ fixed.

B. Blazar luminosity functions and generic models for the neutrino contribution from BL Lacs and radio quasars

In order to obtain a parameterization of blazar neutrino spectra, we need to express $\tau_{n\gamma}$ as a function of the blazar luminosity L, given at some frequency. Here we take into consideration that blazars are assumed to be beamed emitters, with a Doppler factor $\delta \sim 10$. The optical depth intrinsic to the emission region is proportional to L_t/R, where L_t is the intrinsic target photon luminosity and R is the size of the emitter. Since blazars are strongly variable objects [48], we can use the variability time scale $T_{\rm var}$ to estimate the intrinsic size by $R \sim T_{\rm var} c\,\delta$. Using also the relation between intrinsic and observed luminosity, $L = L_t \delta^4$, we obtain $\tau_{n\gamma} \propto L T_{\rm var}^{-1} \delta^{-5}$. Although there is no detailed study of a possible systematic dependence of $T_{\rm var}$ on L, the observations are compatible with no such correlation existing—for example, variability time scales of order 1 day are common in both, moderately bright BL Lacs like Mrk 501 or Mrk 421, and powerful quasars like 3C279, PKS0528+134, or PKS1622-297, whose optical luminosities differ by at least three orders of magnitude [48]. For the Doppler factors, there is no evidence for a systematic dependence on L either [17], although unification models for blazars and radio galaxies (see next section) suggest that BL Lacs are on average slightly less beamed ($\langle \delta \rangle \sim 7$) than radio quasars ($\langle \delta \rangle \sim 11$) [49]. Therefore, we may assume that for both BL Lacs and quasars $\tau_{n\gamma}(E,L) \propto L$ holds on average, and that $\tau_{n\gamma}(E,L,\text{BL Lac}) \sim 10\tau_{n\gamma}$ (E, L, quasar). It is interesting to note that this relation would imply a relation of the break energies of Mrk501 ($L_{\rm opt} \sim 10^{44}\,\mathrm{erg\,s^{-1}}$) and 3C279 ($L_{\rm opt} \sim 10^{47}\,\mathrm{erg/s}$) as $E_b(\text{3C279}) \sim 10^{-2} E_b(\text{Mrk501})$, consistent with our estimate obtained for intrinsic absorption from the gamma ray spectral break. Of course, this does not rule out the possibility that E_b might be systematically lower in 3C279 due to external photons, as argued above.

To determine the contribution of all blazars in the Universe, we have to relate the proton luminosity L_p to the blazar luminosity L in some frequency range, and then integrate over the luminosity function $d\mathcal{N}/dL$, determined for the same frequency range. Here we have to distinguish between the luminosities in the energy range where the target photons are, L_o, and the luminosity of the gamma rays L_γ, which are here assumed to be produced by hadronic interactions. Obviously, the latter implies $L_p \propto L_\gamma$, but on the other hand $\tau_{n\gamma} \propto L_o$. We also have to distinguish between the two classes of blazars: while for BL Lacs $L_\gamma \propto L_o$, observations rather suggest that for quasars $L_\gamma \propto L_o^2$ [44]. This leads to the following simple models:

For BL Lacs, we use the x-ray luminosity functions given by Wolter *et al.* [50] for x-ray selected BL Lacs,

$$d\mathcal{N}_{\rm BL}/dL_X \propto L_X^{-1.6} \quad (3\times 10^{43}\ \mathrm{erg\,s^{-1}} < L_X < 3\times 10^{46}\ \mathrm{erg\,s^{-1}}), \tag{28}$$

and use the relations $L_p \propto L_X$, and $E_b \propto L_X^{-1}$ with $E_b = 10^{11}\,\mathrm{GeV}$ for $L_X = 3\times 10^{44}\,\mathrm{erg\,s^{-1}}$ (Mrk501). Here we have assumed that $L_o \propto L_X$.

For quasars, we use the EGRET luminosity function given by Chiang and Mukherjee [51],

$$d\mathcal{N}_q/dL_\gamma \propto L_\gamma^{-2.2} \quad (10^{46}\ \mathrm{erg\,s} < L_\gamma < 10^{48}\ \mathrm{erg\,s^{-1}}), \tag{29}$$

and the relations $L_p \propto L_\gamma$, and $E_b \propto L_\gamma^{-1/2}$ where we consider two possible normalizations for $L_\gamma = 10^{48}\,\mathrm{erg\,s^{-1}}$ (3C279), which are $E_b = 10^9\,\mathrm{GeV}$ in case that jet-intrinsic photons dominate the target field, and $E_b = 10^8\,\mathrm{GeV}$ for the assumption that external photons are the dominant target. For illustration, we show in Fig. 4 cosmic ray and neutrino spectra on emission for $E_b = 10^8\,\mathrm{GeV}$ and $E_{\max} = 10^{11}\,\mathrm{GeV}$.

Then we obtain the form of the production spectra of cosmic ray protons and neutrinos averaged over the local universe due to these two classes of AGN,

$$\langle Q_{{\rm cr},\nu}(E) \rangle = \frac{\int Q_{{\rm cr},\nu}(E,L)(d\mathcal{N}/dL)dL}{\int L(d\mathcal{N}/dL)dL}, \tag{30}$$

where the input spectra $Q_{\rm cr}$, Q_ν are given by Eqs. (14) and (15), with $E_{\max} = 10^{11}\,\mathrm{GeV}$, and E_b, L_p given as functions of L as discussed above.

To integrate properly over redshift, we note that while quasars show strong evolution similar to galaxies (see below), BL Lacs show little or no evolution [52], and we shall take this into account when propagating these spectra through the Universe from large redshifts. For quasars, their luminosity per comoving volume has a pronounced peak at redshifts of $z \sim 2$, and declines or levels off at higher redshifts [53]. We shall assume that this effect is due to evolution of the number of quasars with z, rather than evolution of the luminosity of individual sources, which keeps the production spectra $\langle Q_{{\rm cr},\nu}(E) \rangle$ independent of z. A particular parameterization of the redshift-dependence of the (comoving frame) UV luminosity density of AGNs as inferred by Boyle and Terlevich [37], assuming an Einstein–de Sitter cosmology and $h_{50} = 1$, is given by

$$\frac{dP_{\rm gal}}{dV_c} = P_0 \begin{cases} [(1+z)/2.9]^{3.4} & (z < 1.9), \\ 1.0 & (1.9 \leq z < 3), \\ \exp[-(z-3)/1.099] & (z \geq 3), \end{cases} \tag{31}$$

where $P_0 = (3.0 \pm 0.3)\times 10^{44}\,\mathrm{erg\,s^{-1}\,Mpc^{-3}}$. Clearly, the normalization plays no role here since L_p is adapted to match the cosmic ray flux at earth. So for BL Lacs (no evolution), we simply use $dP_{\rm gal}/dV_c = 1$.

The result is shown in Fig. 5. As expected, the diffuse neutrino fluxes from AGN jet models exceed the bound for optically thin sources, but fall into the allowed region for sources optically thick for neutron emission. The BL Lac contribution falls below the bound, because it was derived for a nonevolving source distribution—we note that it still exceeds the corresponding Waxman-Bahcall bound for the case of no evolution. In addition to the models discussed above, we have also constructed and estimated an upper

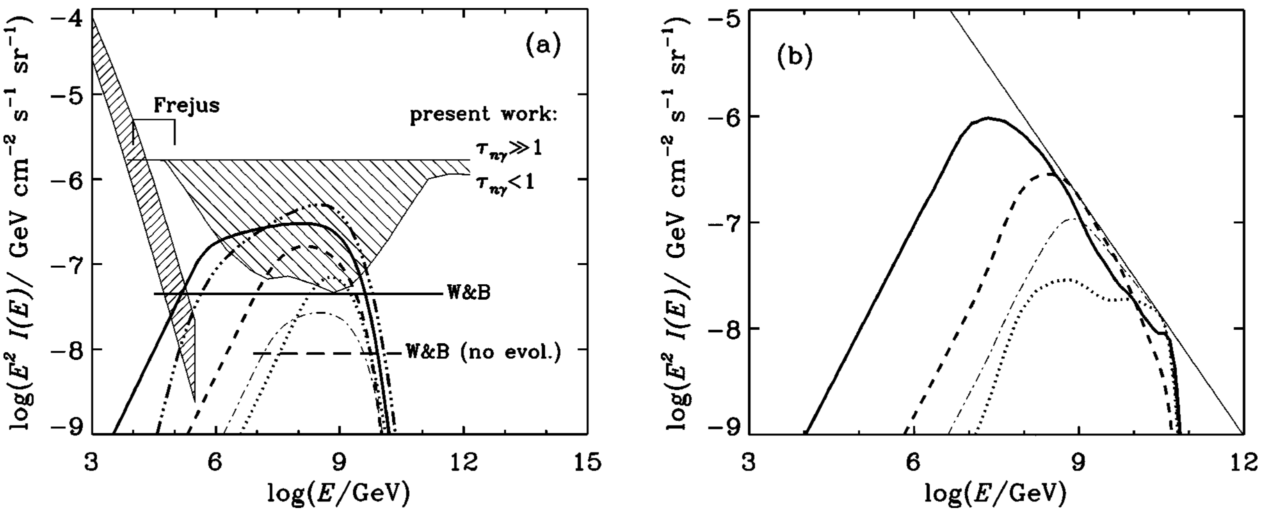

FIG. 5. (a) Comparison of neutrino spectra for optically thick pion photoproduction sources with neutrino upper bounds obtained in the present work and by Waxman and Bahcall [1]. Neutrino intensities obtained in the present work plotted are maximal superposition of spectra having E_b in the range 10^7-10^{11} GeV (thick solid curve); maximum source spectra averaged over the quasar luminosity function assuming $\tau_{n\gamma}\propto L^{1/2}$ with break energies corresponding to $L=10^{48}\,\mathrm{erg\,s^{-1}}$ at $E_b\approx 10^8$ GeV (thick dashed curve) and at $E_b\approx 10^9$ GeV (thick dotted curve); maximum source spectra averaged over the BL Lac luminosity function assuming $\tau_{n\gamma}\propto L$ with the break energy corresponding to $L=3\times 10^{44}\,\mathrm{erg\,s^{-1}}$ at $E_b\approx 10^{11}$ GeV (thin dot-dashed curve). Also shown is the prediction by Protheroe [45] for an external photon optically thick proton blazar model normalized down in accordance with the recent estimates of the blazar contribution to the diffuse gamma ray background [53] (thick chain curve), and the bounds obtained by Waxman and Bahcall [1] with and without source evolution (note that the maximum BL Lac intensity was calculated assuming no evolution). Other symbols shown correspond to Fig. 3. (b) Cosmic ray intensities for the neutrino intensities obtained in the present work shown in part (a) [key to these curves is as in part (a)]. The upper bound to any extragalactic cosmic ray proton spectrum is shown by the thin line, which corresponds to the dotted line in Fig. 1.

bound on the diffuse neutrino contribution on AGN jets. Here we used the same method as for the construction of Fig. 3, but implying input spectra of the form of Eqs. (14) and (15), with variable E_b and fixed $E_{\max}=10^{11}$ GeV. The break energy E_b was then varied in the range $10^7\,\mathrm{GeV}<E_b<E_{\max}=10^{11}$ GeV, and the normalization chosen such that the superposed spectra approximately represented the upper limit on the extragalactic cosmic ray spectrum (Fig. 5, left panel). We note that this upper bound corresponds within a factor of 2 with the prediction of a previously published model by Protheroe [45]. Other published models, for example by Halzen and Zas [11], or Mannheim ([10], model B), exceed our bound by about one order of magnitude at energies $E_\nu\gtrsim 10^8$ GeV, but their predictions for the important energy range below $\sim 10^7$ GeV are compatible with this bound. We return to the discussion of these models later, after we discuss the effect of magnetic fields in the AGN environment in the next section.

V. MAGNETIC FIELDS AND THEIR IMPACT ON COSMIC RAY PROPAGATION

As discussed in Sec. II, protons accelerated inside AGN jets are likely to be trapped in the jet to be released later near the end of the jet, and will consequently suffer severe adiabatic losses as a result of jet expansion. The bound that we calculated for optically thin photoproduction sources shown in Fig. 3 may of course be exceeded for optically thick sources. However, it may even be exceeded for optically thin photoproduction sources if the cosmic rays resulting from photoproduced neutrons are trapped, or suffer severe adiabatic deceleration in the large-scale magnetized environment such as the host galaxy and its halo, a galaxy cluster, or a supercluster. We emphasize that although there is now useful information available about magnetic fields in galaxies and clusters, the magnetic field structure and topology is not sufficiently well known for us to predict reliably magnetic trapping and adiabatic losses in host galaxies and clusters. However, we shall discuss in this section the fate of the cosmic rays resulting from photoproduced neutrons, and show that in some plausible scenarios, our bound for optically thin photoproduction sources can be exceeded for neutrinos from optically thin photoproduction sources.

X-ray observations and measurements of extragalactic Faraday rotation suggest that structures surrounding compact AGN jets carry magnetic fields of the order of 0.1–10 μG [54]. These can influence the propagation of cosmic ray protons in essentially two ways: (a) particles may be physically confined in the structure for a time $t\geqslant t_H=1/H_o$, or (b) the diffusive escape of the particles can lead to adiabatic energy losses. Obviously, magnetic fields on scales larger than the mean free path of a neutron for β decay, $l_n\approx 10\,\mathrm{kpc}\,(E_n/10^9\,\mathrm{GeV})$, also affect cosmic rays which are ejected as neutrons from the source. Here we shall discuss what influence these effects can have on the strength of the measured cosmic ray flux relative to the corresponding neutrino flux. We shall estimate the critical energy E^* below which our bound could plausibly be exceeded for optically thin photoproduction sources for a number of scenarios, in particular for clusters of galaxies and radio galaxies hosting AGN.

A. Particle confinement in clusters and superclusters

Clusters of galaxies have been recently discussed in the literature as a possible ''storage room'' for cosmic rays [55–57], because of their relatively strong magnetic fields ($B_0 \gtrsim 1\ \mu$G) extending over large scales, i.e., cluster radii, $R_{\rm cl} \gtrsim 1$ Mpc. The time scale for diffusive escape of cosmic rays in a turbulent magnetic field of homogeneous strength B_0 within a central radius R_0 is given by

$$t_{\rm esc} \sim \frac{R_0^2}{D} = \frac{3eB_0R_0^2}{cE_{\rm cr}\lambda(E_{\rm cr})}, \tag{32}$$

where we have used for the diffusion coefficient $D = \frac{1}{3}\lambda r_L c$, and $\lambda(E_{\rm cr})$ is the scattering length in units of the Larmor radius, $r_L = E_{\rm cr}/eB_0$ of a cosmic ray proton of energy $E_{\rm cr}$. The function λ depends on the turbulence spectrum of the magnetic field, $[\delta B(k)]^2 \propto k^{-y}$, where k is the wave number. Cosmic ray scattering is dominated by field fluctuations on the scale of the Larmor radius, i.e., $k \sim r_L^{-1}$. This implies that

$$\lambda(E_{\rm cr}) = \left[\frac{eB_0a_0}{E_{\rm cr}}\right]^y \quad \text{for } E_{\rm cr} < eBa_0, \tag{33}$$

where $a_0 = k_{\rm min}^{-1}$ is the largest scale of the turbulence, i.e., the ''cell size'' (or ''reversal scale'') of the magnetic field. For $k > a_0^{-1}$, the magnetic turbulence in clusters of galaxies seems to be well described by the Kolmogorov law for fully developed hydrodynamical turbulence, i.e., $y = \frac{2}{3}$. This can be easily seen from relating the typical turbulent magnetic field and cell size found from Faraday rotation measurements, $a_0 \sim 20$ kpc and $B_0 \sim 1\ \mu$G [58,59], to the diffusion coefficient found for electrons of energy ~ 1 GeV from the synchrotron radio emission spectrum [60], $D \sim 2\times 10^{29}\ {\rm cm^2\,s^{-1}}$. For $E_{\rm cr} > eBa_0$, i.e., $r_L > a_0$, the particle motion is a random walk with scattering angles $\sim a_0/r_L$ in each step. This can also be approximately described by Eqs. (32) and (33) by setting $y = -1$.

Confinement of cosmic rays over the cluster radius, $R_0 = R_{\rm cl} \sim 3$ Mpc, is then obtained for $t_{\rm esc} > t_H$, corresponding to a critical energy

$$E^*_{\rm diff} \approx 5\times 10^8\ {\rm GeV} \times \left[\frac{B_0}{1\ \mu{\rm G}}\right]\left[\frac{R_{\rm cl}}{3\ {\rm Mpc}}\right]^6\left[\frac{a_0}{20\ {\rm kpc}}\right]^{-2}, \tag{34}$$

provided $a_0 > 1$ kpc $(R_{\rm cl}/3\ {\rm Mpc})^2$. This assumes that the magnetic field strength is homogeneous over the entire cluster. However, if it dropped to $\sim 0.1\ \mu$G, and $a_0 \sim 200$ kpc at the edge of the cluster then $E^*_{\rm diff}$ would be a factor $\sim 10^3$ lower. On the other hand, some observations suggest also larger magnetic fields on smaller reversal scales [61,62], which would imply higher confinement energies.

The scenario above assumes that the background plasma filling the cluster is at rest. However, simulations of structure formation suggest that clusters of galaxies are accreting extragalactic hot gas [63], forming inflows of typical speeds $v_{\rm acc} \gtrsim 300\ {\rm km\,s^{-1}}$ downstream of an accretion shock near the outer radius of the cluster [7]. Since particles diffuse relative to the plasma flow, we require that in order to escape from the cluster, the average ''speed'' of particle diffusion, $\sim D/R_{\rm cl}$, exceeds the inflow velocity. If we assume the lower value for the magnetic field at $R_{\rm cl} \sim 3$ Mpc, i.e., $B_0 \sim 0.1\ \mu$G and $a_0 \sim 200$ kpc we obtain a critical energy

$$E^*_{\rm acc} \approx 2\times 10^6\ {\rm GeV} \times \left[\frac{B(R_{\rm cl})}{0.1\ \mu{\rm G}}\right] \times \left[\frac{R_{\rm cl}}{3\ {\rm Mpc}}\,\frac{v_{\rm acc}}{300\ {\rm km\,s^{-1}}}\right]^3\left[\frac{a(R_{\rm cl})}{200\ {\rm kpc}}\right]^{-2}. \tag{35}$$

Of course, with higher magnetic fields and a shorter field-reversal length a_0, $E^*_{\rm acc}$ would be higher. Similar processes would occur in the cooling flows observed in some rich clusters harboring powerful radio galaxies, but the confinement energies for typical cooling flow parameters [64,65] are even lower. Also the effect of drift in the nonhomogeneous cluster field [1] is small compared to diffusion, and cannot lead to larger confinement energies.

Ensslin *et al.* [57] point out that simple diffusion, as assumed above, may not be the best description of cosmic ray propagation in clusters of galaxies, and suggest that a one-dimensional random walk along static, but randomly tangled, magnetic field lines may be more realistic for particles with $r_L \ll a_0$. In this case the field line topology may be considered as arising from a three-dimensional random walk with steps of size a_0, yielding an effective diffusion length of $R_0 = R_{\rm cl}^2/a_0$. For the escape time we then obtain

$$t_{\rm esc} = \frac{R_{\rm cl}^4}{2ca_0^3}\left[\frac{eB_0a_0}{E_{\rm cr}}\right]^{1/3} \gtrsim t_H \quad \text{for } E_{\rm cr} \ll eB_0a_0, \tag{36}$$

which means for typical cluster parameters that essentially all cosmic rays with $E_{\rm cr} \lesssim 3\times 10^9$ GeV are confined to the cluster. Of course, this result neglects cross-field diffusion so that the maximum confinement energy is probably lower, but it could still be higher than the result obtained in Eq. (34) for the case of simple diffusion, in particular in the likely case that the field strength decreases with the distance from the cluster core.

Confinement of cosmic rays in clusters would lead to a decrease of the cosmic ray flux measured at earth relative to the corresponding neutrino flux, causing an increase of the neutrino bound. It is important to note, however, that cosmic ray confinement may exist even on larger scales, i.e., superclusters. It has been shown that magnetic fields $\sim 0.1\ \mu$G in superclusters are consistent with observations, and expected in simulations of structure formation which also predict accretion of gas with speeds $v_{\rm acc} \sim 1000\ {\rm km\,s^{-1}}$ [66,67]. The cosmic ray confinement energies for these larger scales ($R \sim 10$ Mpc, $a \sim 1$ Mpc) could therefore be $\gtrsim 10^8$ GeV, where again topological aspects connected to the detailed structure of the field may allow even higher values. Since our Galaxy itself is located in a supercluster, this latter scenario would tend to *increase* the cosmic ray flux relative to the corresponding neutrino flux, which fills the Universe homogeneously—thus it could actually decrease the bound below $E_\nu \sim 10^7$ GeV. The net effect of confinement in clusters and our local supercluster is obviously strongly model dependent, and therefore difficult to estimate.

An effect of the unknown structure of magnetic fields on galaxy cluster and supercluster scales may therefore most likely affect the observed extragalactic cosmic ray flux at energies below $\sim 2\times 10^6$ GeV, corresponding to $E_\nu \sim 10^5$ GeV. On the other hand, we can exclude an effect only for energies $E_{\rm cr}\gtrsim 3\times 10^9$ GeV ($E_\nu \gtrsim 10^8$ GeV), but we note that cosmic ray confinement at energies higher than $\sim 10^8$ GeV requires extreme assumptions on the strength and topology of the magnetic fields. We therefore agree with the conclusion of Waxman and Bahcall [1] that at $E_\nu \sim 5\times 10^8$ GeV an effect of large scale magnetic fields on the relation of cosmic ray and neutrino fluxes cannot be expected.

B. Adiabatic losses in expanding radio lobes and halos

Although many galaxy clusters have powerful radio galaxies at their centers, it is also a fact that most powerful radio galaxies are not found in such environments [68,69]. For a radio galaxy located in the normal extragalactic medium, evidence has been found that a pressure equilibrium with the external medium cannot be obtained within the lifetime of the source ($\lesssim 10^8$ yr) [70]. Therefore, it can be concluded that the lobes must expand. For a sample of powerful double-lobe (or FR-II) radio galaxies, lobe propagation velocities of order $10^4\,{\rm km\,s^{-1}}$ have been inferred [71]. Since the aspect ratio of the sources is found to be independent of their size, consistent with a propagation of the lobes along a constant opening angle of $\sim 10°$, expansion velocities of the lobes and the connecting ''bridges'' of $v_{\rm ex}\sim 1000\,{\rm km\,s^{-1}}$ can be inferred.

Neutrons produced by pion photoproduction interactions at an acceleration site near the base of the jet will be beamed preferentially along the jet direction decaying farther out along the jet or in the radio lobe, where they decay. The resulting cosmic ray protons will then be advected with the outflowing plasma on a time scale of the galaxy lifetime $\sim 10^8\,{\rm yr}\ll t_H$. Additionally, the protons perform a random walk in the magnetic field, which may even decrease their confinement time. Expanding radio lobes can therefore not confine cosmic rays indefinitely. However, in plasma outflows all particles which are isotropized in the flow due to scattering with plasma turbulence will experience adiabatic losses before their release. Kinetic theory implies that the particle energy at ejection from the flow is related to the injected particle energy by

$$\frac{E_{\rm ej}}{E_{\rm inj}}=\frac{R_{\rm inj}}{R_{\rm max}}. \tag{37}$$

Here, $R_{\rm max}$ is the radius of the outer termination shock of the flow, and $R_{\rm inj}$ is in general given by some minimum radius $R_{\rm min}$ where the outflow starts. For the case of neutron ejection from a central source, as considered here, $R_{\rm inj}=l_n$ if $l_n > R_{\rm min}$ and $R_{\rm inj}=R_{\rm min}$ otherwise. Note that the lobes we are considering here are much more extended than the jets in which the particles are accelerated.

The observed synchrotron spectra from radio lobes imply typical magnetic fields in the range 10–50 μG, turbulent with a maximum scale (or cell-size) of ~ 0.5 kpc [72,71]. The observed asymmetric depolarization in double radio galaxies leads to the suggestion that the magnetized plasma around the radio galaxy extends into a halo of radius ~ 300 kpc, with $B\sim 0.3\,\mu$G and a cell size of ~ 5 kpc at a radius $R\sim 100$ kpc [73]. The properties of the magnetic field in the central lobe and the halo can be connected by assuming that magnetic field and cell size scale as $B=B_0(R/R_0)^{-2}$ and $a=a_0(R/R_0)$, respectively, with $B_0\sim 30\,\mu$G, $a_0\sim 0.5$ kpc, and $R_0=10$ kpc.

Within radius R_0, we assume the properties of the turbulent magnetic field are constant. This corresponds to the assumption of an isotropic magnetic field expanding in a plasma outflow with $R_{\rm min}=R_0$ and $R_{\rm max}\sim 300$ kpc, which will be used as a working hypothesis in the following. We can consider cosmic ray protons as nearly isotropized in the plasma if $r_L(R)<a(R)$, corresponding to energies $E_{\rm cr}(R) < E^*_{\rm ad}(R)$ with

$$E^*_{\rm ad}(R)=10^{10}\,{\rm GeV}\times\left[\frac{B_0}{30\,\mu{\rm G}}\right]\left[\frac{a_0}{0.5\,{\rm kpc}}\right]\left[\frac{R}{10\,{\rm kpc}}\right]^{-1} \tag{38}$$

for $10\,{\rm kpc}<R\lesssim 300\,{\rm kpc}$. To justify the assumption of near isotropy below $E^*_{\rm ad}(R)$, we have to show that advection in the flow indeed dominates over diffusion, i.e., that the centers of the diffusive cloud of cosmic ray protons move approximately steady with the flow. Indeed, for the parameters adopted above we find for the ''speed'' of diffusive escape $D/R\sim \frac{1}{3}\,a(R)v_{\rm cell}/R<v_{\rm ex}$, where $v_{\rm cell}\leqslant \frac{1}{3}\,c$ is the average velocity of the particles to cross the cell, considering that also this motion is diffusive if the magnetic turbulence proceeds to smaller scales as expressed in Eq. (33).

Consequently, we can consider the cosmic rays with $E_{\rm cr}\lesssim E^*_{\rm ad}(R)$ as are advected with the flow at a distance R from the center of the galaxy, so that they suffer adiabatic losses following Eq. (37) leading to $E_{\rm cr}(R)\propto R^{-1}$. This implies that $E_{\rm cr}/E^*_{\rm ad}$ is independent of R, i.e., that cosmic rays confined in the outflow at some radius R remain confined for larger radii. If we consider cosmic rays which are ejected as neutrons, the radius where they couple to the magnetic field is given by the β-decay mean free path, $l_n\approx 10\,{\rm kpc}\,(E_n/10^9\,{\rm GeV})$. The resulting protons are confined to the outflow and are subject to adiabatic losses provided $E_n<E^*_{\rm ad}(l_n)$, or

$$E_n<3\times 10^9\,{\rm GeV}\left[\frac{B_0}{30\,\mu{\rm G}}\right]^{1/2}\left[\frac{a_0}{0.5\,{\rm kpc}}\right]^{1/2}. \tag{39}$$

The energy on ejection from the outflow is then $E_{\rm ej}\approx E_n/30$ for $E_n<10^9$ GeV, and $E_{\rm ej}=E_n^2/(3\times 10^{10}\,{\rm GeV})$ for $10^9\,{\rm GeV}<E_n<3\times 10^9$ GeV. For $E_n>3\times 10^9$ GeV, cosmic rays (neutrons and protons) traverse the lobe-halo along almost straight paths and adiabatic losses do not apply. (Note that the corresponding energies where these modifications may affect the neutrino bound are $E_\nu\sim E_n/20$, thus $E_\nu<10^8$ GeV.)

Energy losses of cosmic rays in the lobes and halos of radio galaxies are of particular relevance for models of neutrino production in AGN jets, which we discussed in Sec. IV. These models apply to radio loud AGN, which are likely to be the beamed counterparts of radio galaxies [49]. The two classes of AGN discussed in the last section correspond to the two Fanaroff-Riley (FR [74]) classes of radio galaxies: radio quasars might be associated to the powerful double-lobed FR-II radio galaxies, while BL Lac objects might correspond to the less luminous FR-I radio galaxies which generally have diffuse lobes centered around the AGN. The parameters used above for the lobes-halos of radio galaxies were mainly obtained from observations of FR-II radio galaxies or radio quasars. Therefore, the cosmic ray ejection from radio quasar (FR-II) sources can be expected to be diminished by more than an order of magnitude below $\sim 10^9$ GeV. For the less luminous FR-I galaxies, it could be that the magnetic fields, turbulence scales, and halo sizes used above are overestimated, so that $E^*_{\rm ad}$ could be lower by about one or two orders of magnitude for these sources.

VI. DISCUSSION AND CONCLUSIONS

We have looked at the problem of to what extent a possible flux of extragalactic neutrinos in the broad energy range $10^5\ \mathrm{GeV} \lesssim E_\nu \lesssim 10^{11}$ GeV is bounded by the observed flux of cosmic rays and gamma rays. As the minimum contribution to the cosmic rays we consider the neutrons produced together with other neutrals such as gamma rays and neutrinos in photohadronic interactions. The most restrictive bound arises for sources which are transparent to the emission of neutrons and where the protons from the decaying neutrons are unaffected by large-scale magnetic fields. The bound for this case is approximately in agreement with the bound previously computed by Waxman and Bahcall [1] in the neutrino energy range $10^7\ \mathrm{GeV} \lesssim E_\nu \lesssim 10^9$ GeV but is higher at lower and at higher energies. The difference is a consequence of the different approaches: Waxman and Bahcall assume a *particular model spectrum*, which could explain the observed UHECR spectrum, but which has a pronounced GZK cutoff that is not clearly confirmed by the present low statistics data, and which does not reproduce the steep slope of the observed CR spectrum below $\sim 10^{10}$ GeV. It thus leaves room for additional contributions from extragalactic sources (with different spectra) outside the narrow energy range where it matches with the observed cosmic ray flux. Our approach considers this possibility by using the best currently available *experimental* upper limit on the extragalactic proton contribution. We have also pointed out in Sec. II C, referring to recent Monte Carlo simulations [15], that the fundamental properties of photohadronic interactions can affect the bound by up to a factor of 5.

At neutrino energies below 10^7 GeV, the cosmic ray bound computed in this work rises and equals the bound inferred from the EGRB flux at about 10^5 GeV, below which the EGRB constraint is tighter. In the same energy region, we have also shown that effects from extragalactic magnetic fields come into play; they could either increase or reduce the bound, depending on the details of the field strength and structure, and the distribution of cosmic ray–neutrino sources. We therefore do not regard the bound based on the cosmic ray flux as a firm upper limit on neutrino fluxes below about 10^5 GeV, the main energy region for underwater-ice Cherenkov experiments such as the Antarctic Muon and Neutrino Detector Array (AMANDA) [75]. Models which predict neutrino emission mainly in this energy range are therefore not rigorously bounded by cosmic ray data, as for example the model of Berezinsky *et al.* [56] predicting neutrino fluxes from cosmic rays stored in clusters of galaxies. Such sources could in principle produce neutrino fluxes almost up to the EGRB limit within the AMANDA range. However, the neutrino contribution from clusters of galaxies is expected to be much lower, as it is limited by their expected contribution to the EGRB of $\sim 1\%$ [76].

At neutrino energies above 10^9 GeV, the upper limit rises up to the point where the energy flux of secondary gamma rays increases above the level of the observed EGRB. The reason for the increasing bound is that while cosmic rays from evolving extragalactic sources above the nominal GZK cutoff reach us exponentially damped due to interactions with the microwave background, neutrinos reach us essentially unattenuated. An observed excess of cosmic ray events above the GZK cutoff would therefore correspond to an even more pronounced excess of neutrinos in the framework of such models. Cosmic rays from a fiducial class of extragalactic photohadronic sources with a sufficiently flat spectrum could thus saturate the rising bound at the highest energies and would still remain below the local cosmic ray flux at all energies. Clearly, this is not the most likely scenario to explain the events beyond the GZK cutoff. For example, if these events are due to a single strong nearby source, then we would not expect the extragalactic neutrino flux above 10^9 GeV to be at a level higher than given by the bound computed by Waxman and Bahcall, even if this source would have the spectral properties which we used to construct our bound in Sec. III C. On the other hand, we note that the flux from non-photohadronic sources, as, for example, from the decay of topological defects, can exceed this bound because of their different branching ratios for the production of baryons relative to mesons [12]. Measuring the neutrino flux in this energy region, as is planned using large air-shower experiments such as the Pierre Auger Observatory [77], would therefore be highly relevant for understanding the nature and cosmic distribution of UHECR sources.

The bound we derive for extragalactic neutrinos is indeed an upper limit, since we do not consider the possibility that ultrarelativistic protons are ejected from their sources without preceding isospin flip interactions (''prompt protons''). Obviously, any additional prompt proton emission would have the effect to lower the relative neutrino flux consistent with the upper bound. However, for the most interesting sources, i.e., GRBs and AGN jets, the ejection of prompt protons can be expected to be strongly suppressed due to the rapid expansion of the emission region and the implied large adiabatic losses. This will allow us in the future to use measured neutrino fluxes, or experimental limits thereof, as a limit to the contribution of these sources to the cosmic ray flux. If, for example, GRBs produce the observed UHECR flux and if the global GRB emissivity follows a similar evolution as star formation and AGNs, bursts would have to emit a neutrino flux on the level of their optically thin bound, which is roughly two orders of magnitude above the original prediction [18]. This is independent of the total fractional burst energy converted by the protons in their interactions. A factor of ~3–6 is due to evolution, a factor of ~10 due to the limited efficiency of $p \rightarrow n$ conversion [78], and an additional factor ~5 due to an accurate treatment of photoproduction yields in the hard photon fields typical for GRBs (note that the yield factors used in our figures correspond to AGN-like target photon spectra). In this case, the gamma rays would also account for most of the EGRB, and one could observe strong TeV bursts from nearby sources [79,80]. Currently, only AGN jets are known to contribute to the EGRB in a major way and this has motivated the original model A of Mannheim [10] (which is exactly on the level of the optically thin bound as a consequence of the assumption that AGN jets produce the observed UHECRs). We note that the neutrino flux due to nucleons escaping the AGN jets and diffusing through the surrounding AGN host galaxies is difficult to assess and limited by the EGRB only [10], although this neutrino flux is the one most relevant in the 1–100 TeV range.

Generally, the relevance of a bound on optically thin sources should not be overestimated. A large number of extragalactic neutrino sources could be opaque to the emission of UHE cosmic rays producing a neutrino flux well above this bound. For the particularly important class of AGN jets, we have therefore developed an upper bound to their total neutrino flux using constraints on their gamma ray and neutron opacity set by present gamma ray data. This *maximum* contribution of AGN jets to the extragalactic diffuse neutrino flux is up to an order of magnitude higher than our optically thin bound in the energy range between 10^7 GeV and 10^9 GeV. Additionally, we have discussed that for energies below $\sim 10^8$ GeV an influence of magnetic fields in the radio lobes of active galaxies or on larger scales cannot be excluded. For example, the model of Halzen and Zas [11] predicts a neutrino flux about one order of magnitude above our upper limit for AGN at $E_\nu \sim 10^8$ GeV, which is certainly extreme, but still cannot be considered as completely ruled out, because at this energy the flux of neutrons after turning into cosmic ray protons may be diminished by an additional factor ~10 due to adiabatic losses in expanding radio halos. At present, the most general bound to EGRET detected AGN is given by their contribution to the EGRB. The most extreme class of neutrino sources opaque to the emission of UHECRs could be AGN jets which have *not* been detected in gamma rays. Prime candidates are the GPS and CSS quasars mentioned in Sec. IV A, which together make up about 40% of the bright extragalactic radio sources.

Finally, we wish to clarify that our result is not in conflict with, but complementary to, the upper limit previously obtained by Waxman and Bahcall [1]. Their result applies to photohadronic sources with a particular spectral shape which are transparent to the emission of UHE cosmic rays, such as may be representative of low-luminosity BL Lacertae objects (like Mrk501) or GRBs. Our result is more general, and therefore less restrictive, indicating that there may be other classes of sources, such as quasars, with a different spectral shape, and/or which are opaque to the emission of UHE cosmic rays, which can produce a higher neutrino flux than the source classes considered by Waxman and Bahcall. We also confirm the claim by Waxman and Bahcall that large scale magnetic fields are unlikely to have an effect for cosmic ray propagation at $\sim 10^{10}$ GeV, but we additionally point out that magnetic field effects cannot be disregarded at lower cosmic ray energies. We show that hadronic processes in AGN which are optically thick to the emission of UHE cosmic rays *could* produce the extragalactic gamma ray background according to present observational constraints. However, the limiting model for EGRET-detected AGN presented in this paper is mainly determined by the current instrumental limits of gamma ray astronomy—future observations in the GeV-to-TeV energy range may impose stricter bounds, and therefore also limit the possible hadronic contribution to the EGRB. Together with an independent estimate on the total contribution of AGN jets to the EGRB, this may allow one to constrain the overall cosmic ray content of AGN jets in the near future.

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