Neutrino bounds on astrophysical sources and new physics

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Ultrahigh energy cosmic neutrinos are incisive probes of both astrophysical sources and new TeV-scale physics. Such neutrinos would create extensive air showers deep in the atmosphere. The absence of such showers implies upper limits on incoming neutrino fluxes and cross sections. Combining the exposures of the Akeno Giant Air Shower Array, the largest existing ground array, with the exposure of the Fly's Eye fluorescence detector integrated over all its operating epochs, we derive 95% confidence level bounds that substantially improve existing limits. We begin with model-independent bounds on astrophysical fluxes, assuming standard model cross sections, and model-independent bounds on new physics cross sections, assuming a conservative cosmogenic flux. We then derive model-dependent constraints on new components of neutrino flux for several assumed power spectra, and we update bounds on the fundamental Planck scale M_D in extra dimension scenarios from black hole production. For large numbers of extra dimensions, we find $M_D > 2.0$ (1.1) TeV for $M_{\rm BH}^{\rm min} = M_D$ ($5M_D$), comparable to or exceeding the most stringent constraints to date.

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I. INTRODUCTION

Cosmic neutrinos provide a unique window on astrophysical processes because they escape from dense regions and typically propagate to the Earth unhindered [1]. At ultrahigh energies, they also provide an important probe of new ideas in particle physics. In contrast with all other standard model (SM) particles, their known interactions are so weak that new physics may easily alter neutrino properties, sometimes drastically. This is especially relevant for neutrinos with energies above 10^6 GeV, which interact with nucleons with center-of-mass energies above 1 TeV, where the SM is expected to be modified by new physics.

The signal for ultrahigh energy neutrinos is quasihorizontal giant air showers initiated deep in the atmosphere [2]. This signal is wellstudied and easily differentiated from air showers initiated by hadrons. The Earth's atmospheric depth rises from about 1000 g/cm^2 vertically to nearly 36000 g/cm^2 horizontally. For all but the most extreme (and typically problematic [3,4]) neutrino cross sections, the mean free path of neutrinos is larger than even the horizontal atmospheric depth. Neutrinos therefore interact with roughly equal probability at any point in the atmosphere and may initiate showers in the volume of air immediately above the detector. These will appear as ''normal'' showers, with large electromagnetic components, curved fronts (a curvature radius of a few km), and signals well spread over time (of the order of microseconds).

On the other hand, hadrons have interaction lengths of the order of 40 g/cm^2 and so always interact at the top of the atmosphere. The electromagnetic component of an air shower has mean interaction length $\sim 45-60$ g/cm^2. For a quasi-horizontal shower initiated by an ordinary hadron, then, this component is absorbed long before reaching the ground, as it has passed through the equivalent of several vertical atmospheres—two at a zenith angle of 60°, three at 70°, and six at 80°. In these showers, only high energy muons created in the first few generations of particles survive past two equivalent vertical atmospheres. The shape of the resulting shower front is therefore very flat (with curvature radius above 100 km), and its time extension is very short (less than 50 ns).

These shower characteristics are exploited by both ground arrays and fluorescence detectors in searches for primary cosmic ray neutrinos. At present, no ultrahigh energy neutrino signal has been reported. Here we determine the total exposure for neutrino detection from existing facilities and derive both model-independent and modeldependent bounds on astrophysical neutrino fluxes and new neutrino interactions.

The outline of the paper is as follows. In Sec. II we examine acceptances for neutrino detection and compute the current combined total exposure using all available data from the Akeno Giant Air Shower Array (AGASA) [5] and Fly's Eye [6] experiments. In Sec. III we determine model-independent bounds on the total neutrino flux, assuming SM cross sections. To derive model-independent results, we assume only that fluxes are confined to a small window around some central neutrino energy and obtain bounds as a function of this central energy. After that, we assume a power law neutrino flux $d\Phi/dE_\nu \propto E_\nu^{-\gamma}$ to obtain stronger, but more model dependent, bounds on the total neutrino flux from integrating over all energies.

In Sec. IV we derive model-independent bounds on high energy neutrino cross sections, assuming a conservative cosmogenic flux. These significantly improve existing limits [7]. We then derive model-dependent bounds on cross sections, focusing on the example of TeV-scale gravity scenarios in Sec. V. We improve existing constraints [8] on the fundamental Planck scale from the non-observation of microscopic black hole production by cosmic neutrinos [9–13]. For large numbers of extra dimensions, these bounds are comparable to or exceed all existing bounds on extra dimensions. Section VI contains our conclusions.

II. NEUTRINO EXPOSURE

To estimate the sensitivity of a ground-based detector to neutrino-initiated showers, we must first compute its exposure for various types of showers. The exposure is the product of the effective aperture and the range of depths within which the shower must originate to trigger the device, integrated over time. The effective aperture is the detector's projected area weighted by detection probability and integrated over solid angle. The exposure is, then,

$$\mathcal{E}(E_{\rm sh}) \approx \int_0^T dt \int_0^{h_{\rm max}} (A\Omega)_{\rm eff}(E_{\rm sh},t) \frac{\rho_0}{\rho_{\rm water}} e^{-h/H} dh, \quad (1)$$

where T is the total observation time of the detector, $h_{\rm max} = 15$ km, $H \approx 8$ km, and $\rho_0 \approx 1.15 \times 10^{-3} \rho_{\rm water}$ is the density of the atmosphere at ground level. The effective aperture is [14]

$$(A\Omega)_{\rm eff}(E_{\rm sh},t) \equiv \int_{\theta_{\rm min}}^{\theta_{\rm max}} A(t) \mathcal{P}(E_{\rm sh},\theta,t) 2\pi \sin\theta d\theta \quad (2)$$

where $A(t)$ is the detector's area, $\mathcal{P}(E_{\rm sh},\theta,t)$ is the probability that a shower that arrives with energy $E_{\rm sh}$ and zenith angle θ triggers the detector, and the angular cuts $\theta_{\rm min}$ and $\theta_{\rm max}$ are chosen to optimize detection efficiency while eliminating hadronic background.

Exposures at a given detector depend on detection method and shower type. We will refer to showers as hadronic or electromagnetic (EM), depending on the nature of their first interaction, irrespective of their later development. With this convention, for example, the SM neutral current process $\nu N \to \nu X$ produces a hadronic shower, while the charged current process for *electron* neutrinos $\nu_e N \to eX$ produces both a hadronic shower and an EM shower.

For ground arrays, exposures for hadronic and EM showers differ [15]. In hadronic showers, the initial hadronic interaction produces a strong muon component that remains until the shower reaches the ground. This muon component is largely absent for EM showers. These muons significantly enhance triggering efficiencies for ground arrays, which are sensitive only to ground level activity, and so exposures for hadronic showers exceed exposures for EM showers at ground arrays. Above some critical energy, which depends on the total effective area of the array, the detector exposure saturates. Exposures for hadronic and EM showers for the AGASA ground array [5] are given in Fig. 1. These exposures are for 1.5×10^8 s of live time between December 1995 and November 2000. The hadronic exposure was derived in Ref. [8], based on results from searches for deeply penetrating showers [16] and conservative assumptions. We derive the EM exposure by comparison with results for the Auger experiment. Effective apertures for Auger have been calculated in Ref. [15]. For AGASA, we adopt the Auger aperture for quasi-horizontal EM showers with axes falling in the array, reduced by a factor of 30, the ratio of surface areas of the arrays. We have checked that, using this EM exposure, we reproduce AGASA's bounds [17] on ν_e fluxes to within 20%.

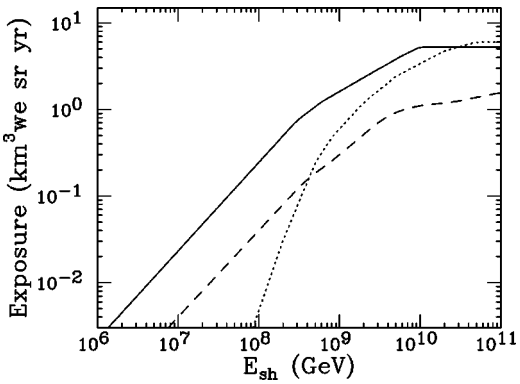

FIG. 1. The figure shows the total exposures for hadronic showers at AGASA (solid), EM showers at AGASA (dashed), and EM and hadronic showers at Fly's Eye (dotted) as functions of shower energy.

In contrast to ground arrays, the hadronic and EM exposures of fluorescence detectors are very similar. Fluorescence detectors are sensitive to the total EM activity along the entire longitudinal development of the shower. In EM showers, essentially all of the energy produces EM activity. In hadronic showers, hadronic collisions produce equal numbers of π^0, π^+, and π^-. The π^0 decays to photons, producing EM energy. However, the charged pions typically interact before decaying. Through successive interactions, most of their energy also becomes EM activity. As a result, roughly 90% of the energy in hadronic showers is EM. The response and efficiency of fluorescence detectors to hadronic and EM showers is therefore expected to be similar [18]. In this work we adopt the total Fly's Eye exposure reported in Ref. [6] for both hadronic and EM showers. This exposure, given in Fig. 1, includes not only data from the first epoch of observation (February 1983 to May 1985), most of which were reported in Ref. [18], but also data from four additional running periods between November 1985 to July 1992. The additional periods enhance the total Fly's Eye exposure by roughly a factor of 3.

The AGASA Collaboration has searched for quasi-horizontal showers that are deeply penetrating, with depth at shower maximum $X_{\rm max} > 2500$ g/cm^2 [17]. At AGASA, the location of the shower maximum is determined through its correlation to two measurable quantities: η, which parametrizes the lateral distribution of charged particles at ground level, and δ, which parametrizes the curvature of the shower front. Deeply penetrating events must satisfy $X^{\eta}_{\rm max}, X^{\delta}_{\rm max} \geqslant 2500$ g/cm^2. The expected background from hadronic showers is $1.72^{+0.14+0.65}_{-0.07-0.41}$, where the first uncertainty is from Monte Carlo statistics, and the second is systematic. Among the six candidate events, five have values of $X^{\eta}_{\rm max}$ and/or $X^{\delta}_{\rm max}$ that barely exceed 2500 g/cm^2, and are well within $\Delta X_{\rm max}$ of this value, where $\Delta X_{\rm max}$ is the estimated precision with which $X_{\rm max}$ can be reconstructed. The AGASA Collaboration thus concludes that there is no significant enhancement of deeply penetrating shower rates given the detector's resolution.

The Fly's Eye observes an air shower as a nitrogen fluorescence light source traveling at the speed of light through

the atmosphere [6]. The light is emitted isotropically with intensity proportional to the number of charged particles in the shower. The received light profile is reconstructed by a three-parameter fit to the charged particle density (largely electrons and positrons) along the shower track. The parameters are the depth of the observed first interaction X_0; the depth of the shower maximum $X_{\max}$; and the density $N_{\max}$ at $X_{\max}$. In a running time of 11 years, Fly's Eye recorded more than 5000 events. However, there are no neutrino candidates with shower maximum deeper than 2500 g/cm^2 [6,18,19].

All in all, given one event that unambiguously passes all cuts with 1.72 events expected from hadronic background, the combined results of AGASA and Fly's Eye imply an upper bound of 3.5 events at 95% confidence level (C.L.) [20] from neutrino fluxes.

The event rate for quasi-horizontal deep showers from ultrahigh energy neutrinos is

$$N=\sum_{i,X}\int dE_i N_A \frac{d\Phi_i}{dE_i}\sigma_{iN\to X}(E_i)\mathcal{E}_{iX}(E_i), \tag{3}$$

where the sum is over all neutrino species $i=\nu_e,\bar{\nu}_e,\nu_\mu,\bar{\nu}_\mu,\nu_\tau,\bar{\nu}_\tau$, and all final states X. $N_A=6.022\times10^{23}$ is Avogadro's number, and $d\Phi_i/dE_i$ is the source flux of neutrino species i. Finally, $\mathcal{E}_{iX}(E_i)$ is the appropriate exposure measured in cm^3 we sr$\cdot$time.

To clarify, we present appropriate exposures for some SM processes. At the ultrahigh energies of interest, on average 20% of the total neutrino energy goes into hadronic recoil for both SM neutral current and charged current events [21]. Exposures for SM charged current events at AGASA are therefore

$$\mathcal{E}_{\nu_e X}(E_{\nu_e})=\min\{\mathcal{E}_{\rm had}(0.2E_{\nu_e})+\mathcal{E}_{\rm EM}(0.8E_{\nu_e}),\mathcal{E}_{\rm sat}\} \tag{4}$$

$$\mathcal{E}_{\nu_\mu X}(E_{\nu_\mu})=\mathcal{E}_{\rm had}(0.2E_{\nu_\mu}) \tag{5}$$

$$\mathcal{E}_{\nu_\tau X}(E_{\nu_\tau})=\mathcal{E}_{\rm had}(0.2E_{\nu_\tau}), \tag{6}$$

with identical expressions for anti-neutrinos. The exposures for ν_μ and ν_τ are identical, because at these energies, τ leptons typically do not decay before arriving at the Earth's surface. For ν_e, the expression includes the effects of saturation, noted previously. For AGASA, the exposure saturates at $\mathcal{E}_{\rm sat}\approx5.3$ km^3 we sr yr at $E_{\rm sh}\approx10^{10}$ GeV (see Fig. 1). For neutral current interactions at AGASA, all of the right-hand entries are $\mathcal{E}_{\rm had}(0.2E_{\nu_i})$.

For the Fly's Eye experiments, the corresponding charged current exposures are simply

$$\mathcal{E}_{\nu_e X}(E_{\nu_e})=\mathcal{E}(E_{\nu_e}) \tag{7}$$

$$\mathcal{E}_{\nu_\mu X}(E_{\nu_\mu})=\mathcal{E}(0.2E_{\nu_\mu}) \tag{8}$$

$$\mathcal{E}_{\nu_\tau X}(E_{\nu_\tau})=\mathcal{E}(0.2E_{\nu_\tau}). \tag{9}$$

For neutral current exposures at the Fly's Eye experiments, all of the right-hand entries are $\mathcal{E}(0.2E_{\nu_i})$.

Given the exposures $\mathcal{E}$ described above, the absence of events implies upper bounds on cross sections, assuming some fixed flux, or upper bounds on fluxes, assuming some fixed cross section. We now consider these two possibilities in turn.

III. BOUNDS ON ASTROPHYSICAL NEUTRINO FLUXES

To derive bounds on the neutrino flux, we assume SM charged and neutral current interactions for all neutrinos. We further assume that the source flux of neutrinos, which in the energy range of interest is dominantly ν_μ, $\bar{\nu}_\mu$, and ν_e at production, is completely mixed in flavor upon arrival at Earth. There is now strong evidence for maximal mixing among all neutrino species [22]. In addition, given mass differences of $\Delta m^2\gtrsim10^{-6}$ eV^2, the neutrino flux is expected to be completely mixed if they travel a distance $\gtrsim0.1$ Mpc. The assumption of equal flavor representation upon arrival is therefore strongly supported by data [23].

We first derive model-independent bounds on the total neutrino flux. Let us start by noting that if the number of events integrated over energy is bounded by 3.5, then it is certainly true bin by bin in energy. Thus, using Eq. (3) we obtain

$$\sum_{i,X}\int_\Delta dE_i N_A \frac{d\Phi_i}{dE_i}\sigma_{iN\to X}(E_i)\mathcal{E}_{iX}(E_i)<3.5, \tag{10}$$

at 95% C.L. for some interval Δ. Here the sum over X takes into account charge and neutral current processes. In a logarithmic interval Δ where a single power law approximation

$$\frac{d\Phi_i}{dE_i}\sigma_{iN\to X}(E_i)\mathcal{E}_{iX}(E_i)\sim E_i^\alpha \tag{11}$$

is valid, a straightforward calculation shows that

$$\int_{\langle E\rangle e^{-\Delta/2}}^{\langle E\rangle e^{\Delta/2}}\frac{dE_i}{E_i}E_i\frac{d\Phi_i}{dE_i}\sigma_{iN\to X}(E_i)\mathcal{E}_{iX}(E_i)$$
$$=\langle\sigma_{\nu_i N\to X}(E_i)\mathcal{E}_{iX}(E_i)E_i d\Phi_i/dE_i\rangle\frac{\sinh\delta}{\delta}\Delta, \tag{12}$$

where $\delta=(\alpha+1)\Delta/2$ and $\langle A\rangle$ denotes the quantity A evaluated at the center of the logarithmic interval. The parameter $\alpha=0.363+\beta-\gamma$, where the 0.363 is the power law index of the SM neutrino cross sections [21], and β and $-\gamma$ are the power law indices (in the interval Δ) of the exposure and flux $d\Phi_i/dE_i$, respectively. Since $\sinh\delta/\delta>1$, a conservative bound may be obtained from Eqs. (10) and (12):

$$N_A\sum_{i,X}\langle\sigma_{\nu_i N\to X}(E_i)\rangle\langle\mathcal{E}_{iX}(E_i)\rangle\langle E_i d\Phi_i/dE_i\rangle<3.5/\Delta. \tag{13}$$

In this work we choose $\Delta=1$, corresponding to one e-folding of energy, as a likely interval in which the single power law behavior is valid. By setting $\langle E_i d\Phi_i/dE_i\rangle=\frac{1}{6}\langle E_\nu d\Phi_\nu/dE_\nu\rangle$, where Φ_ν is the total neutrino flux, we

TABLE I. Model-independent upper limits on the differential neutrino flux at 95% C.L.

E_ν (GeV)	$\langle E_\nu d\Phi_\nu/dE_\nu\rangle$ (km^{-2} sr^{-1} yr^{-1})
1×10^8	1.8×10^5
3×10^8	4.1×10^4
1×10^9	7.9×10^3
3×10^9	2.2×10^3
1×10^{10}	5.0×10^2
3×10^{10}	1.6×10^2
1×10^{11}	6.8×10^1

obtain model-independent upper limits on the total neutrino flux at 95% CL. The results are given in Table I.

These model-independent upper bounds on the total neutrino flux can be strengthened by assuming a particular flux behavior. For example, if the neutrino flux falls like

$$\frac{d\Phi_\nu}{dE_\nu}=J_0\left(\frac{E_\nu}{E_0}\right)^{-\gamma}, \tag{14}$$

Eq. (10) leads to $J_0<1.4\times10^{-5}$ km^{-2} sr^{-1} yr^{-1} GeV^{-1} for $\gamma=2$ and $E_0=10^8$ GeV. Under the same assumptions for $\gamma=1.5$ one obtains $J_0<9.8\times10^{-7}$ km^{-2} sr^{-1} yr^{-1} GeV^{-1}. Figure 2 shows both the model-independent bounds of Table I as well as the bounds on flux under the power law assumptions just discussed. Additionally displayed in the figure are the upper limits on the $\nu_\mu+\nu_e$ flux obtained from the non-observation of microwave Čerenkov pulses from EM showers induced by neutrinos in the Moon's rim, as measured by the Goldstone Lunar Ultrahigh energy neutrino Experiment (GLUE) [24], as well as bounds from the search for radio pulses from EM showers created by electron neutrino collisions in ice at the Radio Ice Čerenkov Experiment (RICE) [25]. Comparing model-independent bounds to model-independent bounds, and model-dependent bounds to model-dependent bounds, we find that the bounds obtained here are significantly more stringent than existing limits.

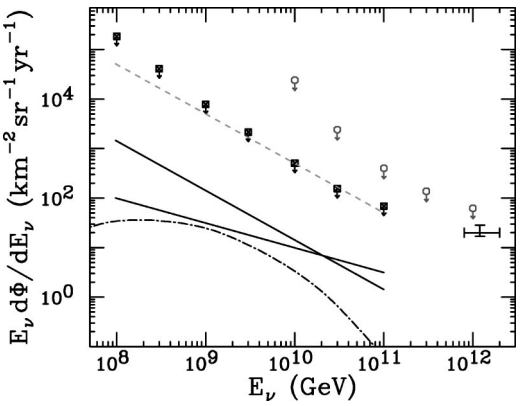

FIG. 2. 95% C.L. model-independent (squares) and model-dependent (solid lines) upper limits on the total neutrino flux $E_\nu d\Phi_\nu/dE_\nu$ derived here using the combined exposures of AGASA and Fly's Eye. The model-dependent bounds assume $d\Phi_\nu/dE_\nu\propto E^{-\gamma}$, with $\gamma=1.5$ and 2. For comparison, we also present limits on the $\nu_\mu+\nu_e$ from GLUE (circles) [24] and the 95% C.L. bound from RICE on the ν_e flux assuming a power law spectrum with $\gamma=2$ (dashed) [25]. The dash-dotted line is the predicted total cosmogenic flux from pion photo-production [26], and the point with error bars is the total neutrino flux required by *Z*-burst models, as derived in [27].

Also displayed in Fig. 2 as a single point with error bars is the total neutrino flux required by *Z* bursts [28]. The *Z*-burst model proposes that the observed ultrahigh energy cosmic rays with energies above 10^{11} GeV are secondaries resulting from resonant annihilation of ultrahigh energy neutrinos on relic neutrinos at the *Z* pole [29]. A recent analysis [27] includes several possibilities for a diffuse background of protons, as distinct from protons resulting from the *Z*-burst itself. The point shown in the figure corresponds to the case in which the proton background originates at distances $\lesssim50$ Mpc. The normalized Hubble expansion rate, the matter and vacuum energy densities, and the maximum redshift are taken as $h=0.71$, $\Omega_M=0.3$, $\Omega_\Lambda=0.7$, and $z_{\max}=2$, respectively. The neutrino flux is assumed to have no cosmological evolution. The horizontal errors result from the 1σ uncertainty in the neutrino mass determination. The errors in the flux reflect the statistical fluctuations in the fits, as well as the uncertainty in the Hubble expansion rate.

A more speculative explanation of the mysterious events at the high end of the spectrum assumes that the cosmic ray primaries arise in the decay of massive elementary *X* particles. Sources of these exotic particles could be either topological defects left over from early universe phase transitions ($m_X\sim10^{16}-10^{19}$ GeV) [30], or some long-lived metastable superheavy ($m_X\gtrsim10^{12}$ GeV) relic particles produced through vacuum fluctuations during the inflationary stage of the universe [31]. The *X* particles typically decay to leptons and quarks. The latter produce jets of hadrons containing mainly pions, together with a 3% admixture of nucleons. The predicted spectrum would thus be dominated by gamma rays and neutrinos produced via pion decay. The neutrino flux bounds derived in this work therefore seriously constrain this type of model. Moreover, recent analyses of Haverah Park data [32] suggest that less than 50% of the primary cosmic rays above 4×10^{10} GeV can be photons at 95% C.L. Note that mechanisms which successfully deplete the high energy photons (such as efficient absorption on the universal and galactic radio background) require an increase in the neutrino flux to maintain the overall normalization of the observed spectrum [33]. Definite quantitative comparison with the neutrino flux bounds presented here can be obtained by specifying the nature of the decay of the *X* particle.

IV. BOUNDS ON NEW PHYSICS INTERACTIONS

In this section we examine the potential of probing the big desert that lies between the electroweak and grand unified theory (GUT) scales using ultrahigh energy neutrino interactions. To derive bounds on possible new physics contributions to neutrino cross sections, we assume the ''guaranteed'' flux of cosmogenic neutrinos arising from pion photoproduction from ultrahigh energy protons propagating through the cosmic microwave background. This flux de-

pends on the cosmological evolution of the cosmic ray sources. Throughout this work, we conservatively adopt the estimates of Protheroe and Johnson [26] with nucleon source spectrum scaling as $d\Phi_{\rm nucleon}/dE \propto E^{-2}$ and extending up to the cutoff energy $10^{12.5}$ GeV. We assume also a cosmological source evolution scaling as $(1+z)^4$ for redshift $z<1.9$ [34]. The total ultrahigh energy cosmogenic neutrino flux is shown in Fig. 2.

We consider this flux highly conservative. For example, one might conjecture that the cosmogenic flux is absent because the observed cosmic rays with energies $\gtrsim 10^{11}$ GeV are protons or nuclei generated by nearby sources within 50 Mpc. However, none of the known nearby candidate sources, such as Virgo, M82, Centaurus A, and galactic pulsars [35] is as powerful as more distant sources, such as Cygnus A and Pictor A [36]. The latter must therefore inject nucleons with energies of at least 10^{12} GeV. Indeed, contributions to the nucleon channel from semi-local sources [37], as well as recently discussed possibilities [38] of source spectra harder than $E_\nu^{-1.5}$ and source evolutions stronger than $(1+z)^4$, would all significantly enhance the cosmogenic neutrino flux. Contributions from decaying topological defects, active galactic nuclei, and other speculative sources would have a similar effect. If realized in nature, any one of these possibilities would strengthen our bounds.

For SM interactions and the cosmogenic neutrino flux given in Fig. 2, the expected rate for deeply penetrating showers at AGASA and Fly's Eye is about 0.02 events per year, and so negligible. New physics may contribute to neutrino cross sections in a multitude of ways [9,11,39,40], with different contributions for different flavors, and a variety of final states producing showers with a variety of hadronic and EM shower components. Here we assume that the new physics is flavor-blind, inducing equivalent cross section modifications for all neutrino species, and that the resulting new physics final state leads to showers with a negligible EM component. We then consider two simple but representative cases: (1) $y=1$ and (2) $y=0.1$, where $y \equiv E_{\rm sh}/E_\nu$ is the average inelasticity. In case (1), the shower energy $E_{\rm sh}$ differs little from the neutrino energy E_ν, as in TeV-scale black hole production (see below). In case (2), $E_{\rm sh}$ is substantially less than E_ν. This holds, for example, when neutral current interactions are enhanced by the exchange of Kaluza-Klein (KK) gravitons [4].

Our starting point to derive model-independent bounds on the total neutrino-nucleon cross section $\sigma_{\nu N}$ is again Eq. (13). For reasons given previously, we choose $\Delta=1$, so that Eq. (13) becomes

$$N_A \langle \sigma_{\nu N \to X}(E_\nu) \rangle \langle \mathcal{E}(yE_\nu) \rangle \langle E_\nu d\Phi_\nu/dE_\nu \rangle < 3.5. \quad (15)$$

Using the exposures in Fig. 1 and the cosmogenic flux in Fig. 2, we find model-independent bounds on the neutrino cross section for different energies and the two inelasticities above. The resulting limits are given in Table II and shown in Fig. 3. For reference, the SM cross sections [21] are also given in the figure.

These bounds assume that neutrinos produce deeply penetrating quasi-horizontal showers. They may be avoided if neutrinos are so strongly interacting that they shower high in the atmosphere. For zenith angles of 70°, this requires an interaction length below 3000 g/cm^2, corresponding to a cross section above 6×10^8 pb. If the interaction length is between 3000 g/cm^2 and the horizontal atmospheric depth of 36000 g/cm^2, corresponding to a νN cross section between 6×10^8 pb and 5×10^7 pb, respectively, an exponential decrease in the event rate for showers of a given total energy will be observed beyond a critical angle. Since this range includes our upper bounds on $\sigma_{\nu N}$ at 10^{11} GeV, it is conceivable that $\sigma_{\nu N}$ could thus be measured directly. Bounds on anomalous νN cross sections which are near-hadronic strength can be obtained through considerations of their feedback to low energy physics via dispersion relations [41].

TABLE II. 95% C.L. model-independent upper limits on the neutrino-nucleon cross section.

E_ν (GeV)	$\sigma_{\nu N}^{y=1}$ (pb)	$\sigma_{\nu N}^{y=0.1}$ (pb)
1×10^{10}	1.8×10^6	7.4×10^6
3×10^{10}	8.8×10^6	2.2×10^7
1×10^{11}	8.1×10^7	1.1×10^8

The bounds shown in Table II strengthen previous bounds [7] by roughly one order of magnitude. This enhancement results from a number of factors. First, the exposure used in Ref. [7] was limited to the first 6×10^6 s of the Fly's Eye operation. The updated exposure used here is roughly 9 times larger. Second, the cosmogenic flux used in Ref. [7] assumes a source cutoff energy of $10^{11.5}$ GeV, and is smaller by roughly a factor of 4 than the one used here. We find it unnatural to assume that the the source cutoff coincides with the maximum observed cosmic ray energy, and so have used the flux corresponding to the larger cutoff energy. Finally, we have presented 95% C.L. limits, corresponding to a limit of 3.5 events.

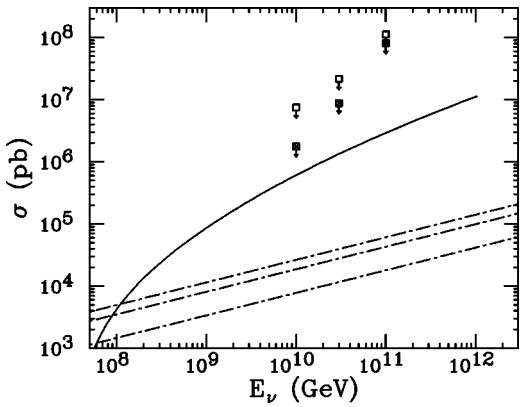

FIG. 3. 95% C.L. model-independent upper limits on $\sigma_{\nu N}$ for inelasticity $y=1$ (filled squares) and $y=0.1$ (open squares). The solid contour is the upper limit on the black hole production cross section for $n=7$ and $x_{\rm min}=5$, corresponding to $M_D=1.1$ TeV. For comparison, the SM neutral current, charged current, and total neutrino nucleon cross sections are also given (dot-dashed, from below). These bounds assume that the neutrino cross sections does not exceed 6×10^8 pb (see text).

As in the case of flux bounds, the model-independent cross section bounds of Table II will be improved if one assumes a cross section shape and so can integrate over all energies. We consider a particular example in the next section.

V. IMPLICATIONS FOR TeV-SCALE GRAVITY

The idea that our universe could be a brane embedded in some higher dimensional world has received a great deal of renewed attention over the last 5 years [42,43]. From a phenomenological point of view, this possibility presents a new perspective on the hierarchy between the gravitational and electroweak mass scales. In these scenarios, the effective 4-dimensional Planck scale $M_{\rm Pl}\sim 10^{19}$ GeV is determined by the fundamental $(4+n)$-dimensional Planck scale $M_D \sim 1$ TeV and the geometry of the n extra dimensions.

Arguably the most fascinating prediction of TeV-scale gravity is the production of black holes (BHs) in observable collisions of elementary particles [44,45]. For cosmic rays, this implies that ultrahigh energy neutrinos may produce BHs in the atmosphere, initiating deep quasi-horizontal showers far above SM rate [9]. BH production therefore provides a specific example of a model-dependent cross section that is bounded by the arguments discussed above.

TeV-scale gravity also has a number of other implications for cosmic rays. The implications of perturbative KK graviton exchange has been considered in a number of scenarios [39]. However, for cosmic rays, in contrast to the case at colliders, there is an abundance of center-of-mass energy. Extra dimensional effects will therefore first appear as nonperturbative BH production in processes with center-of-mass energies above the fundamental Planck scale, rather than through perturbative effects below the Planck scale. This is in stark contrast to the case at colliders, where the sensitivity to KK graviton effects surpasses the sensitivity to BH production.[1]

The sensitivity of current cosmic ray experiments to BH production, as well as that of facilities expected in the not-too-distant future, has been thoroughly investigated [8,9,11–13]. The parton-parton cross section is estimated from the geometric area of the BH horizon and is of order $\hat{\sigma}_i \sim \pi r_s^2$ [44], where

$$r_s(M_{\rm BH}) = \frac{1}{M_D}\left[\frac{M_{\rm BH}}{M_D}\right]^{1/(1+n)} \left[\frac{2^n \pi^{(n-3)/2}\Gamma\left(\frac{n+3}{2}\right)}{n+2}\right]^{1/(1+n)} \tag{16}$$

is the radius of a Schwarzschild BH in $4+n$ dimensions [46]. Criticisms of the absorptive black disk scattering amplitude, which center on the exponential suppression of transitions involving a (few-particle) quantum state to a (many-particle) semiclassical state [47], have been addressed in Refs. [48,49]. The geometric cross section applies for both flat and hyperbolic [50] extra dimensions that are larger than the Schwarzschild radius, and for warped extra dimensions where r_s is small compared to the curvature scale of the geometry associated with the warped subspace [51,52].

The total production cross section for BHs with mass $M_{\rm BH}\equiv\sqrt{sx}$ is then [9]

$$\sigma_{\nu N\to {\rm BH}}(E_\nu) = \sum_i \int_{(M_{\rm BH}^{\rm min})^2/s}^1 dx \hat{\sigma}_i(\sqrt{xs}) f_i(x,Q), \tag{17}$$

where $s=2m_N E_\nu$, x is the parton momentum fraction, the f_i are parton distribution functions (PDFs), $M_{\rm BH}^{\rm min}$ is the minimum BH mass, and the sum is carried out over all partons in the nucleon. The choice of the momentum transfer Q is governed by considering the time or distance scale probed by the interaction. According to Thorne's hoop conjecture [53], the formation of a well-defined horizon in four dimensions occurs when the colliding particles are at a distance $\sim r_s$ apart. (Note that there could be an n-dependent factor for higher number of dimensions.) This has led to the advocacy of the choice $Q\simeq r_s^{-1}$ [11], which has the advantage of a sensible limit at very high energies. However, as has been pointed out by Dimopoulos and Emparan [48], string progenitors can give experimental signals akin to BH for the collision energies and values of $M_{\rm BH}^{\rm min}$ under present consideration. In cosmic ray experiments, these signals are indistinguishable from those of BH decay (more on this below). Detailed calculations [54] show that these ''string ball'' cross sections can exceed BH cross sections. In this region, the dual resonance picture of string theory would suggest a choice $Q\sim M_{\rm res}\sim\sqrt{xs}$. Since we have chosen to use only BH cross sections over the entire energy region, we set $Q=\min\{M_{\rm BH},10\ {\rm TeV}\}$, where the upper limit is from the CTEQ5M1 distribution functions [55]. We are aware that for $Q\gtrsim$ string scale, the PDFs will receive significant corrections from the rapid increase of degrees of freedom. Fortunately, as noted in Ref. [8], the cross section $\sigma_{\nu N\to {\rm BH}}$ is largely insensitive to the details of the choice of Q. For example, the two choices discussed here result in cross sections that differ by only 10% to 20%.

Once produced, BHs will Hawking evaporate with a temperature proportional to the inverse radius $T_H=(n+1)/(4\pi r_s)$. The wavelength $\lambda=2\pi/T_H$ corresponding to this temperature is larger than the BH size. Hence, to first approximation the BH behaves like a point radiator with entropy

$$S=\frac{4\pi M_{\rm BH} r_s}{n+2} \tag{18}$$

and mean lifetime

$$\tau_{\rm BH}\sim\frac{1}{M_D}\left(\frac{M_{\rm BH}}{M_D}\right)^{(3+n)(1+n)}. \tag{19}$$

Microscopic black holes therefore decay almost instantaneously to a thermal distribution of SM particles. As very few SM particles are invisible to cosmic ray detectors, the neutrino energy is almost entirely transformed into shower energy. The EM component of these showers differs substan-

[1]We thank S. Dimopoulos for emphasizing this point.

tially from that of SM neutrino interactions, allowing a good characterization of the phenomenon against background when the BH entropy$\gg 1$ [10]. Recently, it has been noted that BH recoil induced by KK-graviton emission may launch the BH out of the brane [56]. In this case, BH radiation would be prematurely terminated from the perspective of brane observers. This effect would drastically deplete the rate of deeply developing showers, and could be misinterpreted as a sharp cutoff on the ultrahigh energy neutrino spectrum. We assume here that the effect of recoil is negligible.

An important parameter in determining the BH cross section is $x_{\rm min} \equiv M_{\rm BH}^{\rm min}/M_D$, the ratio of the minimal black hole mass to the fundamental Planck scale. The above description of BH production and decay relies on semi-classical arguments, valid for large $x_{\rm min}$ or, equivalently, large BH entropy. For large $x_{\rm min}$, thermal fluctuations due to particle emission are small ($S \gg 1$) [57], statistical fluctuations in the microcanonical ensemble are small ($\sqrt{S} \gg 1$), and quantum gravity effects may be safely neglected. In addition, gravitational effects of the brane on BH production, which are ignored in all analyses to date, are expected to be insignificant for BH masses well above the brane tension, which is presumably $\sim M_D$. None of these is true for $x_{\rm min} \approx 1$.

Cosmic ray experiments are largely insensitive to the exact details of BH decay. Whatever happens near $x_{\rm min} \approx 1$, it seems quite reasonable to expect that BHs or their Planck mass progenitors will decay visibly, triggering deeply atmospheric cascade developments. There is, however, sensitivity to the BH production cross section. In string theory, BH production is expected to gradually pass to the regime of string ball production as $M_{\rm BH}$ approaches M_D. Evidence from this picture suggests that the BH production cross section is not radically altered in this limit [48]. This addresses many of the concerns listed above, but does not address the problem of brane effects on BH production—these may still be large for $x_{\rm min} \approx 1$. In our analysis, we avoid choosing a specific $x_{\rm min}$; rather, we present results for the generous range $1 \leqslant x_{\rm min} \leqslant 10$. As we will see, the bounds are rather insensitive to $x_{\rm min}$, in contrast to the case at colliders such as the CERN Large Hadron Collider (LHC).

In Fig. 4 we show the lower limit on M_D as a function of $x_{\rm min}$ corresponding to <3.5 events (95% C.L.) observed for the combined exposures of AGASA and Fly's Eye. Bounds on M_D from tabletop gravity experiments, as well as from astrophysical and cosmological considerations, greatly exceed 1 TeV, for models with $n \leqslant 4$ flat extra dimensions [58]. For $n > 4$, however, the bounds of Fig. 4 are among the most stringent to date. Note that these bounds do not depend on the shape of the extra dimensions and are valid for both warped and non-warped scenarios [52]. For $x_{\rm min}=1$, the bounds extend up to 2.0 TeV for $n=7$. Moreover, assuming $x_{\rm min}=3$, for which the entropy $S>10$, the bounds derived with the combined exposure, for $n=5,6,7$, are $M_D > 1.26$ TeV, 1.30 TeV, 1.40 TeV, respectively. All of these exceed bounds from the Fermilab Tevatron and CERN e^+e^- collider LEP [59], even in the case where the brane softening

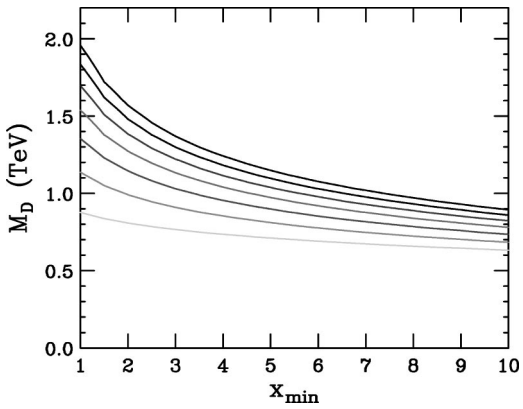

FIG. 4. 95% C.L. lower limits on the fundamental Planck scale as a function of $x_{\rm min}$ for $n=1,\ldots,7$ extra dimensions (from below).

parameter Λ [60] is as large as M_D (see [8] for details).[2]

In Fig. 3 we have also plotted the maximal BH cross section, corresponding to $M_D=1.1$ TeV, for the case $n=7$ and $x_{\rm min}=5$. As expected, given a model for the cross section's energy dependence, the resulting bounds on new interactions are much more stringent than the model-independent limits derived above.

Finally, as noted above, the event rates for black hole production by cosmic rays are fairly insensitive to the choice of $x_{\rm min}$. This contrasts sharply with the case at colliders. Specifically, the total production cross section of a BH of mass $M_{\rm BH} \sim \sqrt{\tau s}$ in a pp collision is given by

$$\sigma_{pp\to {\rm BH}}(\tau_{\rm min},s) = \sum_{ij} \int_{\tau_{\rm min}}^1 d\tau \int_\tau^1 \frac{dx}{x} f_i(x) f_j(\tau/x) \hat{\sigma}_{ij}, \tag{20}$$

where τ is the parton-parton center-of-mass energy squared fraction, and $\sqrt{\tau_{\rm min} s}$ is the minimum center-of-mass energy for which the black disk approximation is valid. The number of BH produced at the LHC ($\sqrt{s}=14$ TeV and luminosity $\mathcal{L}=10^{34}$ cm^{-2} s^{-1} [63]) is then $N^{\rm LHC} = \int \sigma_{pp\to {\rm BH}} \mathcal{L} dt$. In Fig. 5 we show both Auger and LHC event rates for various $x_{\rm min}$. For fixed M_D, the LHC event rates drop by one to two orders of magnitude for every unit increase in $x_{\rm min}$, while the Auger event rates are relatively stable [64]. This may be understood as resulting from a combination of the very high energies available in cosmic neutrinos and the fact that the parton energy is not degraded by a parton distribution function in the cosmic neutrino ''beam.''

VI. CONCLUSIONS

In the first part of this paper, we derived new limits on the cosmic neutrino flux striking the Earth's atmosphere. This was accomplished by searching for quasi-horizontal deeply

[2]The increased exposure used here will also strengthen existing bounds [61] from p-brane production [62] in asymmetric compactifications.

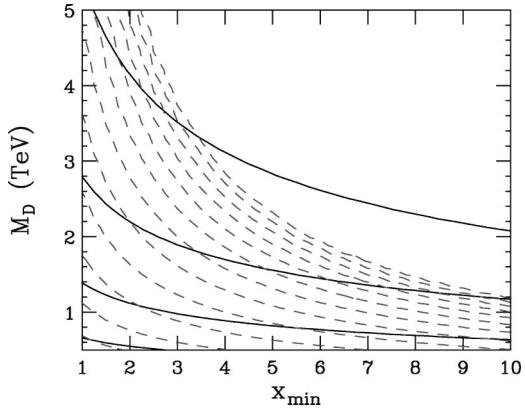

FIG. 5. Black hole events at Auger in 3 years (1, 10, 100, 1000, from above) (solid) and at the LHC for integrated luminosity 100 fb^{-1} (1, 10, 100, . . . ,10^{11}, from above) (dashed) for $n=7$ extra dimensions.

developing showers in ultrahigh energy cosmic ray data, taking into account the combined exposures of the AGASA and Fly's Eye experiments. Our results significantly strengthen existing limits and present serious problems for models where exotic elementary X particles cascade decay to cosmic ray particles. In particular, models where topological defects are responsible for the events detected with energies $\gtrsim 10^{11}$ GeV are severely constrained, because neutrinos are typically a significant component in X decays, and have a hard spectrum extending up to $M_{\mathrm{GUT}} \sim 10^{16}$ GeV. The bounds obtained in this paper will also challenge any attempt to normalize the observed spectrum to the proton flux as predicted by top down models.

In the second part of the paper, we used the atmosphere as a giant calorimeter to probe neutrino-nucleon cross sections at $\sqrt{s} \gtrsim 1$ TeV. We first combined the complete neutrino exposure of the above-mentioned facilities with the flux of cosmogenic neutrinos, to derive model-independent upper bounds on the neutrino-nucleon cross section. These bounds strengthen existing limits by roughly one order of magnitude. We then considered TeV-scale gravity models to study BH production. The upper bounds on the neutrino-nucleon cross section implied lower limits on the fundamental Planck scale, which represent the best existing limits on TeV-scale gravity for $n \geqslant 5$ extra spatial dimensions.

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