

## Exotic Neutrino Interactions at the Pierre Auger Observatory

Luis Anchordoqui,<sup>1</sup> Tao Han,<sup>2</sup> Dan Hooper,<sup>3,4</sup> and Subir Sarkar<sup>5</sup>

<sup>1</sup>*Department of Physics, Northeastern University,  
Boston, MA 02115, USA*

<sup>2</sup>*Department of Physics, University of Wisconsin,  
Madison, WI 53706, USA*

<sup>3</sup>*Particle Astrophysics Center, Fermilab,  
P.O. Box 500, Batavia, IL 60510, USA*

<sup>4</sup>*Astrophysics, University of Oxford,  
Oxford OX1 3RH, UK*

<sup>5</sup>*Rudolf Peierls Centre for Theoretical Physics, University of Oxford,  
Oxford OX1 3NP, UK*

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### Abstract

The Pierre Auger Observatory for cosmic rays provides a laboratory for studying fundamental interactions at energies well beyond those available at colliders. In addition to hadrons or photons, Auger is sensitive to ultra-high energy neutrinos in the cosmic radiation and models for new physics can be explored by observing neutrino interactions at center-of-mass energies beyond the TeV scale. By comparing the rate for quasi-horizontal, deeply penetrating air showers triggered by all types of neutrinos with the rate for slightly upgoing showers generated by Earth-skimming tau neutrinos, any deviation of the neutrino-nucleon cross-section from the Standard Model expectation can be constrained. We show that this can test models of low-scale quantum gravity (including processes such as Kaluza-Klein graviton exchange, microscopic black hole production and string resonances), as well as non-perturbative electroweak instanton mediated processes. Moreover, the observed ratios of neutrino flavors would severely constrain the possibility of neutrino decay.

## I. INTRODUCTION

The Pierre Auger Observatory is the largest cosmic ray detector in the world [1]. Currently under construction at Malargüe (Argentina), it has begun taking data and already accumulated an exposure comparable to previous experiments such as AGASA and HiRes [2].

In addition to studying the highest energy cosmic rays, Auger is also capable of observing ultra-high energy cosmic neutrinos [3]. At present, the AMANDA telescope at the South Pole holds the record for the most energetic neutrino interactions observed [4]; these events have energies up to  $\sim 10^5$  GeV and are consistent with the predicted spectrum of atmospheric neutrinos. Auger, by contrast, is expected to detect neutrinos with energies above  $\sim 10^8$  GeV. The ability to study neutrino interactions at such high energies will open a unique window on possible physics beyond the Standard Model (SM) of strong and electroweak interactions.

A variety of models have been proposed in which neutrino interactions become substantially modified at very high energies, the most interesting being models of low scale quantum gravity (involving the exchange of Kaluza-Klein (KK) gravitons, production of microscopic black holes, and excitation of TeV-scale string resonances), and models featuring non-perturbative electro-weak instanton induced interactions.

Moreover the neutrino flavor ratios predicted by the standard oscillation phenomenology can be modified in propagation over cosmological distances if processes such as neutrino decay occur. Auger is expected to detect the ‘cosmogenic’ neutrino flux from interactions of extragalactic ultra-high energy cosmic rays with the cosmic microwave background. Thus, it will be sensitive to such effects, being capable of measuring the flux of ultra-high energy tau neutrinos in addition to the overall neutrino flux.

In this article, we explore the relevant phenomenology and quantify the sensitivity of Auger to such new physics.<sup>1</sup> In Sec. II, we describe the Auger experiment and its ability to detect quasi-horizontal neutrino-induced air showers, as well as up-going showers induced by Earth-skimming tau neutrinos. In Sec. III we discuss possible sources of cosmic ultra-high energy neutrinos. In section IV, we infer the sensitivity of Auger to the neutrino-nucleon interaction cross-section and to the flavour content of the ultra-high energy cosmic neutrino flux. In Sec. V we consider specific models of physics beyond the SM and their signatures in Auger. We present our conclusions in Sec. VI.

## II. ULTRA-HIGH ENERGY NEUTRINOS AT AUGER

### A. The Pierre Auger Observatory

Auger is a hybrid ultra-high energy cosmic ray detector, with a ground array of water Cerenkov detectors sampling air shower particles, overlooked by air fluorescence detector telescopes which observe the longitudinal development of the showers [1]. When completed in 2005–06, the Southern hemisphere site will have 1600 detectors on the ground covering 3000 km<sup>2</sup>, and 4 fluorescence telescopes having 6 detectors each. A similar facility has been proposed for a Northern hemisphere site in Colorado (USA). In our calculations we consider only a single site and focus on the ground array.

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<sup>1</sup> For a review of exotic neutrino interactions and their signatures in high-energy cosmic neutrino telescopes such as IceCube, see Ref. [5].

## B. Quasi-Horizontal, Deeply Penetrating Showers

At sufficiently high energies cosmic neutrinos can trigger atmospheric air showers similar to those due to high energy cosmic rays (hadrons or photons). However, unlike ordinary cosmic ray showers which are initiated near the top of the atmosphere, those generated by neutrinos can be initiated at any depth since the interaction cross-section is much smaller, hence the probability of interaction per unit length is approximately constant. Neutrino induced showers can thus be distinguished from cosmic ray showers by requiring that they be *deeply penetrating*. This is most useful for distinguishing the two kinds of showers, because within  $\sim 20^\circ$  of the horizon the electromagnetic component of hadron-induced showers is completely absorbed before reaching the detector.

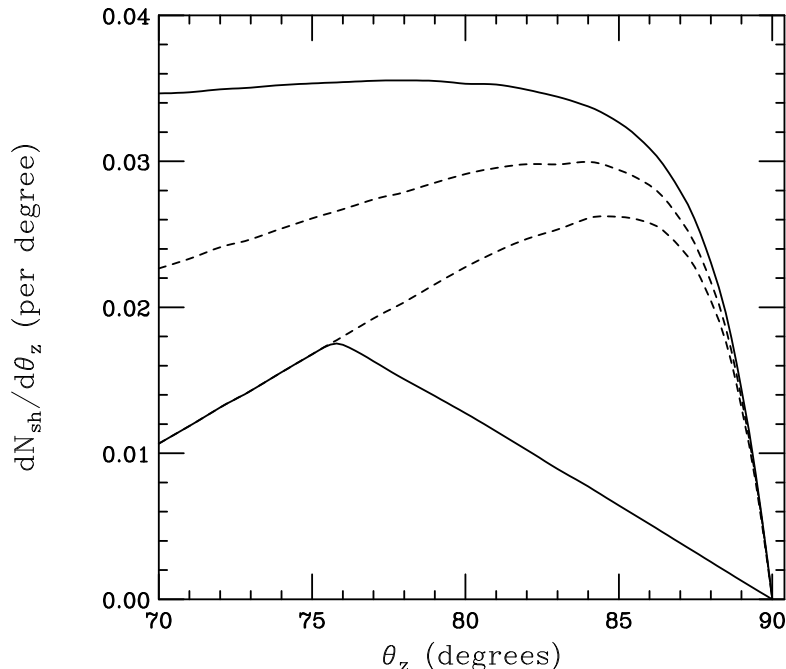


FIG. 1: Angular distribution of neutrino-induced showers expected to be observed by the Auger ground array with different selection criteria. The upper solid line indicates all showers generated by cosmic neutrinos over the range of zenith angles shown, while the upper and lower dashed lines correspond to the cases when the shower is initiated at a depth exceeding 1000 and 2000  $\text{g}/\text{cm}^2$ , respectively. The lower solid line corresponds to the case when the shower is initiated at a depth exceeding 2000  $\text{g}/\text{cm}^2$  and within 2000  $\text{g}/\text{cm}^2$  from the detector — this last set of cuts is adopted in our calculations. For the four cases shown, the fraction of events which survive the various cuts are 1.0, 0.80, 0.60 and 0.33, respectively. We have adopted a neutrino spectrum  $\propto E^{-2}$ , saturating the Waxman-Bahcall flux bound, see Eq. (3).

The rate of neutrino-induced showers expected to be observed in an experiment such as Auger can be written as

$$\begin{aligned} \frac{N_{\text{events}}}{\Delta T_{\text{obs}}} &= 2\pi N_A \int dE_\nu \int_0^1 dy \frac{d\sigma_{\nu N}}{dy}(E_\nu) \int d\cos\theta_z A_\perp(\cos\theta_z) \\ &\times \int_{X_{\text{min}}}^{X_{\text{ground}}} dX P[E_{\text{sh}}, \cos\theta_z, X] \frac{dN_\nu}{dE_\nu}(E_\nu), \end{aligned} \quad (1)$$

where  $\Delta T_{\text{obs}}$  is the observation time,  $N_A$  is Avogadro's number,  $d\sigma_{\nu N}/dy$  is the differential neutrino-nucleon cross-section,  $y$  is the inelasticity,  $\theta_z$  is the zenith angle,  $A_\perp$  is

the cross sectional area of the experiment as seen from a given zenith angle,  $X$  is the atmospheric depth of the interaction (the atmospheric mass per unit area),  $E_{\text{sh}}$  is the total energy dissipated in the shower, and  $dN_\nu/dE_\nu$  is the incoming neutrino flux. The function  $P[E_{\text{sh}}, \theta_z, X]$  is the probability of the experiment detecting a shower created at an atmospheric depth  $X$  of energy  $E_{\text{sh}}$ , at a zenith angle  $\theta_z$ . In order to ensure that such a shower can be distinguished from one initiated by a hadron or photon primary, we require that for the Auger ground array:

- The zenith angle,  $\theta_z > 70^\circ$ ,
- The neutrino penetrates at least  $2000 \text{ g/cm}^2$  into the atmosphere before interacting,
- The interaction takes place within  $2000 \text{ g/cm}^2$  of the detector.

This third requirement is included due to the difficulty in reconstructing events beyond this range at Auger. At  $\theta_z = 70^\circ$ , the total path length traversed before the shower hits the Earth's surface is  $X_{\text{ground}} \approx 3000 \text{ g/cm}^2$ , hence relatively little atmosphere is present in which a neutrino primary can interact and be distinguished from an ordinary cosmic ray shower. At zenith angles larger than  $85^\circ$  however, the slant depth exceeds  $10000 \text{ g/cm}^2$ . In Fig. 1, we show how these selection criteria will affect the observed angular distribution of quasi-horizontal showers in Auger.

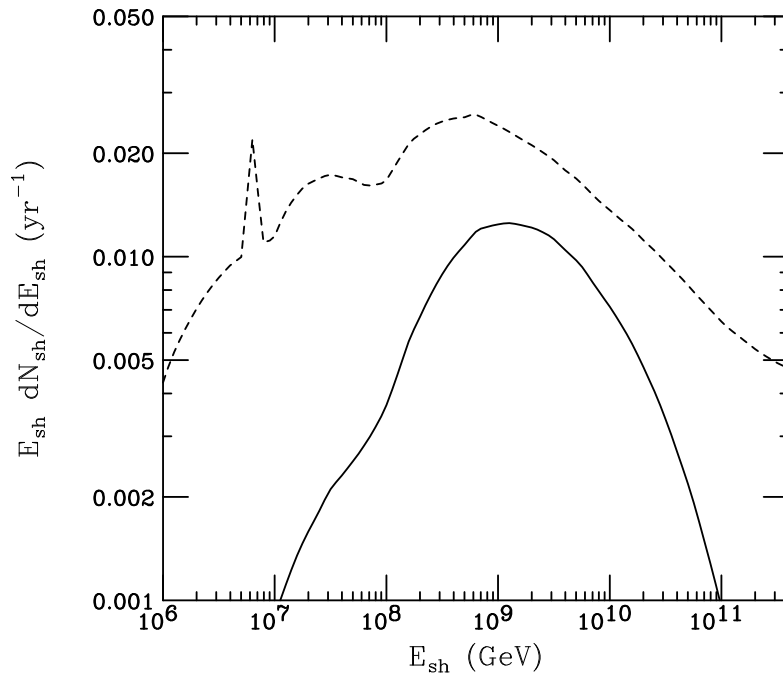


FIG. 2: Spectra of quasi-horizontal, deeply penetrating, neutrino induced showers as would be seen by Auger with the cuts shown in Fig.1. The solid and dashed lines correspond to the cosmogenic flux and the Waxman-Bahcall flux (see section III for details) which yield, respectively, 0.07 and 0.22 events per year. For the latter case, the ‘bumps’ at  $\sim 10^{7.5} \text{ GeV}$  and  $\sim 10^{8.5} \text{ GeV}$  correspond to NC and CC interactions respectively, while the  $W^-$  resonance can be seen at  $6.3 \times 10^6 \text{ GeV}$ .

If the neutrino-nucleon cross-section is enhanced well above the SM prediction at ultra-high energies, it is possible that this depth of atmosphere will significantly attenuate the cosmic neutrino flux; to account for this, an additional factor of  $e^{-N_A (X_{\text{ground}} - 2000 \text{ g/cm}^2) \sigma_{\nu N}}$

should be included [6] in Eq. (1). For SM interactions, this factor is very nearly unity and can safely be neglected but it will be important for some of the exotic models we will consider.

The neutrino-nucleon cross-section in Eq. (1) describes both charged current (CC) neutrino-quark scattering and neutral current (NC) neutrino-quark scattering for which we adopt the cross-sections given in Ref. [7]. The energy of the shower produced depends on the neutrino flavor and the type of interaction [8]. Electron neutrinos undergoing CC interactions produce a shower with both an electromagnetic and hadronic component:  $E_{\text{sh,em}} = (1 - y)E_\nu$ ,  $E_{\text{sh,had}} = yE_\nu$ . Muon neutrinos undergoing CC interactions, as well as all neutrino flavors undergoing NC interactions, produce a hadronic shower with an energy,  $E_{\text{sh,had}} = yE_\nu$ .

Charged current interactions of tau neutrinos are somewhat more complicated but more interesting. The tau lepton produced in the initial CC neutrino interaction has a decay length of  $L_\tau \approx 50 \text{ m} \times (E_\tau/10^6 \text{ GeV})$ . Thus, at sufficiently high energies, the second hadronic shower from the tau decay will be spatially separated and be identifiable as a “double bang” event [9]. Lacking a full-blown simulation, we estimate that a separation of the two bangs by 10 km would be adequate for definitive identification by the Auger ground array — this requires that the primary neutrino energy exceed  $\sim 3 \times 10^9 \text{ GeV}$  and that the first interaction occurs 50 km or more above the ground (which is easily satisfied for neutrino induced showers inclined over  $80^\circ$ ). However, if the energy exceeds  $\sim 10^{10} \text{ GeV}$ , the tau lepton will hit the ground before decaying. A more careful analysis of shower profiles as seen by the Auger fluorescence detectors may allow significant acceptance for such events over a somewhat broader energy range.

The scattering of electron flavor anti-neutrinos with electrons can occur efficiently via the resonant exchange of a  $W^-$  boson [10] at a neutrino energy of  $6.3 \times 10^6 \text{ GeV}$ . Although we include this process in our calculations, the detector acceptance for showers at this energy is expected to be rather low, thus this process is of only marginal importance.

To accurately determine the probability  $P[E_{\text{sh}}, \cos \theta_z, X]$  of the Auger ground array observing a shower with a given energy, zenith angle and initiated at a given depth, a detailed detector simulation is required, which is beyond the scope of this study. To make a reasonable estimate we have modelled the energy dependence such that we reproduce the acceptances found through the simulations performed in Ref. [3]. The probability function we arrive at is of order unity for shower energies of  $\mathcal{O}(10^{12}) \text{ GeV}$ , decreases slowly down to energies of  $\mathcal{O}(10^9) \text{ GeV}$ , and then falls rapidly at lower energies. We treat hadronic and electromagnetic showers separately as in Ref. [3]; in the case of a mixed electromagnetic-hadronic shower, we treat it as two separate showers for the purpose of estimating the probability of detection. In Fig. 2 we plot the spectrum of quasi-horizontal, deeply penetrating, neutrino induced showers as would be seen by the Auger ground array, for two choices of the ultra-high energy cosmic neutrino spectrum.

### C. Earth Skimming Tau Neutrinos

A second class of neutrino events potentially observable at Auger is generated by tau neutrinos which interact while skimming the Earth’s surface [11]. Such interactions can generate tau leptons which escape the Earth and produce a slightly upgoing hadronic

shower when they decay in the atmosphere.<sup>2</sup> This does not happen for electron neutrinos since the electrons produced in CC interactions are invariably absorbed in the Earth. For muon neutrinos, the produced muon can escape the Earth, but will not decay in the atmosphere since the decay length is  $\gtrsim 10^8$  times longer than for a tau.

When a tau lepton is generated in the Earth, it loses energy via electromagnetic processes at a rate per unit length of [11]

$$\frac{dE_\tau}{dx} \approx -\alpha - \beta E_\tau, \quad (2)$$

where  $\alpha = 0.002 \text{ GeV cm}^2/\text{g}$  and  $\beta = 6 \times 10^{-7} \text{ cm}^2/\text{g}$ . At very high energies, this will often dramatically reduce the energy before the tau is able to decay. At moderate energies, this has little effect over a single decay length. In Fig. 3 we show the effect of interactions for tau neutrino ‘Earth skimmers’. For incoming zenith angles only slightly below the horizon, the spectrum is not suppressed until above  $\sim 10^9 \text{ GeV}$ , while there is a noticeable pile-up near  $10^7 \text{ GeV}$ .

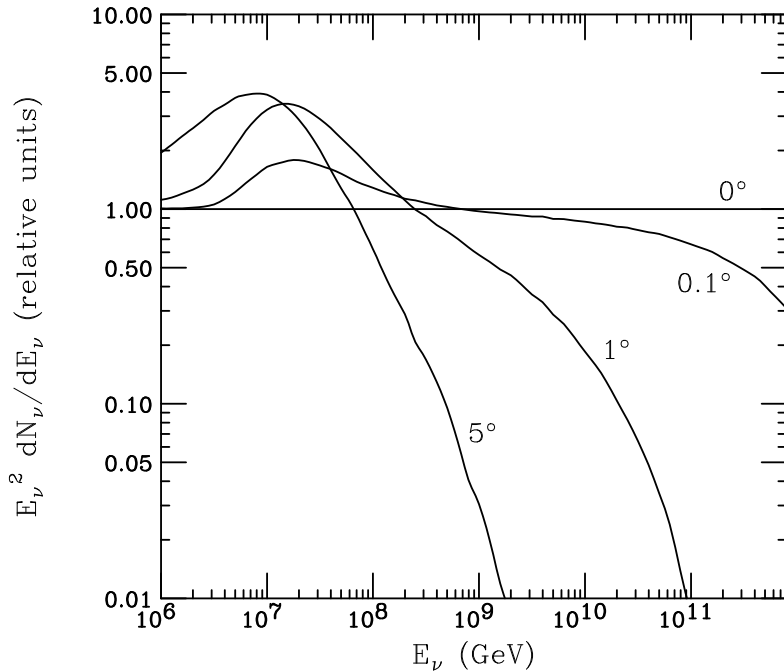


FIG. 3: Effect of interactions in the Earth on cosmic tau neutrinos with a spectrum  $\propto E_\nu^{-2}$  extending up to  $10^{12} \text{ GeV}$ . The horizontal line is the unmodified spectrum and the other lines are for incoming angles  $0.1^\circ$ ,  $1^\circ$  and  $5^\circ$  degrees below the horizon.

Tau leptons produced in CC interactions near the Earth’s surface can occasionally escape the Earth before decaying, and thus produce a hadronic shower which is potentially observable by Auger. We have calculated the spectrum of tau leptons escaping the Earth’s surface by Monte Carlo using the energy loss rate of Eq. (2). To calculate the probability of given tau neutrino induced shower being detected by Auger, we have used the same probabilities as employed in the case of quasi-horizontal showers. In addition to this function, however, we require that the shower be initiated at a height such that the

<sup>2</sup> Even when the tau lepton decays in the Earth, regeneration effects [12] can extend the effective range, such that another tau lepton emerges and decay in the atmosphere.

shower is still able to be detected (for the details of this aspect of the calculation, see Ref. [13]). This is particularly important for very high energy tau leptons which can escape the Earth’s atmosphere before decaying [11]. The Andes mountains near Auger’s southern site are also a possible target for tau neutrinos; however, we have not included this effect in our calculations as the overall correction is less than 10% [14]. In Fig. 4 we show the spectrum of Earth skimmers as would be seen by Auger.

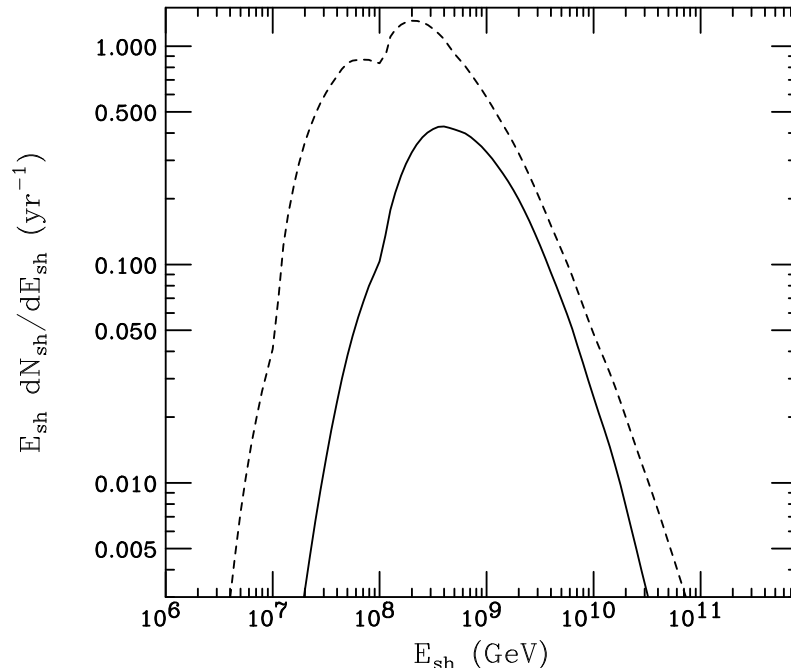


FIG. 4: The spectrum of Earth skimming, tau neutrino induced showers as would be seen by Auger. The solid and dashed lines are, respectively, for the cosmogenic neutrino flux and the Waxman-Bahcall flux, which would yield 1.3 and 4.8 events per year in Auger.

### III. ULTRA-HIGH ENERGY COSMIC NEUTRINO FLUXES

Ultra-high energy neutrinos may be produced in a wide range of astrophysical sources. In this section, we briefly discuss some of these possibilities.

Interactions of ultra-high energy cosmic ray protons propagating over cosmological distances with the cosmic microwave background generates a cosmogenic flux of neutrinos [15] through the decay of charged pions produced in  $p\gamma$  interactions [16], which should also result in a suppression of the cosmic ray spectrum above the ‘GZK cutoff’:  $E_{GZK} \sim 5 \times 10^{10}$  GeV. The intermediate state of the reaction  $p\gamma_{CMB} \rightarrow N\pi$  is dominated by the  $\Delta^+$  resonance, because the  $n$  decay length is smaller than the nucleon mean free path on the relic photons. Hence, there is roughly an equal number of  $\pi^+$  and  $\pi^0$ . Gamma rays, produced via  $\pi^0$  decay, subsequently cascade electromagnetically on the cosmic radiation fields through  $e^+e^-$  pair production followed by inverse Compton scattering. The net result is a pile up of  $\gamma$  rays at GeV energies, just below the threshold for further pair production. On the other hand, each  $\pi^+$  decays to 3 neutrinos and a positron. The  $e^+$  readily loses its energy through synchrotron radiation in the cosmic magnetic fields. The neutrinos carry away about 3/4 of the  $\pi^+$  energy, and therefore the energy in cosmogenic neutrinos is about 3/4 of the one produced in  $\gamma$ -rays.

The normalisation of the neutrino flux depends critically on the cosmological evolution of the cosmic ray sources and on their proton injection spectra [17, 18]. It also depends on the assumed spatial distribution of sources; for example, relatively local objects, such as sources in the Virgo cluster [19], would dominate the high energy tail of the neutrino spectrum. Another source of uncertainty in the cosmogenic neutrino flux is the energy at which there is a transition from galactic to extragalactic cosmic rays as inferred from a change in the spectral slope. While Fly’s Eye data [20] seem to favour a transition at  $10^{10}$  GeV, a recent analysis of the HiRes data [21] points to a lower value of  $\sim 10^9$  GeV. This translates into rather different proton luminosities at the sources [22] and consequently different predictions for the expected flux of neutrinos [23]. A fourth source of uncertainty in the cosmogenic flux is the chemical composition — if ultra-high energy cosmic rays are heavy nuclei rather than protons the corresponding cosmogenic neutrino flux may be somewhat reduced [24]. Throughout this paper, we will adopt the cosmogenic neutrino spectrum as calculated in Ref. [18].

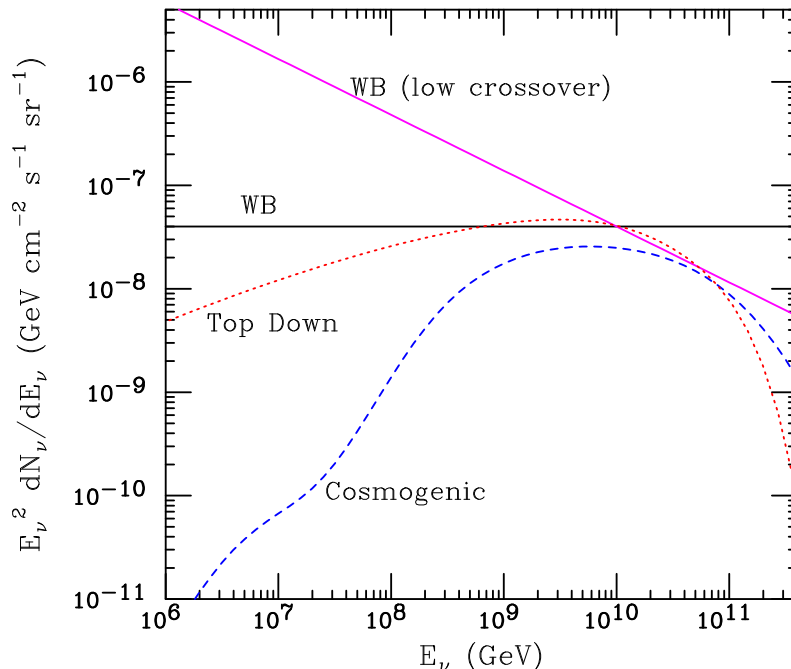


FIG. 5: Different possibilities for the ultra-high energy cosmic neutrino spectrum. The solid horizontal line (WB) corresponds to the Waxman and Bahcall bound [26] assuming proton-proton interactions. The more rapidly falling solid line (WB (low crossover)) is obtained by the same argument but assuming a lower energy transition between galactic and extragalactic cosmic rays [22]. The dotted line (Top Down) is the predicted flux in models where the highest energy cosmic rays arise from the decays of superheavy dark matter particles to many body states [32]. Finally, the dashed line (Cosmogenic) is the spectrum of neutrinos produced in the intergalactic propagation of ultra-high energy protons [18]. In all cases, the curves show the sum of neutrinos and anti-neutrinos of all flavors.

In addition to being produced in the propagation of ultra-high energy cosmic rays, neutrinos are also expected to be generated in their sources, such as gamma-ray bursts or active galactic nuclei [25]. Although the details of the relationship between the cosmic ray spectrum and the cosmic neutrino spectrum are model dependent, some rather general arguments can be applied. In particular, Waxman and Bahcall [26] have shown that for compact, cosmological sources, which are optically thin to proton-proton and proton-



photon interactions, an upper limit can be placed on the flux of neutrinos. We will follow them in adopting a neutrino spectrum arising from proton-proton collisions (with an inelasticity of 60%):

$$E_\nu^2 dN_\nu/dE_\nu \lesssim 4 \times 10^{-8} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}, \quad (3)$$

summed over all flavors. After oscillations during propagation, one finds at Earth a nearly identical flux of the three neutrino flavors [9] with equal number of neutrinos and antineutrinos [27]. If the shape of the neutrino spectrum is not an  $E^{-2}$  power law, or if the other assumptions of the Waxman-Bahcall argument are modified, this bound can be exceeded [28]. For example, if their bound is evaluated under the assumption of a low galactic to extragalactic crossover energy ( $\sim 4 \times 10^8$  GeV rather than the  $\sim 10^{10}$  GeV used by Waxman and Bahcall) a larger flux with a steeper spectrum ( $E^{-2.54}$ ) is obtained [22]. Furthermore, sources which are optically thick such that only neutrinos can escape ('hidden sources'), can easily exceed this bound [29]).

Finally, if the highest energy cosmic rays are not accelerated in distant astrophysical sources but are instead produced relatively locally in the galactic halo in the decays of supermassive dark matter particles, then significantly higher fluxes of ultra-high energy photons and neutrinos will also be generated [30]. This model was motivated by the AGASA observation that the cosmic ray spectrum continues apparently without attenuation beyond  $E_{\text{GZK}}$  but at the same time the events are isotropically distributed on the sky. Such events have not yet been seen by Auger, which has moreover set a restrictive upper limit on the fraction of photons in the cosmic ray flux [31]. We have normalized the theoretical expectations for the neutrino flux from QCD and electroweak fragmentation in heavy particle decay as calculated in Ref. [32] by matching the flux of nucleons observed by Auger, and checked that the photon limits are not violated.

In Fig. 5, we plot the expected spectrum of ultra-high energy cosmic neutrinos in the models discussed above and give the corresponding event rates for Auger with standard QCD parton model calculations in Table I. In this calculation we have truncated the cosmic neutrino spectra above  $10^{12}$  GeV — this choice has only a mild effect on the estimated rates.

	Quasi-horizontal	Earth-skimming $\nu_\tau$
Cosmogenic	0.067	1.3
Waxman-Bahcall	0.22	4.8
Waxman-Bahcall (low crossover)	2.1	35
Top-Down	0.16	4.1

TABLE I: The number of neutrino induced events per year expected in Auger for various choices of the ultra-high energy neutrino spectrum, as shown in Figure 5, calculated using the standard QCD parton model cross-section.

## IV. NEUTRINO PHYSICS WITH AUGER

### A. Prospects for Cross-Section Measurements

Deviations of the neutrino-nucleon cross-sections from the prediction of the simple parton model [7] can signal new physics beyond the SM, but might alternatively be just

due to saturation effects which can substantially modify the parton density at small  $x$  (i.e. small energy fractions) [33]. These effects can significantly reduce the total cross-section at high energies, softening the power law behavior predicted by the simple ‘unscreened’ parton model toward compliance with the Froissart bound [34]. By contrast, new physics such as TeV-scale quantum gravity [35, 36] can enhance the neutrino interaction cross-sections.<sup>3</sup> This has been calculated in various different frameworks, e.g., arising from exchange of Kaluza-Klein (KK) gravitons [38, 39], black hole production [40], and TeV-scale string excitations [41]. In this section, we will discuss the ability of Auger to measure deviations in the neutrino-nucleon scattering cross-section from the SM prediction, without assuming any particular interaction model.

The event rates for quasi-horizontal and Earth-skimming neutrinos have different responses to the inelastic cross-section [42]. The rate of quasi-horizontal showers *rises* proportional to the cross-section (although if this exceeds  $\sim 10^{-28}$  cm<sup>2</sup>, attenuation of the neutrino flux in the upper atmosphere becomes significant [6]). By contrast, the rate of Earth skimming tau events is always *depleted* by an enhanced neutrino-nucleon cross-section because of absorption in the Earth.

In order to probe deviations from the (unscreened) parton model calculation of the cross-section, it is necessary to note that the screening corrections affect CC and NC equally. To assess the experimental sensitivity to such effects, we assume a uniform suppression of the cross-section by a factor of 2 or 5 and show in Fig. 6 the effect on the spectrum of Earth skimmers, as a function of the incoming angle. For a cross section reduced by a factor of 2, the total event rate of Earth skimmers is  $1.4 \text{ yr}^{-1}$ , which is slightly *larger* than for the unscreened parton model. The reduction in cross-section due to screening will be energy dependent in general, but as shown in Table II, the effect is mainly manifest at intermediate energies of  $\sim 10^8 - 10^{10}$  GeV, corresponding to center-of-mass energies  $\sqrt{s} \simeq 10^4 - 10^5$  GeV; at these energies the ratio of quasi-horizontal to Earth-skimming events is a useful diagnostic of any suppression in the cross-section. This is in fact primarily because the cosmogenic neutrino flux peaks at these energies, nevertheless since this represents a reasonable lower limit to the expected flux, this sensitivity is likely to be achieved and even surpassed. A factor of 2 reduction in the cross-section may appear extreme, even so it is clear that Auger can probe the behavior of parton distribution functions (pdfs) in a kinematic region out of reach of foreseeable accelerators. This will be particularly beneficial for calibrating different hadronic interaction models of air shower development, which presently differ significantly in their predictions [43].

With regard to enhancements of the cross-section by new physics, in general this will be different for CC and for NC interactions. To assess the sensitivity of Auger, we consider a toy model in which only the NC cross-section is enhanced by a factor ranging between 3 and 100, while assuming the inelasticity to be the same as in the SM.<sup>4</sup> In Fig. 7 we show that this results in a suppression of Earth-skimming tau spectrum. By contrast the quasi-horizontal showers are enhanced, resulting in a steady increase of the ratio of quasi-horizontals to Earth skimmers, as the NC cross-section is increased (see Table III). Clearly Auger would be sensitive to substantial increases of the NC cross-section.

Thus both an increase and a decrease of the neutrino-nucleon cross-section from the naïve SM value will have distinctive observational signatures. To quantitatively assess the sensitivity of Auger to such effects, the uncertainty in the cosmic neutrino fluxes must

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<sup>3</sup> It is noteworthy that the neutrino-nucleon cross section can also be enhanced in some supersymmetric models through direct channel production of superpartner resonances [37].

<sup>4</sup> This resembles the KK graviton exchange, as we discuss in the next section.

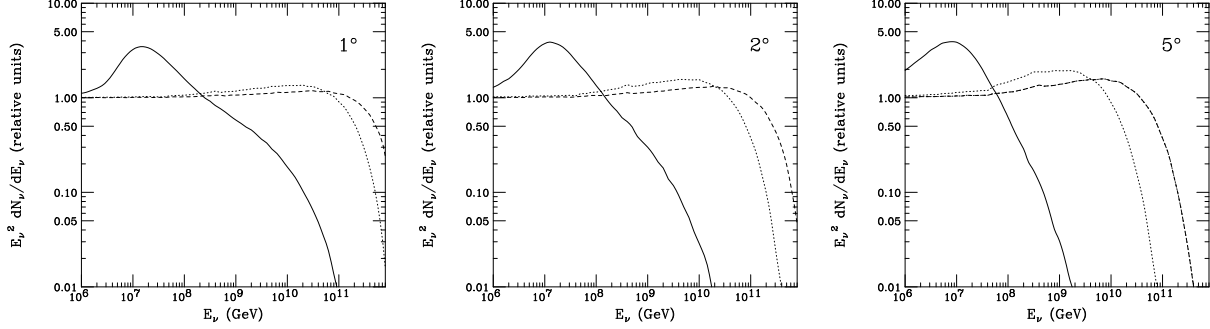


FIG. 6: The effect of interactions in the Earth on the tau neutrino spectrum when the (CC + NC) interaction cross-section is *suppressed*. As in Fig. 3, we adopt a spectrum  $\propto E_\nu^{-2}$ , which extends to  $10^{12}$  GeV. In each frame, results are shown assuming the SM cross-section (solid-line) and a cross-section smaller by a factor of 2 (dotted-line) and a factor of 5 (dashed-line). The three frames are for incoming angles of  $1^\circ$ ,  $2^\circ$  and  $5^\circ$  degrees below the horizon.

also be taken into account [44]. Moreover, to determine the acceptances to different types of events, a full detector simulation is clearly required to improve over the approximate estimates [3] adopted here.

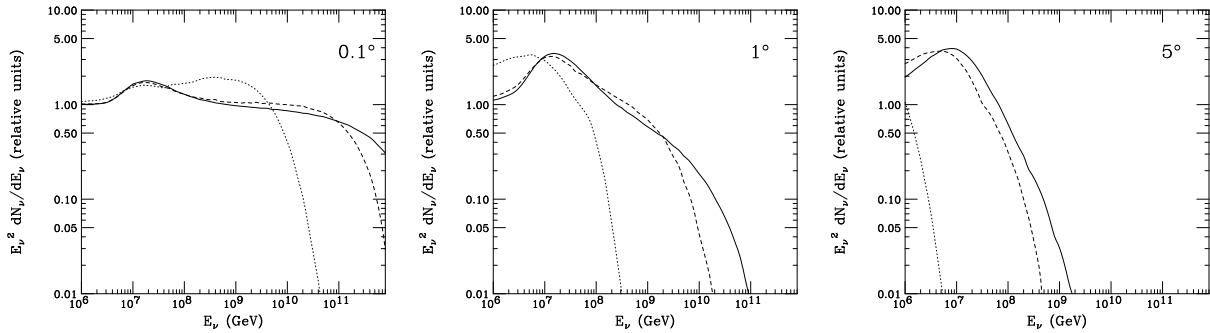


FIG. 7: The effect of interactions in the Earth on the tau neutrino spectrum when the (NC) interaction cross-section is *enhanced*. As in Fig. 3, we adopt a spectrum  $\propto E_\nu^{-2}$ , which extends to  $10^{12}$  GeV. In each frame, results are shown for the cases of the SM cross-section (solid-line), a cross-section 10 times larger (dashed-line) and 100 times larger (dotted-line). The three frames are for incoming angles of  $0.1^\circ$ ,  $1^\circ$  and  $5^\circ$  degrees below the horizon.

$\sigma_{\nu N}$	$10^6 - 10^7$	$10^7 - 10^8$	$10^8 - 10^9$	$10^9 - 10^{10}$	$10^{10} - 10^{11}$	$10^{11} - 10^{12}$
SM	$3.6 \times 10^{-5}$	0.056	0.85	0.41	0.020	$1.1 \times 10^{-4}$
$\text{SM} \times \frac{1}{2}$	$2.1 \times 10^{-5}$	0.057	0.86	0.45	0.026	$1.8 \times 10^{-4}$

TABLE II: Variation of the rate (in  $\text{yr}^{-1}$ ) of Earth-skimming tau neutrino induced events in various energy intervals (in GeV), for the SM (unscreened parton) cross-section, and for a cross section 2 times smaller (for *both* CC and NC). The cosmogenic neutrino flux has been assumed.

$\sigma_{\nu N}$	Quasi-horizontal	Earth-skimming $\nu_\tau$	Ratio
Standard Model	0.067	1.3	0.05
SM $\times 3$	0.096	1.1	0.09
SM $\times 10$	0.20	0.68	0.29
SM $\times 100$	1.5	0.081	19

TABLE III: The energy integrated rate (in  $\text{yr}^{-1}$ ) of quasi-horizontal and Earth-skimmers, as well as their ratio, for the SM cross-section and for different enhancements of the NC component alone. The cosmogenic neutrino flux has been assumed.

## B. Prospects for Flavor Ratio Measurements

In most models of astrophysical neutrino sources, neutrinos are generated through the decay of charged pions:  $\pi^+ \rightarrow \mu^+ \nu_\mu \rightarrow e^+ \nu_e \bar{\nu}_\mu \nu_\mu$  or  $\pi^- \rightarrow \mu^- \bar{\nu}_\mu \rightarrow e^- \bar{\nu}_e \nu_\mu \bar{\nu}_\mu$ , thus the flavor ratio at source is  $\nu_e : \nu_\mu : \nu_\tau = 1/3 : 2/3 : 0$ . However, oscillations modify this ratio as neutrinos propagate to Earth. Given the observed near maximal mixings [45] and the long baselines involved, the predicted flavor ratio at Earth is  $\nu_e : \nu_\mu : \nu_\tau \approx 0.36 : 0.33 : 0.31$  following Ref. [9]. However, cosmic (anti-)neutrinos may also be generated in the decay of neutrons:  $n \rightarrow p^+ e^- \bar{\nu}_e$ . In this case, the initial flavor ratio of  $\nu_e : \nu_\mu : \nu_\tau = 1 : 0 : 0$  becomes  $\nu_e : \nu_\mu : \nu_\tau \approx 0.56 : 0.26 : 0.18$  at Earth [46]. In either case a measured deviation from these predictions could indicate new physics if the neutrino production mechanism is well understood.

To study the sensitivity of Auger to the flavor content, we plot in Fig. 8 the ratio of quasi-horizontal showers to Earth-skimming events as the  $\nu_e$  flux is varied in ratio to the other flavors. As in Table III, this ratio is 0.05 when the flux is equally spread among flavors.

## V. MODELS OF NEW PHYSICS

### A. Low Scale Quantum Gravity

Two of the most important scales in physics are the Planck scale ( $M_{\text{Pl}} = G_{\text{N}}^{-1/2} \simeq 10^{19}$  GeV) and the weak scale ( $M_W = G_{\text{F}}^{-1/2} \simeq 300$  GeV), and a long standing problem is explaining the hierarchy between these scales. The traditional view is to adopt  $M_{\text{Pl}}$  as *the* fundamental scale and attempt to derive  $M_W$  through some dynamical mechanism (e.g. renormalization group evolution). However recently several models [35, 36] have been proposed where  $M_W$  is instead the fundamental scale of nature. In the simplest construction of these models, the SM fields are confined to a 3+1-dimensional ‘brane-world’ (corresponding to our observed universe), while gravity propagates in a higher dimensional ‘bulk’ space.

If space-time is assumed to be a direct product of a 3+1-dimensional manifold and a flat spatial  $n$ -dimensional torus  $T^n$  (of common linear size  $2\pi r_c$ ), one obtains a definite representation of this picture in which the effective 4-dimensional Planck scale is related to the fundamental scale of gravity,  $M_D$ , according to [35]

$$M_{\text{Pl}}^2 = 8\pi r_c^n M_D^{n+2}, \quad (4)$$

where  $D = 4 + n$ . If  $M_D$  is to be not much higher than the electroweak scale, then this requires  $r_c$  to be large in Planck units and thus reformulates the hierarchy problem.

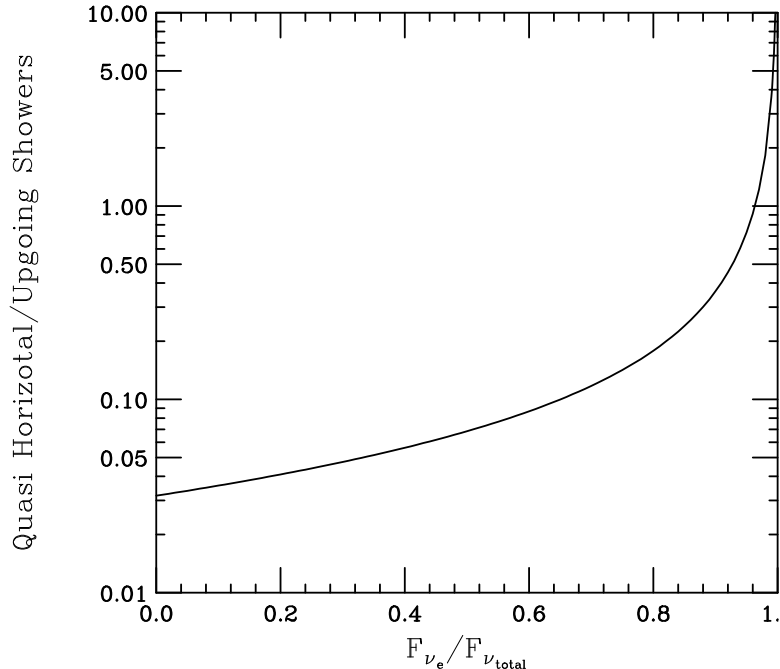


FIG. 8: The ratio of neutrino induced quasi-horizontal showers to Earth-skimming tau neutrino induced upgoing showers as a function of the flavor content of the cosmic neutrino flux. We have assumed equal numbers of muon and tau neutrinos and adopted the cosmogenic neutrino flux.

For illustrative purposes in what follows we will consider only the case of flat extra-dimensions. The consequences of more exoteric scenarios (such as warped extra-dimensions [36]) have been studied in detail by various authors [47].

### 1. Sub-Planckian Regime

From our 4-dimensional point of view, the higher dimensional massless gravitons then appear as an infinite tower of KK modes, of which the lowest is the massless graviton itself, while its excitations are massive. The mass-squared of each KK graviton mode reads,  $m^2 = \sum_{i=1}^n \ell_i^2/r_c^2$ , where the mode numbers  $\ell_i$  are integers. Note that the weakness of the gravitational interaction is compensated by the large number of KK modes that are exchanged: the coupling  $M_{\text{Pl}}^{-2}$  of the graviton vertex is cancelled exactly by the large multiplicity of KK excitations  $\sim \hat{s}^{n/2} r_c^n$ , so that the final product is  $\sim \hat{s}^{n/2}/M_D^{2+n}$  [48]. Here  $\sqrt{\hat{s}}$  is the center-of-mass energy available for graviton-KK emission. Taking brane fluctuations into account, a form factor  $\sim e^{-m^2/M_D^2}$  is introduced at each graviton vertex [49]. This exponential suppression, which parametrizes the effects of a finite brane tension, provides a dynamical cutoff in the (otherwise divergent) sum over all KK contributions to a given scattering amplitude. Altogether, one may wonder whether the rapid growth of the cross-section with energy in neutrino-nucleon reactions mediated by spin 2 particles carries with it observable deviations from SM predictions.

A simple Born approximation to the elastic neutrino-parton cross section (which underlies the total neutrino-proton cross-section) leads, without modification, to  $\hat{\sigma}_{\text{el}} \sim \hat{s}^2$  [38, 39]. Unmodified, this behavior by itself eventually violates unitarity. This may be seen either by examining the partial waves of this amplitude, or by studying the high

energy Regge behavior of an amplitude  $A_R(\hat{s}, \hat{t}) \propto \hat{s}^{\alpha(\hat{t})}$  with spin-2 Regge pole, *viz.*, intercept  $\alpha(0) = 2$ . For the latter, the elastic cross-section is given by

$$\frac{d\hat{\sigma}_{\text{el}}}{d\hat{t}} \sim \frac{|A_R(\hat{s}, \hat{t})|^2}{\hat{s}^2} \sim \hat{s}^{2\alpha(0)-2} \sim \hat{s}^2, \quad (5)$$

whereas the total cross-section reads

$$\hat{\sigma}_{\text{tot}} \sim \frac{\Im[A_R(\hat{s}, 0)]}{\hat{s}} \sim \hat{s}^{\alpha(0)-1} \sim \hat{s}, \quad (6)$$

so that eventually,  $\hat{\sigma}_{\text{el}} > \hat{\sigma}_{\text{tot}}$  [50]. Eikonal unitarization schemes modify this behaviour. Specifically, for large impact parameter, a single Regge pole exchange amplitude yields  $\hat{\sigma}_{\text{tot}} \sim \ln^2(\hat{s}/s_0)$  [51], an estimate which is insensitive to the underlying theory at short distances (UV completion). Recently, the differential cross-section for such gravity-mediated interaction at large distances has been calculated [52]. Because of the large cross-sections, albeit with low inelasticity, there would be distinctive double and/or multiple bang events [53] similar to those discussed in Sec. II B [54]. For small impact parameters, it becomes difficult to respect partial wave unitarity as corrections to the eikonal amplitude are expected to become important. Note that graviton self interactions carry factors of  $\hat{t}$  associated to the vertices, and thus as  $\hat{t}$  increases, so does the attraction among the scattered particles. Eventually it is expected that gravitational collapse to a black hole (BH) will take place, absorbing the initial state in such a way that short distance effects are screened by the appearance of a horizon [40, 55].

## 2. Trans-Planckian Regime

According to Thorne’s hoop conjecture [56], a BH forms in a two-particle collision when and only when the impact parameter is smaller than the radius of a Schwarzschild BH of mass equal to  $\sqrt{\hat{s}} \equiv \sqrt{xs}$ . The total cross-section for BH production is then,

$$\hat{\sigma}_{\text{BH}} = F(n) \pi r_s^2(\sqrt{\hat{s}}), \quad (7)$$

proportional to the area subtended by a “hoop” of radius [57]

$$r_s(\sqrt{\hat{s}}) = \frac{1}{M_D} \left[ \frac{\sqrt{\hat{s}}}{M_D} \right]^{\frac{1}{1+n}} \left[ \frac{2^n \pi^{\frac{n-3}{2}} \Gamma(\frac{n+3}{2})}{n+2} \right]^{\frac{1}{1+n}}, \quad (8)$$

where  $F(n)$  is a form factor of order unity. Recent work has confirmed the validity of Eq. (7) and evaluated the dimension-dependent constant  $F(n)$ , analytically in four dimensions [58] and numerically in higher dimensions [59]. In the course of collapse, a certain amount of energy is radiated in gravitational waves by the multipole moments of the incoming shock waves [60], leaving a fraction  $y \equiv M_{\text{BH}}/\sqrt{\hat{s}}$  available to be emitted through Hawking evaporation [61]. Here,  $M_{\text{BH}}$  is a *lower bound* on the final mass of the BH and  $\sqrt{\hat{s}}$  is the center-of-mass energy of the colliding particles, taken to be partons. This ratio depends on the impact parameter of the collision, as well as on the dimensionality of space-time [62]. Of course, this calculation is purely in the framework of classical general relativity, and is expected to be valid only for energies far above the fundamental Planck scale  $M_D$ , for which curvature is small outside the horizon and strong quantum effects are hidden behind the horizon. Extending this formalism to center-of-mass energies close to  $M_D$  requires a better understanding of quantum gravity.

String theory provides the best hope for understanding the regime of strong quantum gravity, and in particular for computing cross-sections at energies close to the Planck scale [63]. In principle embedding TeV-scale gravity models in realistic string models might facilitate the calculation of cross-sections for BHs (and string excitations) having masses comparable to  $M_D$ . To be specific we will consider embedding of a 10-dimensional low-energy scale gravity scenario within the context of SO(32) Type I superstring theory, where gauge and charged SM fields can be identified with open strings localized on a 3-brane and the gravitational sector consists of closed strings that propagate freely in the internal dimensions of the universe [64]. After compactification on  $T^6$  down to four dimensions,  $M_{\text{Pl}}$  is related to the string scale,  $M_s$ , and the string coupling constant,  $g_s$ , by  $M_{\text{Pl}}^2 = (2\pi r_c)^6 M_s^8 / g_s^2$  (hereafter,  $D = 10$ , i.e.  $n = 6$ ).

Subsequent to formation, the BH proceeds to decay [65]. The decay of an excited spinning BH state proceeds through several stages. The initial configuration loses hair associated with multipole moments in a balding phase by emission of classical gravitational and gauge radiation. Gauge charges inherited from the initial state partons are discharged by Schwinger emission. After this transient phase, the subsequent spinning BH evaporates by semi-classical Hawking radiation in two phases: a brief spin-down phase in which angular momentum is shed [66], and a longer Schwarzschild phase. In the latter the emission rate per degree of particle freedom  $i$  of particles of spin  $s$  with initial total energy between  $(Q, Q + dQ)$  is found to be [67]

$$\frac{d\dot{N}_i}{dQ} = \frac{\sigma_s(Q, r_s) \Omega_{d-3}}{(d-2)(2\pi)^{d-1}} Q^{d-2} \left[ \exp\left(\frac{Q}{T_{\text{BH}}}\right) - (-1)^{2s} \right]^{-1}, \quad (9)$$

where  $T_{\text{BH}} = 7/(4\pi r_s)$  is the BH temperature,

$$\Omega_{d-3} = \frac{2\pi^{(d-2)/2}}{\Gamma[(d-2)/2]} \quad (10)$$

is the volume of a unit  $(d-3)$ -sphere, and  $\sigma_s(Q, r_s)$  is the absorption coefficient (a.k.a. the greybody factor). Recall that SM fields live on a 3-brane ( $d = 4$ ), while gravitons inhabit the entire spacetime ( $d = 10$ ). The prevalent energies of the decay quanta are of  $\mathcal{O}(T_{\text{BH}}) \sim 1/r_s$ , resulting in s-wave dominance of the final state. Indeed, as the total angular momentum number of the emitted field increases,  $\sigma_s(Q, r_s)$  rapidly gets suppressed [68]. In the low energy limit,  $Q r_s \ll 1$ , higher-order terms are suppressed by a factor of  $3(Q r_s)^{-2}$  for fermions and by a factor of  $25(Q r_s)^{-2}$  for gauge bosons. For an average particle energy  $\langle Q \rangle$  of  $\mathcal{O}(r_s^{-1})$ , higher partial waves also get suppressed, although by a smaller factor. This strongly suggests that the BH is sensitive only to the radial coordinate and does not make use of the extra angular modes available in the internal space [69]. A recent numerical study [70] has explicitly shown that the emission of scalar modes into the bulk is largely suppressed with respect to the brane emission. In order to contravene the argument of Emparan–Horowitz–Myers [69], the bulk emission of gravitons would need to exhibit the opposite behavior – a substantial enhancement into bulk modes. There is no *a priori* reason to suspect this qualitative difference between  $s = 0$  and  $s = 2$ , and hence no reason to support arguments [71] favoring deviation from the dominance of visible decay. With this in mind, we assume the evaporation process to be dominated by the large number of SM brane modes. The lower bound on the mass radiated in the Schwarzschild phase could be somewhat reduced at large  $b$  [72] compared to the estimate in Ref. [62] used here. On the other hand, the effective range of  $b$  at which there is trapping is somewhat increase [72], with the result that there is not any significant change.

The total number of particles emitted is approximately equal to the BH entropy,

$$S_{\text{BH}} = \frac{\pi}{2} M_{\text{BH}} r_s. \quad (11)$$

At a given time, the rate of decrease in the BH mass is just the total power radiated

$$\frac{d\dot{M}_{\text{BH}}}{dQ} = - \sum_i c_i \frac{\sigma_s(Q, r_s)}{8\pi^2} Q^3 \left[ \exp\left(\frac{Q}{T_{\text{BH}}}\right) - (-1)^{2s} \right]^{-1}, \quad (12)$$

where  $c_i$  is the number of internal degrees of freedom of particle species  $i$ . Integration of Eq. (12) leads to

$$\dot{M}_{\text{BH}} = - \sum_i c_i f \frac{\Gamma_s}{8\pi^2} \Gamma(4) \zeta(4) T_{\text{BH}}^4 A_4, \quad (13)$$

where  $f = 1$  ( $7/8$ ) for bosons (fermions), and the greybody factor was conveniently written as a dimensionless constant,  $\Gamma_s = \sigma_s(\langle Q \rangle, r_s)/A_4$ , normalized to the BH surface area [69]

$$A_4 = \frac{36}{7} \pi \left(\frac{9}{2}\right)^{2/7} r_s^2 \quad (14)$$

seen by the SM fields ( $\Gamma_{s=1/2} \approx 0.33$  and  $\Gamma_{s=1} \approx 0.34$  [73]). Now, since the ratio of degrees of freedom for gauge bosons, quarks and leptons is 29:72:18 (excluding the Higgs boson), from Eq. (13) one obtains a rough estimate of the mean lifetime,

$$\tau_{\text{BH}} \approx 1.67 \times 10^{-27} \text{ s} \left(\frac{M_{\text{BH}}}{M_{10}}\right)^{9/7} \left(\frac{\text{TeV}}{M_{10}}\right), \quad (15)$$

which indicates that BHs evaporate near-instantaneously into visible quanta.

The semi-classical description outlined above is reliable only when the energy of the emitted particle is small compared to the BH mass, i.e.

$$T_{\text{BH}} \ll M_{\text{BH}}, \text{ or equivalently, } M_{\text{BH}} \gg M_{10}, \quad (16)$$

because it is only under this condition that both the gravitational field of the brane and the back reaction of the metric during the emission process can safely be neglected [74]. For BHs with initial masses well above  $M_{10}$ , most of the decay process can be well described within the semi-classical approximation. However, the condition stated in Eq. (16) inevitably breaks down during the last stages of evaporation. At this point it becomes necessary to introduce quantum considerations. To this end we turn to a quantum statistical description of highly excited strings.

It is well-known that the density of string states with mass between  $M$  and  $M + dM$  cannot increase any faster than  $\rho(M) = e^{\beta_{\text{H}} M}/M$ , because the partition function,

$$Z(\beta) = \int_0^\infty dM \rho(M) e^{-M\beta}, \quad (17)$$

would fail to converge [75]. Indeed, the partition function converges only if the temperature is less than the Hagedorn temperature,  $\beta_{\text{H}}^{-1}$ , which is expected to be  $\sim M_s$ . As  $\beta$  decreases to the transition point  $\beta_{\text{H}}$ , the heat capacity rises to infinity because the energy goes into the many new available modes rather than into raising the kinetic energy of the existing particles [76]. In the limit, the total probability diverges, indicating that the canonical ensemble is inadequate for the treatment of the system. However, one can still employ a microcanonical ensemble of a large number of similar isolated systems, each



with a given fixed energy  $E$ . With the center-of-mass at rest,  $E = M$  so the density of states is just  $\rho(M)$  and the entropy  $S = \ln \rho(M)$ . In this picture, equilibrium among systems is determined by the equality of the temperatures, defined for each system as

$$T \equiv \left( \frac{\partial S}{\partial M} \right)^{-1} = \frac{M}{\beta_{\text{H}} M - 1} . \quad (18)$$

Equilibrium is achieved at maximum entropy when the total system heat capacity,  $C$ , is positive. Ordinary systems (on which our intuition is founded) have  $C > 0$ . However, for a gas of massive superstring excitations the heat capacity,

$$C \equiv -\frac{1}{T^2} \left( \frac{\partial^2 S}{\partial M^2} \right)^{-1} = -\left( \frac{M}{T} \right)^2 , \quad (19)$$

is *negative*, as is the case for BHs [77]. The positivity requirement on the total specific heat implies that strings and BHs cannot coexist in thermal equilibrium, because any subsystem of this system has negative specific heat, and thus the system as a whole is thermodynamically unstable. This observation suggests that BHs may end their Hawking evaporation process by making a transition to an excited string state with higher entropy, avoiding the singular zero-mass limit [78]. The suggestion of a string  $\rightleftharpoons$  BH transition is further strengthened by three other facts: (i) in string theory, the fundamental string length should set the minimum value for the Schwarzschild radius of any BH [79]; (ii)  $T_{\text{BH}} \sim \beta_{\text{H}}^{-1}$  for  $r_s \sim M_s^{-1}$  [80]; (iii) there is an apparent correlation between the greybody factors in BH decay and the level structure of excited strings [81]. The string  $\rightleftharpoons$  BH “correspondence principle” [82] unifies these concepts: When the size of the BH horizon drops below the size of the fundamental string length  $\ell_s \gg \ell_{10}$ , where  $\ell_{10}$  is the fundamental Planck length, an adiabatic transition occurs to an excited string state. Subsequently, the string will slowly lose mass by radiating massless particles with a nearly thermal spectrum at the unchanging Hagedorn temperature [83]. (Note that the probability of a BH radiating a large string, or of a large string undergoing a fluctuation to become a BH is negligibly small [84].)

The continuity of the cross-section at the correspondence point, at least parametrically in energy and string coupling, provides an independent supporting argument for this picture [63]. Specifically, in the perturbative regime, the Virasoro-Shapiro amplitude leads to a “string ball” (SB) production cross-section  $\propto g_s^2 \hat{s} / M_s^4$ . This cross-section saturates the unitarity bounds at  $g_s^2 \hat{s} / M_s^2 \sim 1$  [85], so before matching the geometric BH cross-section  $\propto r_s^2$ , there is a transition region at which  $\hat{\sigma} \sim M_s^{-2}$ . All in all, the rise with energy of the parton-parton  $\rightarrow$  SB/BH cross-section can be parametrized as [63]

$$\hat{\sigma}(\sqrt{\hat{s}}) \sim \begin{cases} \frac{g_s^2 \hat{s}}{M_s^4} & M_s \ll \sqrt{\hat{s}} \leq M_s / g_s , \\ \frac{1}{M_s^2} & M_s / g_s < \sqrt{\hat{s}} \leq M_s / g_s^2 , \\ \frac{1}{M_{10}^2} \left[ \frac{\sqrt{\hat{s}}}{M_{10}} \right]^{2/7} & M_s / g_s^2 < \sqrt{\hat{s}} , \end{cases} \quad (20)$$

where  $M_{10} = (8\pi^5)^{1/8} M_s / g_s^{1/4}$

The inclusive production of BHs proceeds through different final states for different classical impact parameters  $b$  [62]. These final states are characterized by the fraction  $y(z)$  of the initial parton center-of-mass energy,  $\sqrt{\hat{s}} = \sqrt{xs}$ , which is trapped within the

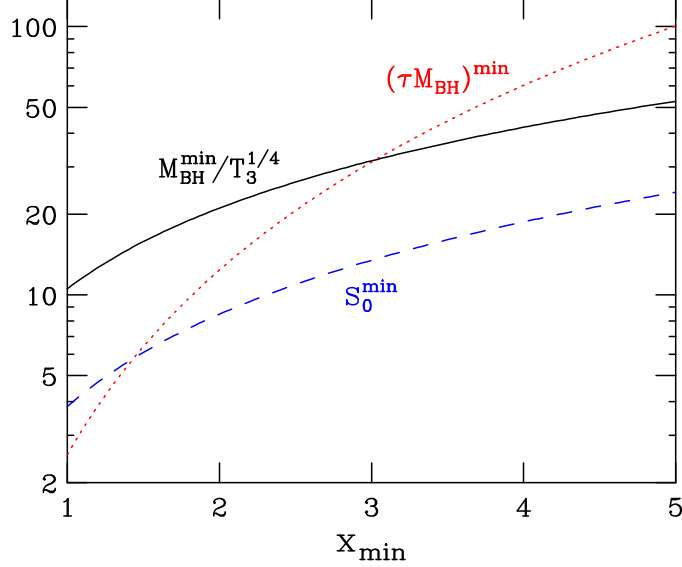


FIG. 9: Quantitative measures of the validity of the semi-classical analysis of BH production for  $n = 6$  extra dimensions, where  $x_{\min} \equiv M_{\text{BH},\min}/M_{10}$ .

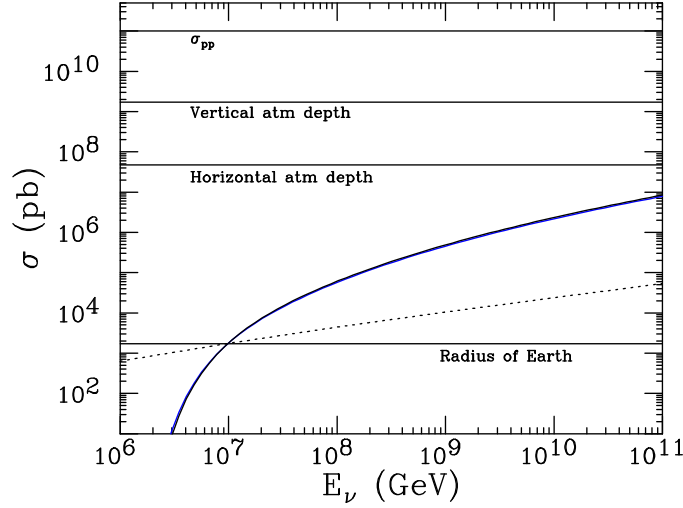


FIG. 10: The cross-section for BH production in neutrino nucleon collisions, for  $n = 6$  extra dimensions, assuming  $M_{10} = 1$  TeV and  $M_{\text{BH},\min} = M_{10}$ . Energy losses by gravitational radiation have been included. The SM  $\nu N$  cross-section is indicated by the dotted line. For comparison the typical  $pp$  cross-section is shown, as well as the cross-section required for triggering vertical and horizontal atmospheric showers. The cross-section for absorption by the Earth is also shown [91].

horizon. Here,  $z = b/b_{\max}$ , where  $b_{\max} = 1.3 r_s(\sqrt{\hat{s}})$  [62]. With a lower cutoff  $M_{\text{BH},\min}$  on the BH mass required for the validity of the semi-classical description, this implies the joint constraint

$$y(z) \sqrt{xs} \geq M_{\text{BH},\min} \quad (21)$$

on the parameters  $x$  and  $z$ . Because of the monotonically decreasing nature of  $y(z)$ , Eq. (21) sets an *upper* bound  $\bar{z}(x)$  on the impact parameter for fixed  $x$ . The corresponding parton-parton BH cross-section is  $\hat{\sigma}_{\text{BH}}(x) = \pi \bar{b}^2(x)$ , where  $\bar{b} = \bar{z} b_{\max}$ . The total BH

production cross-section is then [86]

$$\sigma_{\text{BH}}(E_\nu, M_{\text{BH,min}}, M_{10}) \equiv \int_{\frac{M_{\text{BH,min}}^2}{y^2(0)_s}}^1 dx \sum_i f_i(x, Q) \hat{\sigma}_{\text{BH}}(x), \quad (22)$$

where  $i$  labels parton species and the  $f_i(x, Q)$  are pdfs [87]. The momentum scale  $Q$  is taken as  $r_s^{-1}$ , which is a typical momentum transfer during the gravitational collapse process. The parameter  $M_{\text{BH,min}}$  plays an important role in interpreting the results derived below. The validity of the semi-classical calculation requires at least three criteria to be satisfied. First,  $S_0$ , the initial entropy of the produced BH should be large enough to ensure a well-defined thermodynamic description [74]. Second, the BH lifetime  $\tau_{\text{BH}}$  should be large compared to its inverse mass so that the black hole behaves like a well-defined resonance. Third, the BH mass must be large compared to the scale of the 3-brane tension  $T_3$  so that the brane does not significantly perturb the BH metric [88]. Quantitative measures of these three criteria are given in Fig. 9 for  $n = 6$ , assuming  $T_3 = \sqrt{8\pi}/(2\pi)^6 M_{10}^4$  for 6 toroidally-compactified dimensions [89]. We see that all three criteria are adequately satisfied for  $M_{\text{BH,min}} \gtrsim 3M_{10}$  [90]. The resulting  $\nu N \rightarrow \text{BH}$  production cross-section is shown in Fig. 10.

In the perturbative string regime, i.e.  $M_{\text{SB,min}} < \sqrt{\hat{s}} \leq M_s/g_s$ , the SB production cross-section is taken to be

$$\sigma_{\text{SB}}(E_\nu, M_{\text{SB,min}}, M_{10}) = \int_{\frac{M_{\text{SB,min}}^2}{s}}^1 dx \sum_i f_i(x, Q) \hat{\sigma}_{\text{SB}}(\hat{s}), \quad (23)$$

where  $\hat{\sigma}_{\text{SB}}(\hat{s})$  contains the Chan-Paton factors which control the projection of the initial state onto the string spectrum. In general, this projection is not uniquely determined by the low-lying particle spectrum, so there are one or more arbitrary constants. The analysis in the  $\nu q \rightarrow \nu q$  channel illustrates this point [92]. The  $\nu g$  scattering, relevant for  $\nu N$  interactions at ultra-high energies, introduces additional ambiguities. In our calculations we adopt the estimates given in Ref. [93] considering the saturation limit and including both neutrino-quark and neutrino-gluon scattering. The resulting  $\nu N \rightarrow \text{SB}$  cross-section is shown in Fig. 11, setting the Chan-Paton factors equal to 1/2.

As can be seen in Figs. 10 and 11, although the neutrino interaction length is reduced below the SM value due to BH/SB production, it is still far larger than the Earth's atmospheric depth. Neutrinos therefore would produce BH/SBs with roughly equal probability at any point in the atmosphere. As a result, the light descendants of the BH/SB may initiate low-altitude, quasi-horizontal showers at rates significantly higher than SM predictions.<sup>5</sup> Because of this the atmosphere provides a buffer against contamination by hadronic showers (for which the electromagnetic component is completely attenuated at such large zenith angles) allowing a good characterization of BH-induced showers when  $S \gg 1$  [40, 95].

If the quasi-horizontal deep shower rate is found to be anomalously large, it can be ascribed either to an enhancement of the incoming neutrino flux, or to an enhancement in the neutrino-nucleon cross-section. However, these possibilities may be distinguished by focusing on events which arrive at very small angles to the horizon. An enhanced flux will increase *both* the quasi-horizontal and Earth-skimming event rates, whereas a large BH cross-section *suppresses* the latter, because the hadronic decay products of BH

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<sup>5</sup> Additionally, neutrinos that traverse the atmosphere unscathed may produce black holes via interactions in the ice or water and be detected by neutrino telescopes [94].

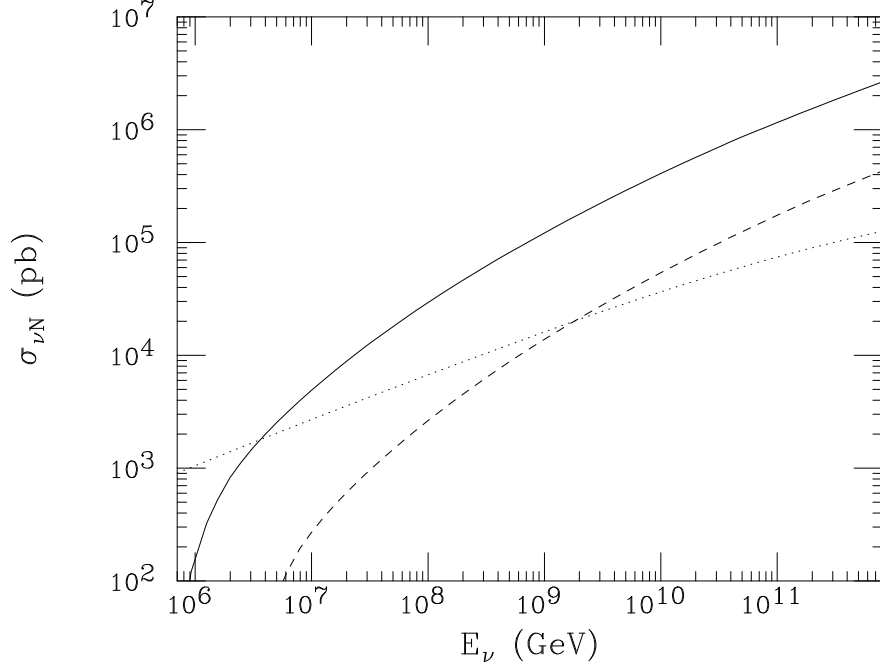


FIG. 11: The neutrino-nucleon cross-section including the effects of TeV scale string resonances. The solid and dashed lines correspond to models with string tension  $M_s = 1$  and  $2$  TeV, respectively. The Standard Model cross-section is shown as a dotted line for comparison.

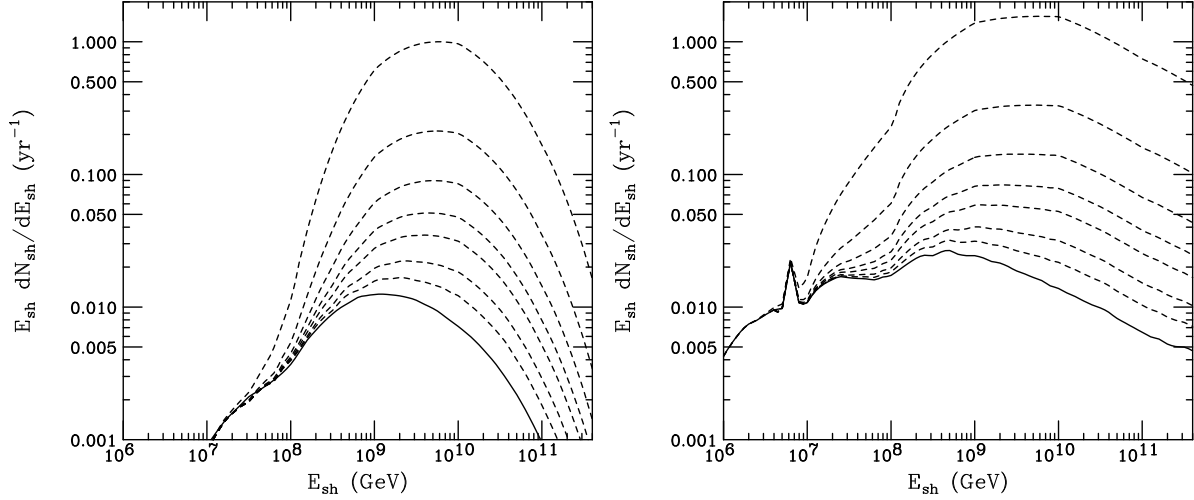


FIG. 12: The spectrum of quasi-horizontal, deeply penetrating, black hole induced showers as would be seen by Auger for the cosmogenic flux (left) and the Waxman-Bahcall flux (right). The dashed lines indicates different values of the fundamental Planck scale (from below  $M_{10} = 10, 7, 5, 4, 3, 2, 1$  TeV; in all cases  $M_{\text{BH,min}} = 3M_{10}$ ) while the solid line is the SM prediction.

evaporation do not escape the Earth's crust [96]. To quantify the potential of Auger in discriminating BH/SB induced showers, we show separately the BH production event rates for quasi-horizontal and Earth skimming neutrinos in Figs. 12 and 13. The SB production rates are similarly given in Figs. 14 and 15 and a summary of these event rates is provided in Tables IV and V respectively.

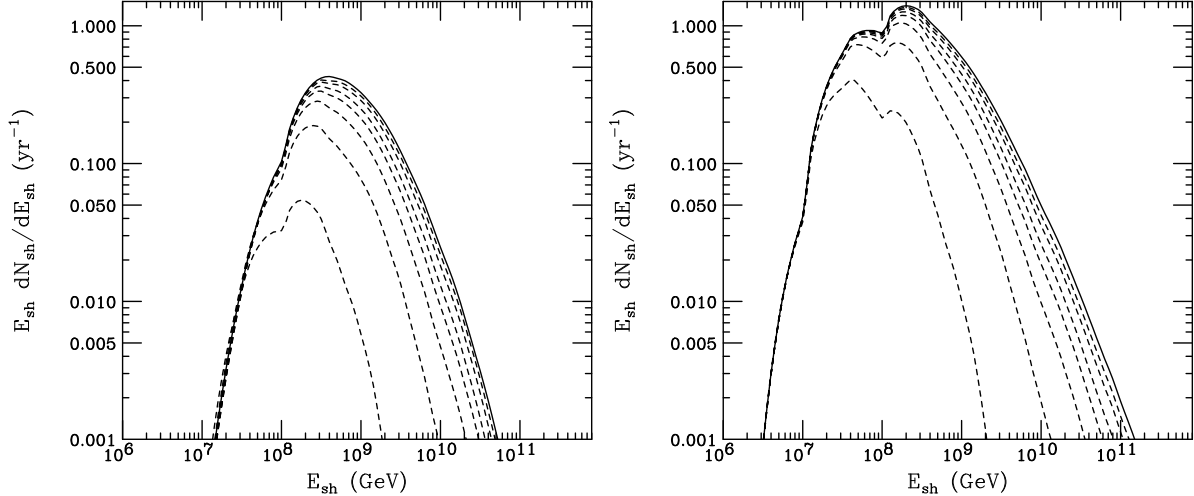


FIG. 13: The spectrum of Earth skimming, tau neutrino black hole induced showers as would be seen by Auger for the cosmogenic flux (left) and the Waxman-Bahcall flux (right). The dashed lines indicates different values of the fundamental Planck scale (from below  $M_{10} = 1, 2, 3, 4, 5, 7, 10$  TeV; in all cases  $M_{\text{BH,min}} = 3M_{10}$ ), while the solid line is the SM prediction.

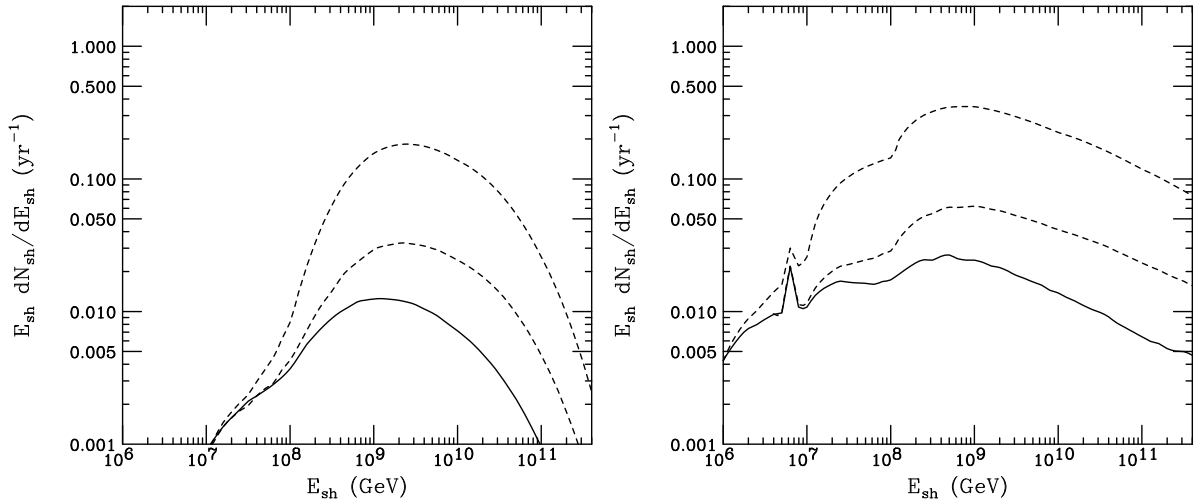


FIG. 14: The spectrum of quasi-horizontal, deeply penetrating, neutrino string-ball induced showers as would be seen by Auger for the cosmogenic flux (left) and the Waxman-Bahcall flux (right). The dashed lines refer to the string scale  $M_s = 1$  TeV (upper) and  $M_s = 2$  TeV (lower), while the solid line is the SM prediction.

## B. Non-perturbative Electroweak Interactions

The transition probability between two flat space vacua can be calculated in a Minkowski framework in analogy with WKB tunneling through non-vacuum fluctuations, or by evaluating the minimal action appropriate to a classical solution of Euclidean space in a given topological sector [97]. As is well known, in Yang-Mills theories the inclusion of massless fermions fundamentally alters the picture [98]: transitions between vacua (separated by energy barriers whose minimum height is set by the sphaleron energy  $E_{\text{sp}}$  [99]) will be totally suppressed unless accompanied by the simultaneous emission or absorption of *all* fermions coupled to the gauge field. In the Minkowski description, these fermions

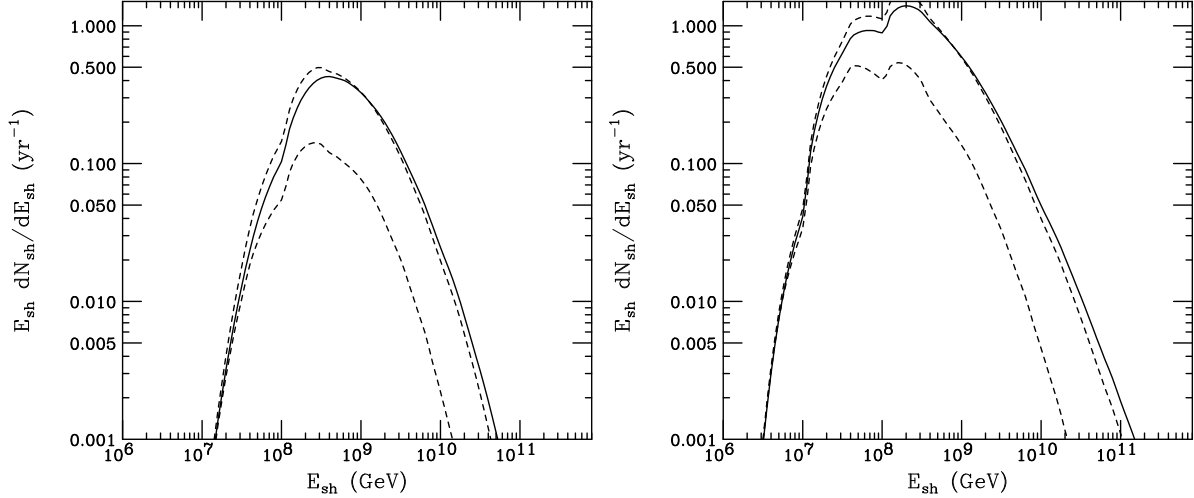


FIG. 15: The spectrum of Earth skimming, tau neutrino string ball induced showers as would be seen by Auger for the cosmogenic flux (left) and the Waxman-Bahcall flux (right). The dashed lines refer to the string scale  $M_s = 1$  TeV (upper) and  $M_s = 2$  TeV (lower), while the solid line is the SM prediction.

$\sigma_{\nu N}$	Quasi-horizontal		Earth-skimming $\nu_\tau$		Ratio	
	Cosmogenic	Waxman-Bahcall	Cosmogenic	Waxman-Bahcall	Cosmo	WB
Standard Model	0.067	0.22	1.3	5.0	0.050	0.044
$M_{10} = 1$ TeV	4.4	10.6	0.13	1.0	36	10.2
$M_{10} = 2$ TeV	0.95	2.4	0.48	2.6	2.0	0.91
$M_{10} = 3$ TeV	0.42	1.1	0.77	3.5	0.54	0.3
$M_{10} = 4$ TeV	0.25	0.66	0.96	4.1	0.26	0.16
$M_{10} = 5$ TeV	0.18	0.48	1.1	4.4	0.16	0.11
$M_{10} = 7$ TeV	0.12	0.34	1.2	4.7	0.1	0.073
$M_{10} = 10$ TeV	0.089	0.27	1.3	4.8	0.08	0.056

TABLE IV: Black hole producing event rates of quasi-horizontal showers and Earth-skimming tau neutrino induced showers expected to be observed per year by Auger for both the cosmogenic neutrino flux and the Waxman-Bahcall flux. In all cases  $M_{\text{BH},\text{min}} = 3M_{10}$ .

$\sigma_{\nu N}$	Quasi-horizontal		Earth-skimming $\nu_\tau$		Ratio	
	Cosmogenic	Waxman-Bahcall	Cosmogenic	Waxman-Bahcall	Cosmo	WB
Standard Model	0.067	0.22	1.3	5.0	0.05	0.044
$M_s = 1$ TeV	0.86	2.5	0.4	2.0	2.1	1.3
$M_s = 2$ TeV	0.17	0.48	1.5	5.7	0.11	0.084

TABLE V: String ball producing event rates of quasi-horizontal showers and Earth-skimming tau neutrino induced showers expected to be observed per year by Auger for both the cosmogenic neutrino flux and the Waxman-Bahcall flux.

emerge during level-shifting in the strong  $\mathcal{O}(1/g)$  gauge fields interpolating between vacua ( $g =$  coupling constant). In the Euclidean description, the presence of a zero mode  $\omega$  for each light fermion coupled to the gauge field will, because of the rules of Grassman

integration, generate a 't Hooft vertex [98] with all the different fermions appearing as legs,

$$\mathcal{L}_{\text{eff}} \propto \prod_{i=1\dots N} \bar{\psi} F_i + \text{h.c.} , \quad (24)$$

where  $F_i$  is a chiral fermion field. Whether these exotic processes occur with sizeable rates in high energy particle collisions is a long-standing open question [100].

At center-of-mass energies  $\sqrt{\hat{s}} < E_{\text{sp}} \approx \pi M_W / \alpha_W \approx 7.5$  TeV, the cross-section for electroweak instanton mediated processes is known to have an exponential form [101]. Here,  $m_W = 80.423$  GeV is the  $W^\pm$  boson mass and  $\alpha_W(m_W) = 0.0338$  is the  $SU(2)$  fine structure constant [45]. Including essential pre-exponential factors [102], one has, for the phenomenologically interesting case of fermion-fermion scattering  $f + f \xrightarrow{I}$  all,

$$\begin{aligned} \hat{\sigma}_{\text{ff}}^{(I)} &\approx \frac{1}{m_W^2} \left( \frac{2\pi}{\alpha_W} \right)^{7/2} \exp \left[ -\frac{4\pi}{\alpha_W} F_{\text{hg}} \left( \frac{\sqrt{\hat{s}}}{4\pi m_W / \alpha_W} \right) \right] \\ &\simeq 5.3 \times 10^3 \text{ mb} \exp \left[ -\frac{4\pi}{\alpha_W} F_{\text{hg}} \left( \frac{\sqrt{\hat{s}}}{4\pi m_W / \alpha_W} \right) \right]. \end{aligned} \quad (25)$$

where  $F_{\text{hg}}$  is the ‘‘holy-grail’’ function [103]. By means of perturbative calculations of the relevant exclusive amplitudes about the instanton ( $I$ ), squaring them and summing over the final states, or, alternatively, by means of a perturbative calculation of the forward elastic scattering amplitude about the widely separated instanton anti-instanton ( $I\bar{I}$ ) pair and determining the imaginary part to get the total cross-section via the optical theorem, one may calculate the decisive tunneling suppression exponent  $F_{\text{hg}}$ , as a series in fractional powers of  $\epsilon \equiv \sqrt{\hat{s}} / (4\pi m_W / \alpha_W) \simeq \sqrt{\hat{s}} / (30 \text{ TeV})$  [102],

$$F_{\text{hg}}(\epsilon) = 1 - \frac{3^{4/3}}{2} \epsilon^{4/3} + \frac{3}{2} \epsilon^2 + \mathcal{O}(\epsilon^{8/3}). \quad (26)$$

Therefore, the total cross-section given in Eq. (25) is exponentially growing for  $\epsilon \ll 1$ . At  $\epsilon$  of  $\mathcal{O}(1)$ , however, the perturbative expression in Eq. (26) no longer applies. In this energy regime, only extrapolations of, and lower bounds on, the tunneling suppression exponent are available [104].

Interestingly, at  $\sqrt{\hat{s}} \sim 100$  TeV, the cross-section can rise to values characteristic of QCD interactions. Since the electroweak instanton-induced interaction applies equally to all fermions, neutrinos can thus acquire hadron-like cross sections at high energies. Moreover, the inelasticity of the process is *high*. Together, these facts imply that neutrino interactions mediated by instantons would induce air showers in the upper atmosphere with characteristics similar to those of proton-induced showers [105]. Conversely, the non-observation to date of deeply penetrating air showers constrains any sudden rise of the neutrino-nucleon cross-section [106, 107]. Figure 16 shows the allowed region for transition from electroweak to QCD-like neutrino-nucleon cross-section, consistent with existing data. The dashed line indicates the neutrino-nucleon cross-section taken from Ref. [108] obtained taking  $\hat{\sigma}_{ff} \gtrsim 1$  mb [104]. As can be seen in Fig. 16, this prediction is marginally consistent with the region allowed by current data. For this cross-section, the expected event rate at Auger would be 4.3 quasi-horizontal showers per year assuming the cosmogenic neutrino flux, and 14 quasi-horizontal showers per year assuming the Waxman-Bahcall neutrino flux; the rate of Earth-skimmers is 1.3 per year in both cases. As shown in Fig. 17, the suppression of Earth-skimmers due to absorption in the Earth is negligible. However, the rate of quasi-horizontal showers is increased by about 2 orders of

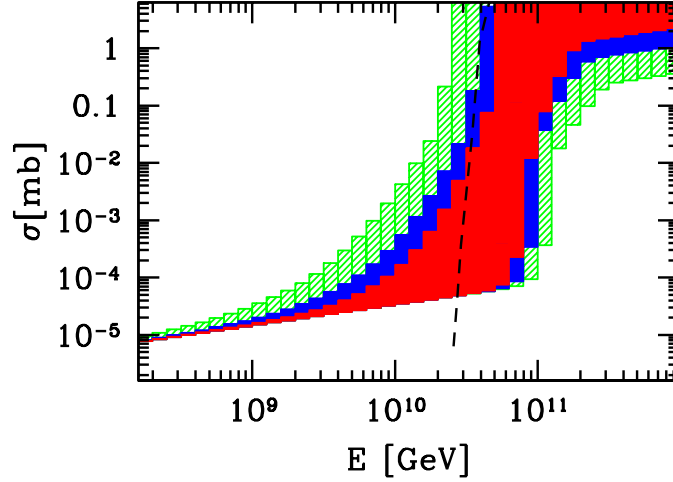


FIG. 16: The allowed 90%, 95%, and 99% CL regions for interpolation between the electroweak and QCD-like neutrino-nucleon cross-section consistent with existing data. Also shown with a dashed line is the predicted enhancement of the cross-section by electroweak sphalerons. For details see Ref. [107, 108].

magnitude, and such events would be concentrated in a small energy range, as indicated in Fig. 18. This would provide a *clean* signal for electroweak instanton-induced interactions. Thus, if no deeply developing showers are observed, tighter constraints can be placed on this model, and more generally on any sudden rise in the neutrino-nucleon cross-section.

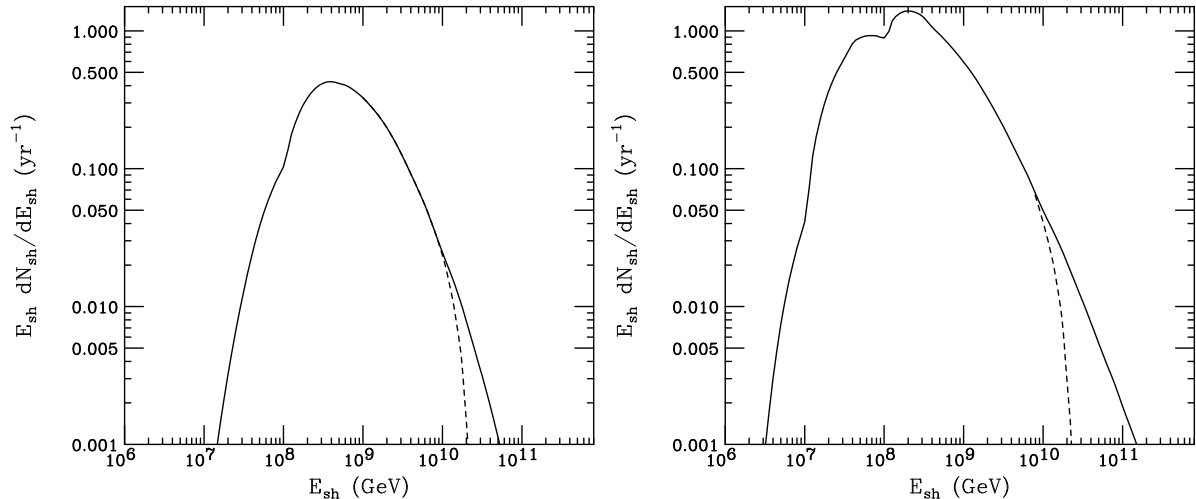


FIG. 17: Suppression of Earth skimming events due to electroweak sphalerons as would be seen by Auger for the cosmogenic neutrino flux (left), and for the Waxman-Bahcall flux (right). The solid line is the SM prediction.

### C. Neutrino Decay

Neutrinos are known to be sufficiently light that they are stable against tree-level electroweak decays. Moreover, decays of the form  $\nu_i \rightarrow \nu_j \gamma$  or  $\nu \rightarrow \nu \nu \bar{\nu}$  are severely constrained by experiment [45]. However, some models of lepton number violation postulate



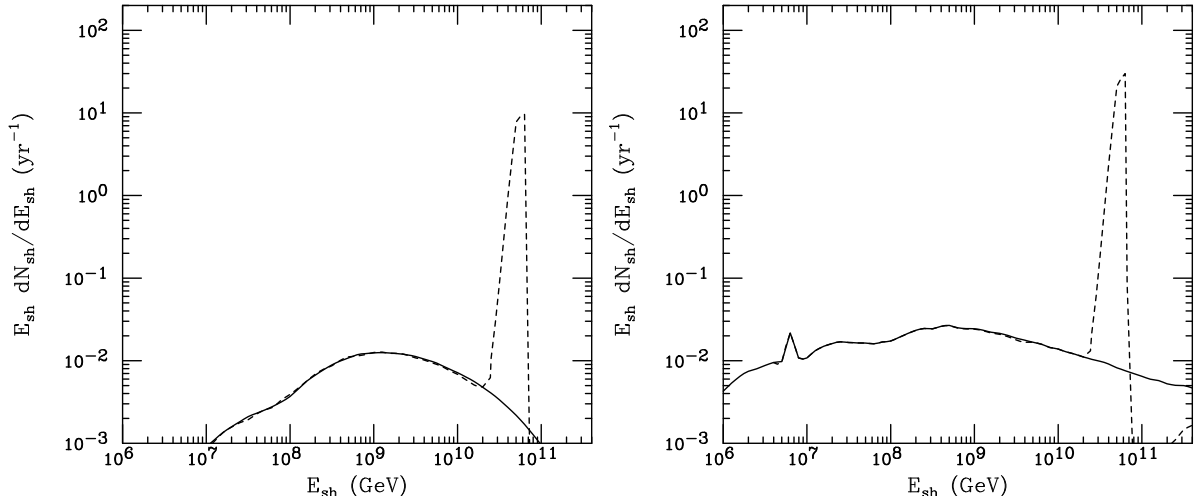


FIG. 18: The spectrum of quasi-horizontal showers mediated by electroweak sphalerons as would be seen by Auger for the cosmogenic neutrino flux (left), and for the Waxman-Bahcall flux (right). The solid line is the SM prediction.

the existence of a massless Goldstone boson, the Majoron,  $X$ . Consequently, decays such as  $\nu_i \rightarrow \nu_j X$  or  $\nu_i \rightarrow \bar{\nu}_j X$ , are then possible, where  $\nu_{i,j}$  denote mass eigenstates [109]. Presently such possibilities are only weakly constrained by Solar neutrino data, which set the bound  $\tau/m \gtrsim 10^{-4}$  s/eV [110]. However, by studying cosmic neutrinos which have travelled over far longer baselines Auger can be more sensitive to their instability by a factor of  $\sim 10^2 - 10^4$ , if an effective flavor ratio measurement can be made. Because of the extremely large energies probed by Auger, it will be complementary in this regard to the IceCube neutrino telescope [111, 112].

The ratio of flavors observed in the cosmic neutrino spectrum depends on whether any species of neutrinos have decayed and on the decay channel. In the simple situation where all heavy neutrino species decay into the lightest mass eigenstate (or into non-interacting states, such as a sterile neutrino), we would expect to observe at Earth the flavor ratio

$$\phi_{\nu_e} : \phi_{\nu_\mu} : \phi_{\nu_\tau} = \cos^2 \theta_\odot : \frac{1}{2} \sin^2 \theta_\odot : \frac{1}{2} \sin^2 \theta_\odot \approx 6 : 1 : 1, \quad (27)$$

where  $\theta_\odot$  is the solar neutrino mixing angle and we have assumed the normal hierarchy as well as  $U_{e3} = 0$ . This result is independent of the flavor ratio at source. However, for the case of an inverted hierarchy, the predicted flavor ratio at Earth is

$$\phi_{\nu_e} : \phi_{\nu_\mu} : \phi_{\nu_\tau} = U_{e3}^2 : U_{\mu 3}^2 : U_{\tau 3}^2 \approx 0 : 1 : 1, \quad (28)$$

where  $U_{\alpha i}$  is the neutrino mixing matrix and we have taken the atmospheric mixing angle to be maximal. These results are in striking contrast to the expectation for stable neutrinos discussed earlier in Sec. IV B.

These two cases represent the most extreme deviations from the usual phenomenology and are robust in that they do not depend on the flavor composition at source. A variety of other (more baroque) possibilities have been considered, e.g. only the heaviest neutrino eigenstate decays, but the predicted flavor ratios after propagation then depend on the assumed flavor ratio at source [111]. In Table VI we list some of these possibilities assuming the usual mass hierarchy and source flavor ratios as for pion decay.

We cannot measure the flavor ratios directly at Auger. However, as discussed earlier, Earth-skimming events are generated uniquely by tau neutrinos, while quasi-horizontal

showers can be generated by all neutrino flavors. Furthermore, because of maximal mixing of  $\nu_\mu$  and  $\nu_\tau$  we expect their fluxes to be always comparable. Therefore, by combining these two measurements Auger can potentially determine the flavor ratios of ultra-high energy neutrinos.

Decaying Mass Eigenstates	Decay Products	$\phi_{\nu_e} : \phi_{\nu_\mu} : \phi_{\nu_\tau}$
$\nu_3, \nu_3$	Irrelevant	6:1:1
$\nu_3$	Invisible	2:1:1
$\nu_3$	$\nu_2$	1.4–1.6:1:1
$\nu_3$	$\nu_1$	2.4–2.8:1:1
$\nu_3$	50% $\nu_1$ , 50% $\nu_2$	2:1:1

TABLE VI: The neutrino flavor ratios predicted for a variety of neutrino decay models with decay mode as indicated [111].

Of course this requires a substantial event rate. As shown in Table I the standard cosmogenic neutrino flux is expected to generate only about 0.7 quasi-horizontal shower events over 10 years. This is certainly insufficient for making the precision measurements needed to identify the effects of neutrino decay. For the nominal Waxman-Bahcall flux, Auger is expected to detect about 2.2 quasi-horizontal events and about 48 Earth-skimming events in 10 yr. However, if the cosmic ray galactic–extragalactic transition happens at around  $10^9$  GeV [113], then the required proton luminosity in the extragalactic sources increases significantly. Then Auger would detect as many as 21 quasi-horizontal and 350 Earth-skimming events over 10 years. This corresponds to a  $2\sigma$  measurement of their ratio of  $0.06 \pm 0.026$ , which would exclude anomalous flavor composition with a  $\nu_e$  content greater than  $\phi_{\nu_e} : \phi_{\nu_\mu} : \phi_{\nu_\tau} \simeq 2.5 : 1 : 1$ .

Other possibilities for altering neutrino flavor ratios have been explored [114, 115]. If Lorentz invariance is violated through modification of the usual dispersion relation for neutrinos by non-renormalizable operators induced by quantum gravity effects, then the fraction of tau neutrinos may be suppressed [115]. Auger would then observed the ratio of Earth-skimmers to quasi-horizontal events to decrease from about 20 to close to zero.

## VI. CONCLUSIONS

Our knowledge of fundamental interactions has largely been limited to the energies up to which collider experiments have been able to probe. The Tevatron, currently operating at Fermilab, produces collisions with a center-of-mass energy slightly below 2 TeV, while the Large Hadron Collider, under construction at CERN, will reach 14 TeV. By contrast, a typical neutrino observed at the Pierre Auger Observatory will have an energy of  $\mathcal{O}(10^9)$  GeV, corresponding to a neutrino-nucleon center-of-mass energy exceeding 40 TeV. Although the number of collisions which will be observed (i.e. the beam luminosity) is far below that of collider experiments, Auger and other experiments sensitive to ultra-high energy cosmic neutrinos have in principle the ability to provide unique information on new physics beyond the reach of any planned accelerator.

In addition to this advantage, cosmic neutrinos have traveled over very great distances before reaching Earth, thus their detection also constitutes an extremely long-baseline oscillation experiment. Instead of being limited to phenomena which occur over minuscule fractions of a second, cosmic neutrinos provide an exceptional window into phenomena only evident over cosmological scales of length or time.

The Pierre Auger Observatory is capable of detecting two primary classes of neutrino induced events — quasi-horizontal, deeply penetrating showers and (slightly) upgoing showers induced by Earth-skimming tau neutrinos. Used separately, the rates of such events are of limited use in probing new physics; since the spectrum of cosmic neutrinos is currently unknown, an event rate cannot by itself be used to determine the neutrino-nucleon interaction cross-section. However by combining these two classes of neutrino-induced events, it becomes possible to make a crude cross-section measurement. As this cross-section is increased (decreased), the rate of quasi-horizontal showers increases (decreases) accordingly, while by contrast, the rate of slightly upgoing showers is reduced (enhanced) since Earth-skimming tau neutrinos become absorbed more (less). Thus the ratio of quasi-horizontal, deeply penetrating showers to slightly upgoing showers provides a check of the behaviour of the neutrino-nucleon cross-section at ultra-high energies. The details of such a measurement, of course, depend on the energy dependence of such interactions, as well as their inelasticity and other characteristics.

These two types of neutrino-induced events also provide the opportunity to constrain the ratios of flavors present in the ultra-high energy cosmic neutrino spectrum. If these neutrinos are generated through the decay of charged pions (as they are in most models), they will reach Earth in nearly equal quantities of each flavor after oscillations are taken into account. A larger than expected rate of quasi-horizontal, deeply penetrating showers, in comparison to the slightly upgoing shower rate, would thus indicate a suppression of the tau neutrino component in the ultra-high energy cosmic neutrino spectrum, due, for example, to neutrino decay.

In this study, we have considered several specific models in which either the neutrino-nucleon cross-section, or the ratio of cosmic neutrino flavors, deviates substantially from the expectation of the perturbative Standard Model. We have studied enhancements in the neutrino-nucleon cross-section in models with low-scale gravity, variously described as due to the exchange of Kaluza-Klein gravitons, the production of microscopic black holes, and/or string resonances. We have also considered increases in the cross-section due to non-perturbative Standard Model electroweak instanton induced processes, which in contrast do not lead to a decrease in the inelasticity. Regarding flavor ratio measurements, we have discussed several models of decaying neutrinos.

It is difficult to precisely delineate the reach of these techniques as this depends on the unknown flux of cosmic neutrinos at the energies to which Auger is sensitive. We have considered both the “guaranteed” cosmogenic flux which sets a lower bound and the Waxman-Bahcall flux which sets an upper bound. Further observations of ultra-high energy cosmic rays by Auger itself will help to pin down the expected neutrino flux.

Over the next few years, the Pierre Auger Observatory may well identify the world’s first ultra-high energy neutrino event. We have attempted to illustrate the exciting new possibilities for probing new physics that will be opened up by such a detection.

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