

Black holes from cosmic rays: Probes of extra dimensions and new limits on TeV-scale gravityLuis A. Anchordoqui,¹ Jonathan L. Feng,^{2,3} Haim Goldberg,¹ and Alfred D. Shapere⁴¹*Department of Physics, Northeastern University, Boston, Massachusetts 02115*²*Center for Theoretical Physics, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139*³*Department of Physics and Astronomy, University of California, Irvine, California 92697*⁴*Department of Physics, University of Kentucky, Lexington, Kentucky 40502*

(Received 13 February 2002; published 19 June 2002)

If extra spacetime dimensions and low-scale gravity exist, black holes will be produced in observable collisions of elementary particles. For the next several years, ultrahigh-energy cosmic rays provide the most promising window on this phenomenon. In particular, cosmic neutrinos can produce black holes deep in the Earth's atmosphere, leading to quasihorizontal giant air showers. We determine the sensitivity of cosmic ray detectors to black hole production and compare the results to other probes of extra dimensions. With $n \geq 4$ extra dimensions, current bounds on deeply penetrating showers from AGASA already provide the most stringent bound on low-scale gravity, requiring a fundamental Planck scale $M_D > 1.3\text{--}1.8$ TeV. The Auger Observatory will probe M_D as large as 4 TeV and may observe on the order of a hundred black holes in 5 years. We also consider the implications of angular momentum and possible exponentially suppressed parton cross sections; including these effects, large black hole rates are still possible. Finally, we demonstrate that even if only a few black hole events are observed, a standard model interpretation may be excluded by comparison with Earth-skimming neutrino rates.

DOI: 10.1103/PhysRevD.65.124027

PACS number(s): 04.70.-s, 04.50.+h, 13.15.+g, 96.40.Tv

I. INTRODUCTION

Tiny black holes (BHs) can be produced in particle collisions with center-of-mass energies above the fundamental scale of gravity [1,2], where they should be well-described semiclassically and thermodynamically [3]. In conventional 4-dimensional theories, viz., where the Planck scale $\sim 10^{19}$ GeV is fundamental and the weak scale ~ 1 TeV is derived from it via some dynamical mechanism, the study of such BHs is far beyond the realm of experimental particle physics. Over the past few years, however, physicists have begun exploring an alternative approach to the long-standing gauge hierarchy problem, wherein the weak scale becomes the fundamental scale of nature and the Planck scale is derived from this, with the hierarchy in scales a consequence of large or warped extra dimensions [4,5]. If this is the case, the fundamental scale of gravity can be $\mathcal{O}(\text{TeV})$, and BH production and evaporation may be observed in collisions of elementary particles [6–10].

If gravity indeed becomes strong at the TeV scale, ultrahigh-energy cosmic rays provide a powerful opportunity to probe BH production at super-Planckian energies [11]. Cosmic rays with energies $\sim 10^{19}$ eV have been observed [12]. They interact in the Earth's atmosphere and crust with center-of-mass energies ~ 100 TeV, far beyond the reach of present and planned manmade colliders. These cosmic rays may therefore produce BHs, allowing cosmic ray detectors to test the existence of TeV-scale gravity and extra dimensions by searching for evidence of BH production [11,13–16]. A particularly promising signal is provided by ultrahigh-energy cosmic neutrinos, which may produce BHs with cross sections two or more orders of magnitude above their standard model (SM) interactions. These BHs will decay promptly in a thermal distribution of SM particles. Of the order of a hundred BH events may be detected at the Auger Observa-

tory [11] as quasihorizontal, deeply penetrating showers with distinctive properties [13]. The possibility of BH production by cosmic rays supplements possible sub-Planckian signatures of low-scale gravity [17–21].

In this article we extend previous work to derive bounds from the nonobservation of BH-initiated showers in current data at the Akeno Giant Air Shower Array (AGASA). We also extend previous analyses of BH discovery prospects at Auger, and discuss in detail the possibility of distinguishing BH events from SM events. A preliminary version of some of these results was presented in Ref. [22].

We begin in Sec. II with an overview of TeV-scale gravity. We collect and review existing bounds on the fundamental Planck scale in a uniform convention. In Sec. III we discuss semiclassical BH production, including the effects of angular momentum and the production of Kerr BHs, as well as the proposed exponential suppression advocated by Voloshin [23,24]. This is followed in Secs. IV and V by detailed discussions of cosmogenic neutrino fluxes and ground array experiments, respectively.

Our results for event rates and new limits on the scale of higher-dimensional gravity are presented in Secs. VI and VII. We begin with current data from AGASA. The AGASA Collaboration has already reported no significant signal for neutrino air showers during an observation time (live) of 1710.5 days [25]. Given the standard assumption of a geometric black hole cross section, we find that this data implies the most stringent bound on the fundamental Planck scale to date for $n \geq 4$ extra dimensions, exceeding limits derived [16] from Fly's Eye data [26] and also more stringent than the constraints from graviton emission and exchange obtained by the CERN e^+e^- collider LEP [27] and $D\bar{O}$ [28] Collaborations. In Sec. VII we then consider the prospects for BH production at the Auger Observatory. Tens of black holes may be observed per year; conversely, nonobservation of

BHs will imply bounds as large as 4 TeV on the fundamental Planck scale.

In Sec. VIII we note that comparison to Earth-skimming neutrino event rates [29–32] allows one to distinguish BH events from SM events. This point was noted already in Ref. [11], but was not considered in Ref. [16], leading to weaker conclusions. Here, we consider this point quantitatively and find that, even with a handful of BH events, a SM explanation may be excluded based on event rates alone. If seen, black holes created by cosmic rays will provide the first evidence for extra dimensions and TeV-scale gravity, initiating an era of detailed study of black hole properties at both cosmic ray detectors and future colliders, such as the CERN Large Hadron Collider (LHC) [9,10,33–37]. Our conclusions are collected in Sec. IX.

II. EXISTING LIMITS ON LOW-SCALE GRAVITY

Depending on the dimensionality and the particular form of spacetime, the gauge hierarchy problem may be reexpressed as a hierarchy in length scales. In the canonical example [4], spacetime is a direct product of a noncompact 4-dimensional spacetime manifold and a flat spatial n -torus of common linear size $2\pi r_c$ and volume $V_n = (2\pi r_c)^n$. Only gravity propagates in the full $(4+n)$ -dimensional spacetime; all other fields are confined to a 3-brane extended in the noncompact dimensions. Here, the low energy 4-dimensional Planck scale M_{Pl} is related to the fundamental scale of gravity in $(4+n)$ dimensions, M_* , according to

$$M_{\text{Pl}}^2 = M_*^{2+n} V_n = V_n / G_{(4+n)}, \quad (1)$$

with $G_{(4+n)}$ defined by the $(4+n)$ -dimensional Einstein field equation $R_{AB} - \frac{1}{2}g_{AB} = -8\pi G_{(4+n)} T_{AB}$. In what follows it will be convenient to work with the mass scale [38]

$$M_D = [(2\pi)^n / 8\pi]^{1/(n+2)} M_*, \quad D = 4 + n. \quad (2)$$

If r_c is significantly larger than the Planck length, a hierarchy is introduced between M_{Pl} and M_D , and gravity becomes strong in the entire $(4+n)$ -dimensional spacetime at the scale M_D far below the conventional Planck scale $M_{\text{Pl}} \sim 10^{19}$ GeV. Our conclusions will be essentially unchanged for more general “asymmetric” compactifications, with, e.g., p “small” dimensions with sizes $\lesssim \text{TeV}^{-1}$ and $n_{\text{eff}} = n - p$ large extra dimensions [39]. (Note, however, that in this case, the production of brane configurations wrapped around small extra dimensions may be competitive with black hole production [40].) Many of our results for black hole production and detection also apply for warped compactifications [5] in which the curvature length is much larger than a TeV^{-1} . Hereafter, we will focus our discussion on bounds in flat compactification scenarios. In the figures, for $n=1$ results, warped scenarios are implicit.

A. Bounds from Newtonian gravity

The provocative new features of these scenarios have motivated many phenomenological studies to assess their experimental viability. Naturally, the most obvious conse-

quence of the existence of large compact dimensions is the deviation from Newtonian gravity at distances of order r_c . For $n=1$ and $M_D \sim 1$ TeV, $r_c \sim 10^{13}$ cm, implying deviations from Newtonian gravity over solar system distances, so this case is empirically excluded. For $n=2$, sub-millimeter tests of the gravitational inverse-square law constrain $M_D > 1.6$ TeV [41]. For $n \geq 3$, r_c becomes microscopic and therefore eludes the search for deviations in gravitational measurements.

B. Astrophysical bounds

In the presence of large compact dimensions, however, the effects of gravity are enhanced at high energies, due to the accessibility of numerous excited states of the graviton [referred to as Kaluza-Klein (KK) gravitons [42]], corresponding to excitations of the graviton field in the compactified dimensions. For low numbers of extra dimensions, by far the most restrictive limits on the radii of large compact dimensions come from the effects of KK graviton emission on cooling of supernovae, and from neutron star heating by KK decays [43]. For $n=2$ the latter requires $M_D > 600\text{--}1800$ TeV, far above the weak scale; for $n=3$, the bound is $M_D > 10\text{--}100$ TeV. These limits apply only for the situation where all extra dimensions have the same compactification radius. In the general case, the bounds could be less restrictive.

C. Collider bounds

For $n \geq 4$ extra dimensions, only high energy collisions are useful as probes. The effects of direct graviton emission, including production of single photons or Z 's, were sought at LEP [44]. The resulting bounds are fairly model-independent, as the relatively low energies at LEP imply a negligible dependence on the soft-brane damping factor discussed below. For $n=4(6)$, these null results imply $M_D > 870(610)$ GeV [27].

The effects of low-scale gravity can also be seen through virtual graviton effects. These are most stringently bounded by the DØ Collaboration, which recently reported [28] the first results for virtual graviton effects at a hadron collider. The data collected at $\sqrt{s} = 1.8$ TeV for dielectron and diphoton production at the Fermilab Tevatron agree well with the SM predictions and provide the most restrictive limits on an *effective* extra-dimensional Planck scale for $n \geq 4$. This scale (called Λ_T in [38], and related to M_S in [45,46]) simply parametrizes the KK graviton exchange amplitudes for these processes: except for the different conventions used, they simply convey the experimental limit in terms of an energy-independent four-point function. In the context of low-scale gravity, the effective scale depends on both $G_{(4+n)}$ and on an ultraviolet cutoff on the contributing KK modes [38,45,46]. This cutoff represents the energy where emission of graviton modes from the brane into the extra dimensions are damped by the effects of a nonrigid brane, and it is expected to be of order $G_{(4+n)}^{-1/(n+2)}$.

In this work we will use a Gaussian cutoff [47,48], which emerges if one includes in the interaction the brane Gold-

stone modes. With this cutoff, a form factor $e^{-m^2/2\Lambda^2}$ is introduced at each graviton-matter vertex, where m is the mass of the graviton and Λ parametrizes the cutoff. In real graviton emission processes, the effect of the cutoff is somewhat alleviated because of finite cuts on the missing energy. However, at LHC energies, the corrections become significant [49] in the expected region $\Lambda < M_D$, and are of order 100% when $\Lambda/M_D \approx 0.5$. For virtual processes, the s -channel diphoton or dielectron amplitude has the form [38,45]

$$\begin{aligned} \mathcal{A} &= \mathcal{S}(\hat{s}) \mathcal{T} \\ \mathcal{S}(\hat{s}) &= \frac{8\pi}{M_{\text{Pl}}^2} \sum_{\vec{\ell}} \frac{1}{m^2 - \hat{s}} \\ \mathcal{T} &= T_{\mu\nu} T^{\mu\nu} - \frac{1}{n+2} T_{\mu}^{\mu} T_{\nu}^{\nu}, \end{aligned} \quad (3)$$

where the sum on $\vec{\ell}$ denotes a sum over the KK graviton modes, labeled by an n -dimensional lattice vector $\vec{\ell}$, with graviton masses $m = |\vec{\ell}|/r_c$. The first and second $T_{\mu\nu}$'s are the stress tensors for the incoming $q\bar{q}, gg$ and outgoing $e^+e^-, \gamma\gamma$ states, respectively, and $\sqrt{\hat{s}}$ is the parton center-of-mass energy. The sum on $\vec{\ell}$ may be approximated by a continuous integration over KK masses, modified by the cutoff, with the result [38,45]

$$\mathcal{S}(\hat{s}) = \frac{S_{n-1}}{M_D^{2+n}} \int_0^{\infty} \frac{m^{n-1} dm e^{-m^2/\Lambda^2}}{m^2 - \hat{s}}, \quad (4)$$

where explicit integration over the $n-1$ angular variables leads to the factor $S_{n-1} = 2\pi^{n/2}/\Gamma(n/2)$. The connection to an effective four-point contact interaction in Refs. [38,45,46] is made by setting $\hat{s} = 0$ in Eq. (4). This allows an explicit evaluation of the integration over m , with the result (for $n \geq 3$)

$$\mathcal{S}(0) = \pi^{n/2} \frac{2}{n-2} \left(\frac{\Lambda}{M_D} \right)^{n-2} \frac{1}{M_D^4} \equiv \frac{4\pi}{\Lambda_T^4}. \quad (5)$$

In the last line we have used the convention of Ref. [38] to parametrize the four-point amplitude. At 95% C.L., the Tevatron data require $\Lambda_T > 1.2$ TeV. With the use of Eq. (5), this allows us to generate Table I, which shows the bounds on M_D for $n=3, \dots, 7$ and $0.5 \leq \Lambda/M_D \leq 1$. It is important to note that (except for small variations for the case of Ref. [46], which permits a sign ambiguity in the amplitude) *Table I is independent of the conventions in [38,45,46]*. We can see that the lower bounds on M_D depend on both n and Λ . Typically, $M_{D,\text{min}} \lesssim 1$ TeV.

III. BH PRODUCTION IN PARTICLE COLLISIONS

The preceding section discussed some potentially observable consequences of scenarios with TeV-scale gravity, and the limits on the scale of higher-dimensional gravity result-

TABLE I. 95% C.L. lower limits on M_D from the DØ Collaboration at the Tevatron.

Λ/M_D	$M_{D,\text{min}}$ (TeV)				
	$n=3$	$n=4$	$n=5$	$n=6$	$n=7$
0.5	0.98	0.80	0.70	0.63	0.58
0.6	1.02	0.88	0.80	0.76	0.73
0.7	1.06	0.95	0.90	0.89	0.88
0.8	1.10	1.01	1.00	1.01	1.04
0.9	1.13	1.07	1.09	1.14	1.21
1.0	1.16	1.13	1.18	1.26	1.38

ing from their nonobservation. In particular, for $n=4$ to 7, the quoted lower limit on M_D comes from the nonobservation at DØ of processes involving KK gravitons. The spectrum and interactions of KK gravitons are model-dependent, to an increasing degree at increasing scales above M_D . Here we describe a more universal and model-independent prediction of low-scale gravity scenarios: the production in particle collisions of microscopic BHs.

A. Geometric cross section

It has been argued [6] that BH formation should occur in the scattering of two incident particles when their impact parameter is approximately less than the Schwarzschild radius of a BH of mass equal to their center-of-mass energy $\sqrt{\hat{s}}$. This suggests a geometric cross section

$$\hat{\sigma} \approx \pi r_s^2, \quad (6)$$

where

$$r_s(M_{\text{BH}}) = \frac{1}{M_D} \left[\frac{M_{\text{BH}}}{M_D} \right]^{1/(1+n)} \left[\frac{2^n \pi^{(n-3)/2} \Gamma\left(\frac{n+3}{2}\right)}{n+2} \right]^{1/(1+n)} \quad (7)$$

is the radius of a Schwarzschild BH of mass $M_{\text{BH}} = \sqrt{\hat{s}}$ [50,51] in $4+n$ dimensions. Even if the incident particles are stuck on the SM brane, the black hole formed should be treated as a fully $(4+n)$ -dimensional object in an asymptotically Minkowskian spacetime, as long as r_s is small compared to r_c .

The cross section Eq. (6) grows like $\hat{s}^{1/(n+1)}$, more rapidly than any SM cross section. Thus, at energies sufficiently far above M_D , BH production is expected to become the dominant process. In pp collisions at the LHC, rates as high as 10^8 events per year have been predicted in scenarios with $M_D \sim 1$ TeV [9,10].

In our investigation of BH production by cosmic rays, we will be most interested in collisions of neutrinos with atmospheric nucleons. Since, at the energy scale of interest, gravitational cross sections will be far smaller than the geometric area of a parton, we write the νN cross section as [11]

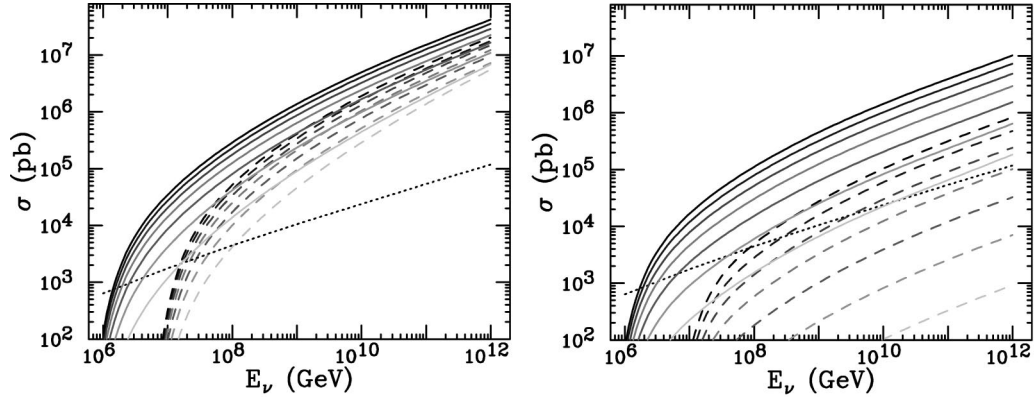


FIG. 1. Cross sections $\sigma(\nu N \rightarrow \text{BH})$ for $n=1, \dots, 7$ from below for $M_D=1$ TeV, $x_{\min}=1$ (solid) and 3 (dashed), and parton cross sections πr_s^2 (left) and $\pi r_s^2 e^{-1E}$ (right). The SM cross section $\sigma(\nu N \rightarrow \ell X)$ (dotted) is also shown.

$$\sigma(\nu N \rightarrow \text{BH}) = \sum_i \int_{(M_{\text{BH}}^{\min})^{2/s}}^1 dx \hat{\sigma}_i(\sqrt{xs}) f_i(x, Q), \quad (8)$$

where $s=2m_N E_\nu$, the sum is over all partons in the nucleon, and the f_i are parton distribution functions. We set the momentum transfer $Q = \min\{M_{\text{BH}}, 10 \text{ TeV}\}$, where the upper limit is from the CTEQ5M1 distribution functions [52]. The cross section $\sigma(\nu N \rightarrow \text{BH})$ is highly insensitive to the details of this choice [11]. For example, choosing instead $Q = \min\{r_s^{-1}, 10 \text{ TeV}\}$ [14] changes BH production rates by only 10% to 20%. For the conservative ν fluxes considered below, our results are also insensitive to low x . [For concreteness, however, we extrapolate to $x < 10^{-5}$ assuming a power law behavior $f_i(x, Q) \propto x^{-[1+\lambda_i(Q)]}$.] Finally, M_{BH}^{\min} is the minimal BH mass for which Eq. (6) is expected to be valid. The appropriate choice of M_{BH}^{\min} is subject to theoretical uncertainties, as discussed below. We define

$$x_{\min} \equiv M_{\text{BH}}^{\min}/M_D \quad (9)$$

and present results for various $1 \leq x_{\min} \leq 5$. The dependence of our event rates on x_{\min} is found to be rather mild.

Cross sections for BH production by cosmic neutrinos are given in Fig. 1 for $M_D=1$ TeV; they scale as

$$\sigma(\nu N \rightarrow \text{BH}) \propto \left[\frac{1}{M_D^2} \right]^{(2+n)/(1+n)}. \quad (10)$$

The SM cross section for $\nu N \rightarrow \ell X$ is also included for comparison. (Note that cross sections rise with increasing n for fixed M_D , whereas they decrease for increasing n for fixed M_* .) As noted in Ref. [11], in contrast to the SM process, BH production is not suppressed by perturbative couplings and is enhanced by the sum over all partons, particularly the gluon. As a result of these effects, BH production may exceed deep inelastic scattering rates in the SM by two or more orders of magnitude.

Although greatly reduced by the cross section for BH production, neutrino interaction lengths

$$L = 1.7 \times 10^7 \text{ kmwe} \left(\frac{\text{pb}}{\sigma} \right) \quad (11)$$

are still far larger than the Earth's atmospheric depth, which is only 0.36 kmwe even when traversed horizontally. Neutrinos therefore produce BHs uniformly at all atmospheric depths. As a result, the most promising signal of BH creation by cosmic rays will be quasihorizontal showers initiated by neutrinos deep in the atmosphere. For showers with large enough zenith angles, the likelihood of interaction is maximized and the background from hadronic cosmic rays is eliminated, since these shower high in the atmosphere.

Once produced, the BH will Hawking evaporate, provided the semiclassical approximation is valid. In this case, a Schwarzschild BH will behave like a thermodynamic system with temperature

$$T_H = \frac{n+1}{4\pi r_s} \quad (12)$$

and entropy

$$S = \frac{2\pi^{(n+3)/2}}{4G_D \Gamma\left(\frac{n+3}{2}\right)} r_s^{n+2} = \frac{4\pi M_{\text{BH}} r_s}{n+2}. \quad (13)$$

According to the semiclassical description, a BH produced in a scattering event should be regarded as an intermediate state, which decays on a time scale

$$\tau \sim \frac{1}{M_D} \left(\frac{M_{\text{BH}}}{M_D} \right)^{(3+n)/(1+n)}. \quad (14)$$

Since $\tau < 10^{-25}$ s for $M_D \geq 1$ TeV and $M_{\text{BH}} \leq 10$ TeV, the decay is effectively instantaneous. In the decay process, particles will be radiated into all available SM channels, into quanta with energies typically of order T_H or above. These decays are predicted to lead to highly distinctive signals in collider events [9,10,33], with high multiplicity, large transverse energy, hard leptons and jets, and a characteristic ratio of hadronic to leptonic activity.

The magnitude of the entropy determines the validity of this picture. Thermal fluctuations due to particle emission are small when $S \gg 1$ [53], and statistical fluctuations in the microcanonical ensemble are small for $\sqrt{S} \gg 1$ [9]. For

$M_{\text{BH}}/M_D=5$, Eq. (13) gives S ranging from 29 for $n=4$ to 25 for $n=7$. For $M_{\text{BH}}/M_D=3$ (or 1), S is about 13 (or 3) for n between 4 and 7.

In searches for BH mediated events at colliders, it is essential to set x_{min} high enough that the decay branching ratios predicted by the semiclassical picture of BH evaporation are reliable, as there are very large QCD backgrounds, and the extraction of signal from background relies on knowing the BH decay branching ratios reliably. This is especially true if one is attempting to determine discovery limits, where the overall rates for BH production are not necessarily large. Thus, in collider searches, a cutoff of $x_{\text{min}}=5$ or more may be appropriate.

By contrast, the search for deeply penetrating quasihorizontal showers initiated by BH decays can afford to be much less concerned with the details of the final state, since the background is, relative to colliders, almost nonexistent. As a result, the signal relies only on the existence of visible decay products, which, in this context, includes all particles other than neutrinos, muons, and gravitons. Indeed, there is very little about the final state, other than its total energy and to some degree its multiplicity and electromagnetic component [13], that we can reasonably expect to observe, since detailed reconstruction of prompt decay particles is not possible at cosmic ray detectors. Thus, it seems reasonable to choose a significantly lower value of $M_{\text{BH}}^{\text{min}}$ than is needed for collider searches; in our estimates of rates below we will take x_{min} as low as 1. While BHs of mass around M_D will be outside the semiclassical regime, it seems quite reasonable to expect that they, or their stringy progenitors, will nevertheless decay visibly, whatever stringy or quantum gravitational description applies.

As an illustration, we examine the scattering in the string regime. For $n=6$ large extra dimensions, $M_s \sim g_s^{1/4} M_*$ (M_s = string scale; g_s = string coupling). As shown in [1], the string cross section σ saturates to $\sim 1/M_s^2$ for $\sqrt{\hat{s}} > M_s/g_s$ (or in terms of M_* , $\sigma \sim M_*^{-2} g_s^{-1/2}$ for $\sqrt{\hat{s}} > M_*/g_s^{3/4}$). As noted in [36] this matches the classical black hole cross section at an energy $M_s/g_s^2 \sim M_*/g_s^{7/4}$. Thus, if g_s is not too small (implying a small hierarchy between M_s and M_*), the transition to the geometric cross section is rapid. In this work we adopt a minimum energy $\sqrt{\hat{s}} \sim M_D \approx 3M_*$ (for $n=6$), so that while there are probably stringy corrections, the cross section should be substantially geometric.

B. Uncertainties in the cross section

Although the details of the process by which BHs decay are not of great concern to us, the production process is of central importance, since our rates (and the lower limits we will be able to set on M_D) will depend directly on the form of the cross section. It should be emphasized that the heuristic arguments on which Eq. (6) is based only determine σ up to an overall factor of order one. Even at the classical level, our conclusions could be significantly affected by theoretical uncertainties in this factor, four sources of which we now discuss.

1. Mass ejection

Classical general relativity calculations [54] indicate that the mass of a BH formed in a head-on collision is somewhat less (about 16% less) than the total center-of-mass energy. At least in four dimensions, this suggests that the formula Eq. (6) should be modified by replacing $r_s(M_{\text{BH}})$ by $r_s(0.84M_{\text{BH}})$, leading to a slight reduction of σ . Very recently, corresponding calculations in more than four dimensions have been presented [55].

2. Angular momentum

The analytic techniques used to study head-on collisions are not applicable to collisions at nonzero impact parameter. Thus the claim that a BH will be produced when $b \lesssim r_s(\sqrt{\hat{s}})$ can only be expected to be true up to a numerical factor.

One issue that arises at nonzero impact parameter that we can address is that the BHs formed will have angular momentum [34]. In particular, the Schwarzschild radius appearing in the formula Eq. (7) should more accurately be replaced by the radius of a Kerr BH of the appropriate angular momentum. This will alter the critical impact parameter at which a BH will form, for given \hat{s} . For two particles each of energy E in the center-of-mass frame colliding with impact parameter b , the total angular momentum with respect to the center of mass is $J=2(b/2)E=bM_{\text{BH}}/2$. So the maximum impact parameter at which a BH will form should occur at a value of b for which the radius $r_k(M, J=bM_{\text{BH}}/2)$ of a Kerr BH and b are equal. The Kerr radius satisfies [50]

$$\begin{aligned} M_{\text{BH}} &= c_n r_k^{n-1} [r_k^2 + (n+2)^2 J^2 / 4M_{\text{BH}}^2] \\ &= c_n r_k^{n-1} [r_k^2 + (n+2)^2 b^2 / 16] \end{aligned} \quad (15)$$

with c_n an n -dependent constant. Setting $r_k=b$ we get

$$M_{\text{BH}} = c_n r_k^{n+1} [1 + (n+2)^2 / 16]. \quad (16)$$

Since for a Schwarzschild BH, $M_{\text{BH}} = c_n r_s^{n+1}$, the cross section in Eq. (6) should be corrected to

$$\begin{aligned} \hat{\sigma} &\approx \pi r_k^2(M_{\text{BH}}, J) \\ &= [1 + (n+2)^2 / 16]^{-2(n+1)} \pi r_s^2(M_{\text{BH}}). \end{aligned} \quad (17)$$

For $1 \leq n \leq 7$, the correction factor is remarkably stable, ranging from 0.62 to 0.64. This result has been recently confirmed [55] in a classical analysis of black hole formation for collisions with nonzero impact parameter. We see, then, that including the effect of angular momentum also leads to a small reduction of σ .

3. Subrelativistic limit

While the corrections related to mass ejection and angular momentum both appear to decrease σ by factors of order 1, another potential correction to the naive cross section Eq. (6) could enhance it by a more-than-compensating factor. Namely, the critical impact parameter may be somewhat larger than the radius of the BH formed. An argument that

supports this conjecture may be given in the case where the incident particles are subrelativistic with rest mass approximately $E = M_{\text{BH}}/2$. In this case, the incident particles may be treated as BHs with mass $M_{\text{BH}}/2$. If $b \leq 2r_s(M_{\text{BH}}/2)$, they will touch as they pass, and thus merge into a BH of mass M_{BH} . In this regime, we expect as an approximate lower bound on σ of

$$\sigma \geq \pi b^2 = \pi [2r_s(M_{\text{BH}}/2)]^2 = 4^{n/(n+1)} \pi r_s(M_{\text{BH}})^2. \quad (18)$$

For large n , the correction factor approaches 4. Of course, the situation may change considerably in the ultrarelativistic limit, but this estimate, in a limit that we understand, at least makes it plausible that the correct coefficient in Eq. (6) might be larger than 1. Thus our choice to take $\sigma = \pi r_s^2$ may well turn out to be conservative.

4. Gravitational infall and capture

It could also be argued that the cross section of Eq. (6) is too small, because it supposes that a black hole only forms when the two particles come within a Schwarzschild radius of each other, when in fact we expect that gravitational collapse will occur for somewhat larger impact parameters as well. Another problem with Eq. (6) is that while the cross section is measured with respect to the flat geometry of the asymptotic region, the Schwarzschild radius is a property of the highly curved region close to the singularity [56].

A better measure of the cross-sectional area associated with a black hole of given mass, which overcomes these objections, is given by the classical cross section for photon capture [57]. If a beam of parallel light rays is sent in towards a Schwarzschild black hole from the asymptotically Minkowskian region of spacetime, the black hole's classical cross section is defined to be the cross-sectional area of the portion of the beam that gets captured. In four dimensions, one finds from the geodesic equation (see, e.g., Ref. [57]) that

$$\sigma = \pi b_c^2 = 27\pi G_4^2 M_{\text{BH}}^2, \quad (19)$$

independent of the energies of the incoming photons. (The relevance of this cross section for black hole production has been independently argued in Ref. [58].) The maximum impact parameter b_c at which capture occurs is about 2.6 times as large as the Schwarzschild radius, and the cross section is enhanced by a factor of 27/4. A straightforward extension of this calculation to Schwarzschild black holes in $4+n$ dimensions gives [7]

$$b_c = \frac{(3+n)^{(1/2)((3+n)/(1+n))}}{2^{1/(1+n)} \sqrt{1+n}} r_s. \quad (20)$$

For $n=1$, the cross section is 4 times larger than the cross section of Eq. (6), and for $n=7$, it is still 87% larger. Thus

the enhancement suggested by this definition of the cross section could be significant and could easily offset other possible reductions.¹

C. Exponential suppression?

Quantum mechanical corrections to the amplitude for BH production may be even more significant than classical uncertainties. In particular, Voloshin [23,24] has proposed that the cross section of Eq. (6) should be modified by an exponential suppression factor

$$\sigma \sim \pi r_s^2 e^{-I_E}, \quad (21)$$

with I_E the Euclidean Gibbons-Hawking action for the BH, which, in terms of the entropy of Eq. (13), is

$$I_E = \frac{S}{n+1}. \quad (22)$$

In part, Voloshin's critique is based on previous attempts to calculate amplitudes for the production of classical field configurations (in which the multiplicity is greater than the inverse coupling) from initial quantum states. The intrinsically nonperturbative nature of such processes suggests employing an instantonlike approximation, which could lead to exponential suppression. Such an approach can be taken even for processes that are semiclassically allowed (such as multi-Higgs production [59,60]) that do not require tunneling in order to take place. The problem is not yet solved. For example, a recent lattice simulation [61] shows no evidence for the enhancement of large multiplicity amplitudes manifest in perturbation theory, perhaps counter-indicating the formation of a classical field state in the quantum collision.

We are cognizant of the uniqueness of gravitation (such as the onset of strong coupling for $\hat{s} > M_D^2$), and the support in favor of the geometric cross section based on classical calculations for both vanishing [54] and nonvanishing [55] impact parameter. An additional supporting argument based on a string calculation has been given [36], and the applicability of CPT arguments when comparing black hole formation and decay, an element in Voloshin's criticism [23], has been questioned [58,62]. Nevertheless, for completeness, we will also present results below for the exponentially suppressed cross section of Eq. (21). For cosmic rays, we will see that even with Voloshin's suppression factor included, useful bounds will emerge after 5 years of operation of Auger.

¹A somewhat more refined estimate of the cross section would take into account the rotation of the black hole. One may derive a geodesic equation for null geodesics in the equatorial plane of a rotating black hole (since the incoming particles are in the equatorial plane), and calculate the impact parameter $b(M_{\text{BH}}, a)$ at infinite distance from a black hole of mass M_{BH} and rotation parameter a (as defined in Refs. [50,57]). In four dimensions, the extremal value $a = M_{\text{BH}}$ gives $b = 2M_{\text{BH}}$, reproducing the cross section of Eq. (6).

IV. THE COSMOGENIC NEUTRINO FLUX

Among the many possible sources of ultrahigh-energy neutrinos, the cosmogenic flux is the most reliable. This neutrino flux relies only on the assumption that the observed extremely high-energy cosmic rays contain nucleons and are primarily extragalactic in origin. If the charge of the primaries satisfies $Z \lesssim \mathcal{O}(1)$, as recently reported [63], an extragalactic origin is almost guaranteed, as the observed nearly isotropic angular distribution strongly disfavors galactic disk sources [64]. Moreover, even if the absence of the Greisen-Zatsepin-Kuzmin (GZK) cutoff [65] on cosmic ray energies is a reflection of our coincidental position near a nucleus or nucleon-emitting source, one still expects the full cosmogenic neutrino flux.

Briefly, the argument for this is as follows [66]: the known astrophysical environments (within a few Mpc of the Earth) are not among the most powerful, but (in principle) can produce hadronic cosmic rays with the desired energies when parameters are stretched to their limits. Thus if these less powerful sources can accelerate particles above 10^{20} eV, it must be that more powerful distant sources (such as Fanaroff-Riley II radiogalaxies) can accelerate protons above the photopion threshold, giving rise to the cosmogenic neutrino flux. Moreover, the approximately smooth power law behavior of the observed spectrum above 10^{19} eV [12] seems to indicate that any “local source” contribution should be comparable to that of all other sources in the universe. Otherwise, one should invoke an apparently miraculous matching of spectra to account for the smoothness of the spectrum. This smoothness will provide the basis for obtaining the cosmogenic neutrino flux, as discussed in what follows.

The chain reaction generating these cosmogenic neutrinos, triggered by GZK pion photo-production, is well known [67,68]. The resulting neutrino flux depends critically on the cosmological evolution of the cosmic ray sources and on their proton injection spectra [69–71]. The high energy tail of the neutrino spectrum can also receive a significant contribution from semilocal sources, such as the Virgo cluster [72]. Additionally, there is a weak dependence on the details of the cosmological expansion of the universe. For example, a small cosmological constant tends to increase the contribution to neutrino fluxes from higher redshifts [71].

In our analysis we adopt the cosmogenic neutrino flux estimates of Protheroe and Johnson (PJ) [70]. We consider their $\nu_\mu + \bar{\nu}_\mu$ estimate with an injection spectrum with $E_{\text{cutoff}} = 3 \times 10^{21}$ eV. In addition to ν_μ and $\bar{\nu}_\mu$, electron neutrinos also contribute to black hole production. In the high energy peak, the ν_e , ν_μ , and $\bar{\nu}_\mu$ fluxes are nearly identical [71], and we include this ν_e flux in our analysis. The study of PJ incorporates the source cosmological evolution from estimates [73] of the power per comoving volume injected in protons by powerful radio galaxies, taking into account the radio luminosity functions given in Ref. [74]. The shape of the resulting neutrino spectrum is shown in Fig. 2. The flux peaks around $E \approx 2 \times 10^{17}$ eV, which is roughly the same energy suggested by other analyses following a source evolution proportional to $(1+z)^4$ [69,71]. To explore possible

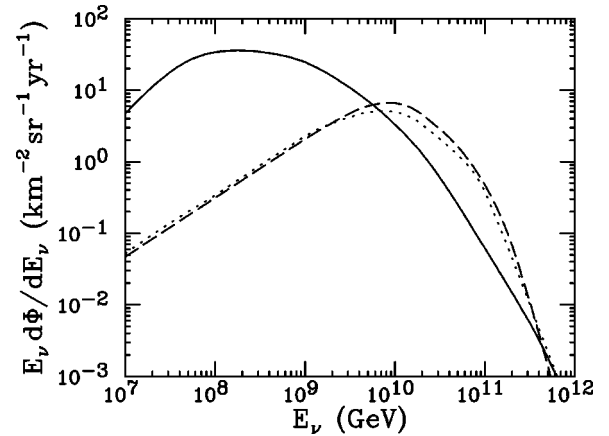


FIG. 2. Cosmogenic $\nu_\mu + \bar{\nu}_\mu + \nu_e$ fluxes from Protheroe and Johnson with an energy cutoff of 3×10^{21} eV (solid) [70], Hill and Schramm (dashed) [72], and a previous estimate by Stecker without source evolution (dotted) [68]. See text for discussion.

additional contributions from semilocal nucleon sources, we also consider below the cosmogenic neutrino flux estimates of Hill and Schramm (HS) [72], which are also given in Fig. 2. A flux estimate of Stecker is also given there. As noted in [11], the PJ, HS, and Stecker fluxes all yield approximately the same rates for BH production.

We stress that the PJ flux agrees with the most recent estimate [71] in the entire energy range, whereas the spectrum obtained in earlier calculations [69] is somewhat narrower, probably as a result of different assumptions regarding the propagation of protons. The PJ analysis is performed within Friedmann cosmology with vanishing cosmological constant Λ , $q_0 = 0.5$, and $H_0 = 75 \text{ km s}^{-1} \text{ Mpc}^{-1}$, assuming an extragalactic magnetic field of 1 nG and a source spectrum proportional to E^{-2} up to redshifts $z = 9$. The extension to cosmological models with $\Lambda \neq 0$ would not produce remarkable changes. For example, for $\Omega_M = 0.3$ and $\Omega_\Lambda = 0.7$, the neutrino flux is increased by a factor of < 1.7 [71].

V. ACCEPTANCE OF SURFACE ARRAYS FOR NEUTRINO SHOWERS

Ultrahigh-energy cosmic neutrinos may be detected by ground arrays and fluorescence detectors on the surface of the Earth, as well as by space-based fluorescence detectors, and neutrino telescopes beneath the Earth’s surface. Here we concentrate on ground arrays, and consider two prominent examples: AGASA, the largest surface array currently in operation, and the Auger Observatory now under construction.

AGASA consists of 111 scintillation detectors each of area 2.2 m^2 , spread over an area of 100 km^2 with 1 km spacing [75]. The array detectors are connected and controlled through a sophisticated optical fiber network. The array also contains a number of shielded scintillation detectors which provide information about the muon content of the showers. The full AGASA experiment has been running since 1992, and has recorded the majority of events claimed to have energies above the GZK cutoff.

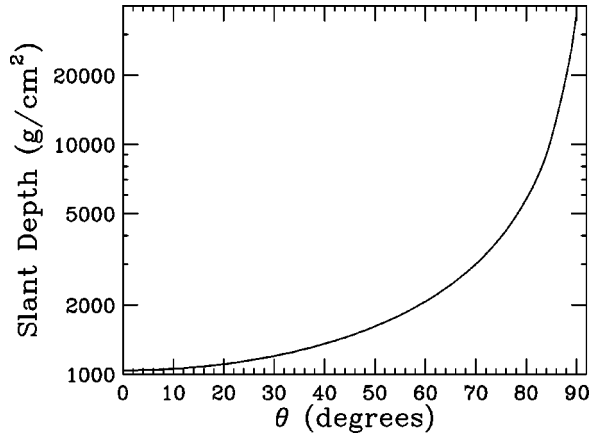


FIG. 3. Slant depths corresponding to various zenith angles θ .

The Auger Observatory is a hybrid experiment, with two sites (one in the northern hemisphere and one in the southern), each covering an area of 3000 km² and consisting of 1600 particle detectors overviewed by 4 fluorescence detectors [76]. The surface array stations are cylindrical water Čerenkov detectors with area 10 m², spaced 1.5 km from each other in a hexagonal grid. Event timing is made possible through global positioning system receivers. The optical system uses the fluorescence technique pioneered by the University of Utah's Fly's Eye detector [77]. "Golden events," events detected by both methods simultaneously, will be extremely valuable for experimental calibration. However, atmospheric fluorescence detection is possible only on clear, dark nights, and so the golden event rate is expected to be less than 10% of the total event rate. We consider only the ground array below. The full southern site is scheduled for completion in 2003. Its engineering array, at 1/40 of the full size, is now complete and is already detecting giant air showers [78].

A surface array's acceptance for neutrino detection may be expressed, in units of km³ water equivalent steradians (km³we sr), as [79]

$$A(E) = S \int_{\theta_{\min}}^{\theta_{\max}} 2\pi \sin \theta d\theta \int_0^{h_{\max}} \frac{\rho_0}{\rho_{\text{water}}} \times e^{-h/H} \mathcal{P}(E, \theta, h) dh, \quad (23)$$

where S is the area of the ground array, $\rho_0 \approx 1.15 \times 10^{-3} \rho_{\text{water}}$ is the density of the atmosphere at ground level, $H \approx 8$ km, $h_{\max} = 15$ km, and $\mathcal{P}(E, \theta, h)$ is the probability of detecting a shower with energy E and zenith angle θ that begins at altitude h . The minimum zenith angle is set by the desire to separate deep neutrino-initiated showers from far showers initiated by hadronic primaries. Typically, a minimum zenith angle in the range $60^\circ < \theta_{\min} < 75^\circ$ is imposed. This range corresponds to atmospheric slant depths of 2000 to 4000 g/cm². (See Fig. 3.) The maximum zenith angle θ_{\max} varies from analysis to analysis. For example, in an analysis

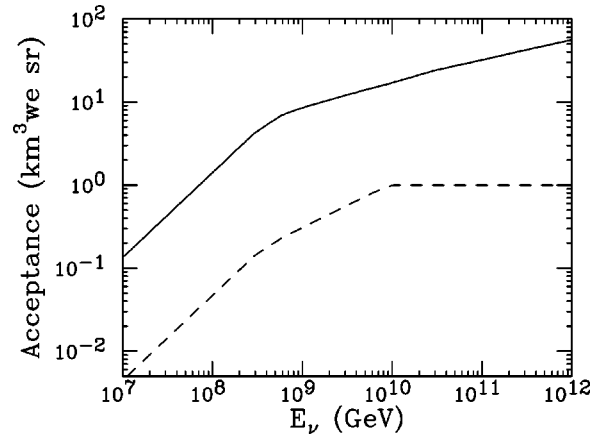


FIG. 4. Ground array acceptances for quasihorizontal air showers at the Auger Observatory (solid) and AGASA (dashed). See text for discussion.

of fully contained showers, a value of θ_{\max} below 90° is required for showers to deposit all of their energy within the array.

Reliable Monte Carlo simulations to determine Auger's acceptance for quasihorizontal showers have been performed by several groups [79,80]. Of course, the acceptance depends on the amount and type of energy generated by neutrino interactions in the atmosphere. For example, the charged current interaction $\nu_\mu p \rightarrow \mu^+ X$ produces a muon that carries approximately 80% of the incoming energy and is not detectable at Auger. Acceptances for both electromagnetic and hadronic showers have been determined in Ref. [80]. BHs decay thermally, according to the number of degrees of freedom available, and so their decays are mainly hadronic. We therefore adopt the hadronic acceptance of Ref. [80] including partially contained showers with zenith angles $\theta > 75^\circ$. The acceptance is given in Fig. 4. Partially contained showers, where the shower axis does not pass through the array, do not contribute significantly to the Auger acceptance for shower energies below 10^{10} GeV.

The AGASA Collaboration has searched for deeply penetrating showers [25,81]. In these studies, they find that, for showers with energy above 10^{10} GeV and the requisite zenith angle, the detection probability \mathcal{P} becomes effectively 100% and independent of altitude [25]. For these energies, then, Eq. (23) may be rewritten as

$$A(E > 10^{10} \text{ GeV}) \approx (S\Omega)_{\text{eff}}(E) \int_0^{h_{\max}} \frac{\rho_0}{\rho_{\text{water}}} e^{-h/H} dh, \quad (24)$$

where

$$(S\Omega)_{\text{eff}}(E) \equiv S \int_{\theta_{\min}}^{\theta_{\max}} 2\pi \sin \theta \mathcal{P}(E, \theta) d\theta \quad (25)$$

is the "effective area \times solid angle" [82]. Acceptances for extremely energetic showers may then be quoted in terms of $(S\Omega)_{\text{eff}}$. The AGASA Collaboration has searched for deeply penetrating showers of any origin in Ref. [81]. They find none with energy above 10^{10} GeV in 9.7×10^7 s of expo-

sure. Given an upper bound of 2.44 events at 90% C.L., then, they derive a flux limit for deeply penetrating showers of $1.9 \times 10^{-10} \text{ km}^{-2} \text{ sr}^{-1} \text{ s}^{-1}$, which implies $(S\Omega)_{\text{eff}} = 132 \text{ km}^2 \text{ sr}$ for quasihorizontal air showers. Equivalently, given Eq. (24), the AGASA acceptance for neutrino initiated events is

$$A(E > 10^{10} \text{ GeV}) \approx 1.0 \text{ km}^3 \text{ we sr}. \quad (26)$$

This acceptance is roughly 30 times smaller than the neutrino acceptance of Auger, as one would naively guess from the ratio between the Auger and AGASA surface areas. For lower energies, since the separation between detectors is smaller at AGASA than at Auger, a conservative approach is to model the AGASA acceptance as that of Auger reduced by a factor of 30. We adopt this estimate for energies below $5 \times 10^8 \text{ GeV}$, and interpolate smoothly between this and Eq. (26) for energies $5 \times 10^8 \text{ GeV} < E_\nu < 10^{10} \text{ GeV}$. The resulting AGASA acceptance is shown in Fig. 4.

VI. NEW BOUNDS FROM AGASA

Given the cross sections, apertures, and fluxes discussed above, the number of BHs detected by a given experiment is

$$N = \int dE_\nu N_A \frac{d\Phi}{dE_\nu} \sigma(E_\nu) A(E_\nu) T, \quad (27)$$

where $A(E_\nu)$ is the experiment's acceptance in $\text{cm}^3 \text{ we sr}$, $N_A = 6.022 \times 10^{23}$ is Avogadro's number, $d\Phi/dE_\nu$ is the source flux of neutrinos, and T is the running time of the detector. We now determine current bounds on BH production from the AGASA experiment. In the next section, we examine future prospects for BH detection at the Auger Observatory.

The AGASA Collaboration has searched for deeply penetrating quasihorizontal showers [25]. The depth at which a shower is initiated is, of course, not directly measurable in a ground array. However, the electromagnetic components of far showers are extinguished by ground level, leaving only a muon component, whereas for deeply penetrating showers, both electromagnetic and muon components are detected. By exploiting this difference, deeply penetrating quasihorizontal showers may be distinguished from showers induced by hadronic cosmic rays.

Relative to showers with muon components only, showers with electromagnetic components have charged particle densities that are more concentrated near the shower axis, and their shower fronts are more curved. The depth at shower maximum X_{max} may then be determined through its correlation to two measurable quantities: η , which parametrizes the lateral distribution of charged particles, and δ , which parametrizes the curvature of the shower front. The values of X_{max} as determined by these correlations, denoted X_{max}^η and X_{max}^δ , are then required to be large to distinguish candidate neutrino events from showers induced by hadronic cosmic rays.

In 1710.5 days of data recorded from December 1995 to November 2000, the AGASA Collaboration found 6 candidate events with $X_{\text{max}}^\eta, X_{\text{max}}^\delta \geq 2500 \text{ g/cm}^2$. The expected

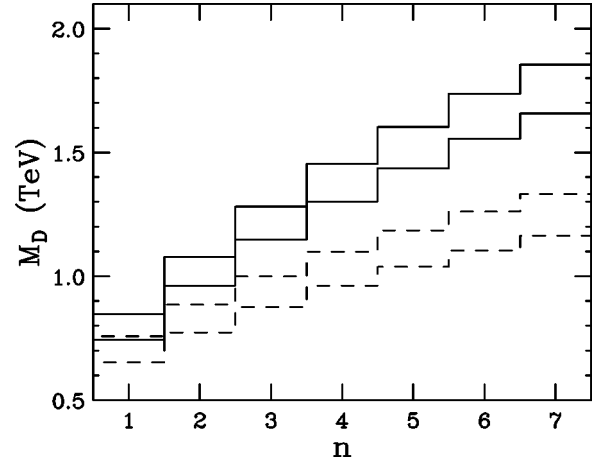


FIG. 5. 95% C.L. lower bound on M_D from nonobservation of quasihorizontal air showers in 1710.5 live days at AGASA for $x_{\text{min}}=1$ (solid) and 3 (dashed), assuming the cosmogenic neutrino flux of Protheroe and Johnson (lower) and Hill and Schramm (upper).

background from hadronic showers is $1.72_{-0.07-0.41}^{+0.14+0.65}$, where the first uncertainty is from Monte Carlo statistics, and the second is systematic. Of the 6 candidate events, however, 5 have values of X_{max}^η and/or X_{max}^δ that barely exceed 2500 g/cm^2 , and are well within ΔX_{max} of this value, where ΔX_{max} is the estimated precision with which X_{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 AGASA results imply lower bounds on the scale of low-scale gravity, assuming the conservative cosmogenic fluxes of Sec. IV. For these fluxes, the expected rate for deeply penetrating showers at AGASA from SM neutrino interactions is about 0.02 events per year, and so negligible. Given 1 event that unambiguously passes all cuts, and the central value of 1.72 background events, the AGASA results imply an upper bound of 3.5 black hole events at 95% C.L. [83].

The 3.5 event contour is given for various dimensions n in Fig. 5. For $x_{\text{min}}=1$, the absence of deeply penetrating showers in the AGASA data implies

$$\begin{aligned} n=4: \quad M_D > 1.3\text{--}1.5 \text{ TeV} \\ n=7: \quad M_D > 1.6\text{--}1.8 \text{ TeV}. \end{aligned} \quad (28)$$

Results for $x_{\text{min}}=3$ are also given in Fig. 5. They imply $M_D > 1.0\text{--}1.1 \text{ TeV}$ for $n=4$, and $M_D > 1.1\text{--}1.3 \text{ TeV}$ for $n=7$; even for $x_{\text{min}}=3$, these bounds exceed or are competitive with all existing collider and astrophysical bounds. As argued in Sec. III A, $x_{\text{min}}=1$ is a reasonable assumption for the present application, as the derivation of limits relies only on the assumption that BHs or their lighter progenitors with mass around M_D decay visibly. This assumption is violated only if their decays are limited to neutrinos, gravitons, and muons.

The range in Eq. (28) is from considering both PJ and HS fluxes. As noted in Ref. [11], the dependence of the bounds

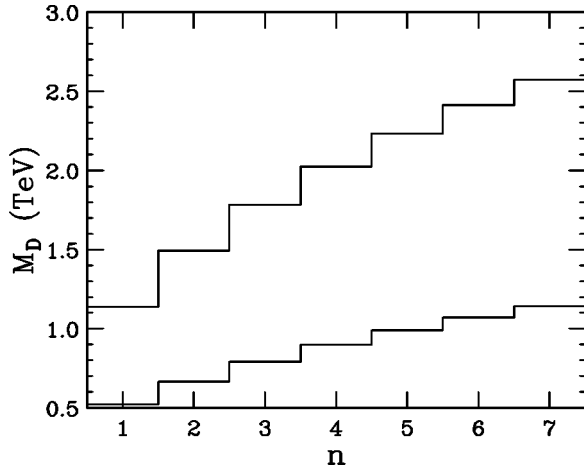


FIG. 6. 95% C.L. upper and lower bounds on M_D for various n , given 6 candidate events above a background of 1.72 in 1710.5 live days at AGASA, and ascribing the excess to BH production. We assume $x_{\min}=1$ and the cosmogenic neutrino flux of Protheroe and Johnson.

on variations in the evaluations of cosmogenic fluxes is weak. These bounds are conservative in that larger noncosmogenic fluxes, as predicted by some models and, as may be indicated by super-GZK cosmic rays, will strengthen them, possibly dramatically. Note also that we have neglected enhancements to cosmic neutrino interactions from sub-Planckian extra-dimensional physics, which are more model-dependent, but can only serve to strengthen these bounds.

The bounds of Eq. (28) are, of course, subject to the $\mathcal{O}(1)$ uncertainties inherent in the parton level cross section. Given this cross section, however, they are direct bounds on the fundamental Planck scale M_D , and are not subject to the uncertainties inherent in collider bounds, such as the choice of brane softening parameter Λ discussed in Sec. II C. Any comparison of collider and cosmic ray bounds is then subject to the independent uncertainties associated with each bound. Nevertheless, for $n \geq 4$, given the geometric BH cross section, the AGASA limit is more stringent than all existing collider bounds for all choices of $\Lambda/M_D \leq 1$.

Before leaving the AGASA results, we derive their implications for extra dimensions if taken at face value. Given 6 events with an expected background of 1.72 events, the expected signal is 0.86 to 11 events at 95% C.L. The preferred region of the (n, M_D) plane is given in Fig. 6. (In Fig. 6, and all following figures, we use the PJ flux. The HS flux yields slightly larger rates.) The evidence for BH production (or any other anomaly) is speculative, given the statistics and the peculiarities of the data noted above. However, this analysis shows the power of cosmic ray measurements for probing extra dimensions. The preferred Planck scales are not probed by any other experiment. At the same time, they will be thoroughly explored in the near future at larger cosmic ray experiments, such as the Auger Observatory, to which we now turn.

VII. FUTURE PROBES AT AUGER

Given the apertures discussed in Sec. V, it is a simple matter to estimate the BH event rate for Auger. The number of detected BH events are given in Fig. 7 for various n as a function of M_D . The Auger ground array is expected to become fully operational in 2003. We assume a running time of 5 years, roughly the data expected before the LHC begins. For $x_{\min}=1$, Auger will probe fundamental Planck scales as large as $M_D=4$ TeV. For $M_D \approx 1$ TeV and $n \geq 4$, 100 BHs could be detected.

Given the prospects for fairly high statistics, detailed BH studies are in principle possible. While BHs with mass near M_D are in some sense of the greatest interest, for detailed studies, one might first restrict attention to more massive BHs (more energetic showers), where the semiclassical description of BHs is expected to be justified. The distribution of BH masses in cosmic ray collisions is given in Fig. 8. They are concentrated near M_D , but the event rate is reduced by only $\mathcal{O}(1)$ factors for x_{\min} as large as 5. This contrasts strongly with the case at colliders, where there is little energy to spare, cross sections are suppressed by two parton distribution functions, and event rates are reduced by two orders of magnitude for $x_{\min}=5$ relative to $x_{\min}=1$ [10]. Total

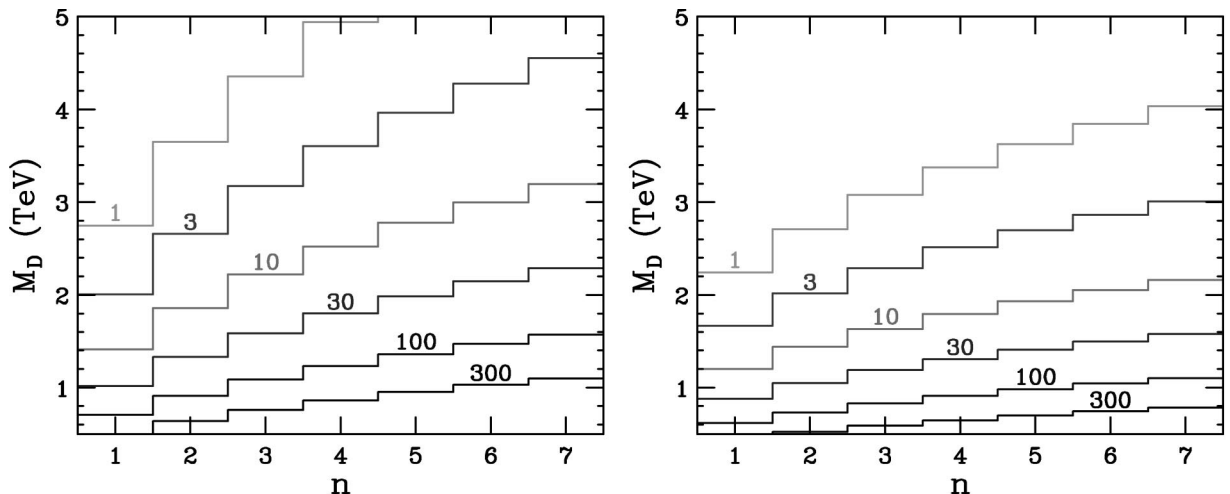


FIG. 7. Event rates in 5 years for the Auger ground array for $x_{\min}=1$ (left) and 3 (right).

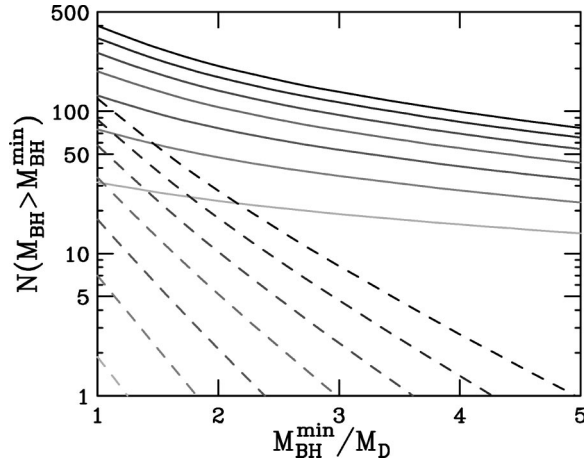


FIG. 8. Event rates for BHs with mass above $M_{\text{BH}}^{\text{min}}$ at the Auger ground array as a function of $M_{\text{BH}}^{\text{min}}/M_D$ for $n=1, \dots, 7$ from below, assuming $M_D=1$ TeV, 5 years running time, and parton cross section πr_s^2 (solid) and $\pi r_s^2 e^{-I_E}$ (dashed).

event rates for $x_{\text{min}}=3$ are also given in Fig. 7. Even for $x_{\text{min}}=3$, we find that ~ 100 BHs may be detected for M_D near 1 TeV.

The dependence of BH event rates on running time T is given in Fig. 9. The event rate contours rise rapidly at first—in even the first few months, Auger will be sensitive to values of M_D beyond present experiments.

If no enhancement of quasihorizontal showers is seen, Auger will set stringent limits on low-scale gravity and scenarios with extra dimensions. To determine these limits, we again assume the cosmogenic fluxes of Sec. IV and that only SM sources of deeply penetrating showers are observed. In contrast to AGASA, SM neutrino interactions lead to observable rates—given the cross section of Fig. 1, 0.5 events per year are expected. In addition, as at AGASA, hadronic show-

ers may fake deeply penetrating showers. As noted above, the Auger aperture of Sec. V assumes zenith angles $\theta > 75^\circ$, corresponding to slant depths of $X_{\text{max}} \geq 4000$ g/cm², significantly more stringent than for the AGASA study [25]. Nevertheless, hadronic showers may be a significant background. We know of no detailed study, but consider the possibility of n_B background events from hadronic showers in 5 years below.

Given these assumptions, the expected background in 5 years is roughly $2+n_B$ events. To determine the expected limit on BH production we assume that $2+n_B$ deeply penetrating events are in fact observed. At 95% C.L., then, the upper bound on signal events for $n_B=0, 5$, and 10 is 4.7, 6.8, and 8.3 events, respectively. In Fig. 9, contours for these event rates are also given. We find that, for $n_B \leq 10$, $x_{\text{min}}=1$, and $n=6$, if no events above background are observed, Auger will extend current bounds on M_D to above 2 TeV after the first year of live time. After 5 years, for $x_{\text{min}}=1$, Auger will set a limit of $M_D \geq 3$ TeV for $n \geq 4$. In conjunction with astrophysical bounds, this will require $M_D \geq 3$ TeV for all n , significantly straining attempts to identify the Planck scale with the weak scale in scenarios of large extra dimensions. Note that we have neglected model-dependent sub-Planckian effects that may increase the rates and strengthen the bounds presented here.

Finally, we consider the impact of the proposed exponential suppression of BH production cross sections. In Fig. 8, we show the dependence on $x_{\text{min}}=M_{\text{BH}}^{\text{min}}/M_D$ for black hole event rates including this suppression. For $x_{\text{min}}=1$, the exponential suppression is not particularly severe, reducing event rates by factors of 3 for large n . Of course, the impact is much larger for larger x_{min} . In Fig. 10, we show the number of BHs observed in time T for parton cross section $\pi r_s^2 e^{-I_E}$. For $x_{\text{min}}=1$, Auger may still see tens of BHs in 5 years, and will extend current bounds to $M_D \approx 2.5$ TeV. For $x_{\text{min}}=3$,

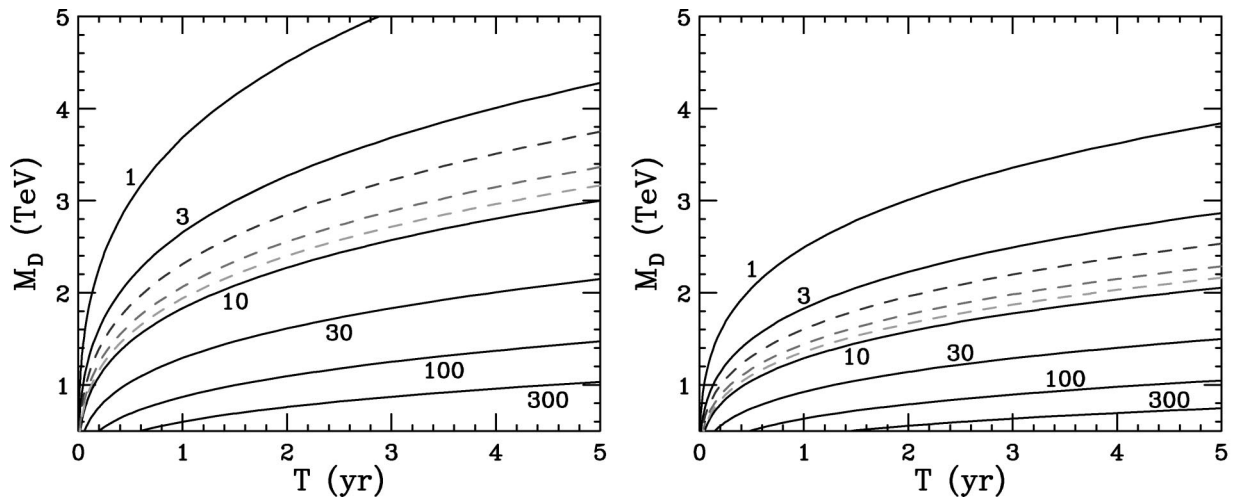


FIG. 9. BH event rates at the Auger ground array for $n=6$ and $x_{\text{min}}=1$ (left) and 3 (right). The dashed contours indicate the expected 95% C.L. lower bound on M_D in the absence of physics beyond the SM and assume $n_B=0, 5$, and 10 background events from hadronic showers (from above). The geometric cross section πr_s^2 is assumed.

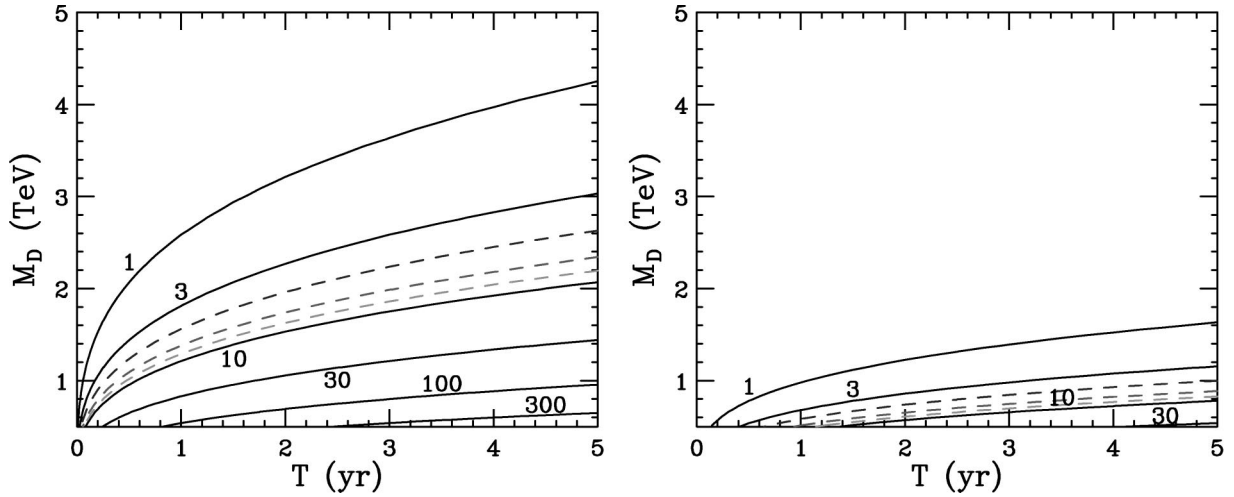


FIG. 10. BH event rates as in Fig. 9, but for parton cross section $\pi r_s^2 e^{-I E}$.

the event rates are quite suppressed, but a few BH events are still observable in 5 years.

VIII. DISTINGUISHING BLACK HOLES FROM SM EVENTS WITH EARTH-SKIMMING NEUTRINOS

If an excess of quasihorizontal showers is observed, how can it be identified as arising from BH events? After all, at first sight, an excess may arise simply from an enhanced flux. With sufficient statistics, a SM explanation may be excluded based on shower properties, as black hole showers differ markedly from those produced by SM charged and neutral current neutrino interactions [13]. It may also be possible to confirm specific predictions of BH production by verifying Hawking radiation through correlations between X_{\max} and shower energy [11].

It is also possible, however, to differentiate BH from SM events by considering additional constraints on ultrahigh-energy neutrino properties. In particular, comparison with Earth-skimming neutrino rates may allow one to distinguish BH and SM interpretations [11]. In this section, we develop this possibility quantitatively, focusing on the question of excluding a SM interpretation for BH events. We also comment briefly on the task of differentiating black hole events from other new physics possibilities at the end of this section.

At ultrahigh energies, even the SM neutrino cross section is large enough that upward-going neutrinos are blocked by the Earth. However, neutrinos that skim the Earth, traveling at low angles along chords with lengths of order their interaction length, are not [29–32]. These Earth-skimming neutrinos may then convert to charged leptons in the Earth’s crust, and the resulting charged leptons may emerge into the atmosphere, producing a signal in cosmic ray detectors. A schematic picture of such an event is given in the top panel of Fig. 11. The best signal is from τ leptons. Unlike electrons that do not escape the Earth’s crust, or muons that do not produce any visible signal in the atmosphere, taus can travel for tens of km in rock, escape, and then decay in the atmosphere, leading to spectacular showers and observable rates

of order 1 per year in both ground arrays [29] and fluorescence detectors [30]. The optimal angle for Earth-skimming neutrinos is energy dependent. For $E_\nu \sim 10^8$ (10^{10}) GeV, the optimal angle relative to the horizon is $\sim 3^\circ$ (1°). Given the angular resolution of cosmic ray detectors, these Earth-skimming events are easily differentiated from standard horizontal neutrino showers.

The scenario changes radically in the presence of a significant cross section for BH production. First, BHs decay largely to hadrons, which do not escape the Earth. Such an event is pictured in the bottom panel of Fig. 11. Of course, BHs also have a significant leptonic branching fraction, but leptons from BH decay carry only a fraction of the initial neutrino energy, and their detection rate is therefore highly suppressed. The probability of detecting BHs produced in the Earth by ground arrays and surface fluorescence detectors is therefore insignificant. Second, a sufficiently large BH cross section also depletes the original neutrino beam through absorption, leading to a substantial suppression of all Earth-skimming events, including those in the top panel of Fig. 11.

To determine the effects of BH production on Earth-skimming rates, we consider here a simple analysis that is nevertheless sufficient to isolate the functional dependence of Earth-skimming rates on cross section parameters. The

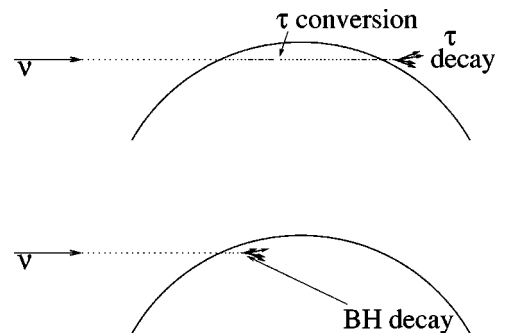


FIG. 11. Top: A neutrino enters the Earth and converts into a charged lepton, which exits the Earth and may be detected. Bottom: A neutrino enters the Earth and produces a BH, which is captured in the Earth.

analysis extends the discussion of Ref. [30], where additional details and discussion may be found.

Consider a flux of neutrinos with energy E_0 . Given the high energies required for detection, the most relevant energies are $E_0 \sim 10^9 - 10^{10}$ GeV, even for cosmogenic flux evaluations peaked at somewhat lower energies, and we may therefore limit the discussion to this rather narrow band of energy. Earth-skimming events occur in the Earth's crust, and so the relevant neutrinos and taus sample only the Earth's surface density, $\rho_s \approx 2.65$ g/cm³. In the SM, the neutrino's path length is

$$L_{CC}^\nu = [N_A \rho_s \sigma_{CC}^\nu]^{-1}, \quad (29)$$

where $N_A \approx 6.022 \times 10^{23}$ g⁻¹ and σ_{CC}^ν is the charged current cross section for $E_\nu = E_0$. (We neglect neutral current interactions, which at these energies serve only to reduce the neutrino energy by approximately 20%.) For $E_0 \sim 10^{10}$ GeV, $L_{CC}^\nu \sim \mathcal{O}(100)$ km. Supplemented by the possibility of BH production, the neutrino's path length is

$$L_{\text{tot}}^\nu = [N_A \rho_s (\sigma_{CC}^\nu + \sigma_{\text{BH}}^\nu)]^{-1}, \quad (30)$$

where σ_{BH}^ν is the BH production cross section for $E_\nu = E_0$.

At these energies, the tau's propagation length is determined not by its decay length but by its energy loss. The τ lepton loses energy in the Earth according to

$$\frac{dE_\tau}{dz} = -(\alpha_\tau + \beta_\tau E_\tau) \rho_s, \quad (31)$$

where, for these energies, α_τ is negligible, and we take $\beta_\tau \approx 0.8 \times 10^{-6}$ cm²/g [84]. The maximal path length for a detectable τ is, then,

$$L^\tau = \frac{1}{\beta_\tau \rho_s} \ln(E_{\text{max}}/E_{\text{min}}), \quad (32)$$

where $E_{\text{max}} \approx E_0$ is the energy at which the tau is created, and E_{min} is the minimal energy at which a τ can be detected. For cosmogenic neutrino fluxes and other reasonable sources, and the acceptances of typical cosmic ray detectors, taus cannot lose much energy and be detected. For $E_{\text{max}}/E_{\text{min}} = 10$, $L^\tau = 11$ km.

Given an isotropic $\nu_\tau + \bar{\nu}_\tau$ flux, the number of taus that emerge from the Earth with sufficient energy to be detected is proportional to an "effective solid angle"

$$\Omega_{\text{eff}} \equiv \int d \cos \theta d \phi \cos \theta P(\theta, \phi), \quad (33)$$

where

$$P(\theta, \phi) = \int_0^\ell \frac{dz}{L_{CC}^\nu} e^{-z/L_{\text{tot}}^\nu} \Theta[z - (\ell - L^\tau)] \quad (34)$$

is the probability for a neutrino with incident nadir angle θ and azimuthal angle ϕ to emerge as a detectable τ . [In Eq. (34), for the reasons noted above, we have neglected the

possibility of detectable signals from BH production by Earth-skimming neutrinos.] Here $\ell = 2R_\oplus \cos \theta$ is the chord length of the intersection of the neutrino's trajectory with the Earth, with $R_\oplus \approx 6371$ km the Earth's radius. Evaluating the integrals, we find [31]

$$\begin{aligned} \Omega_{\text{eff}} &= 2\pi \frac{L_{\text{tot}}^\nu}{L_{CC}^\nu} [e^{L^\tau/L_{\text{tot}}^\nu} - 1] \\ &\times \left[\left(\frac{L_{\text{tot}}^\nu}{2R_\oplus} \right)^2 - \left(\frac{L_{\text{tot}}^\nu}{2R_\oplus} + \left(\frac{L_{\text{tot}}^\nu}{2R_\oplus} \right)^2 \right) e^{-2R_\oplus/L_{\text{tot}}^\nu} \right]. \end{aligned} \quad (35)$$

At the relevant energies, the neutrino interaction length satisfies $L_{\text{tot}}^\nu \ll R_\oplus$. In addition, for $L_{\text{tot}}^\nu \gg L^\tau$, valid when the BH cross section is not very large, Eq. (35) simplifies to

$$\Omega_{\text{eff}} \approx 2\pi \frac{L_{\text{tot}}^\nu L^\tau}{4R_\oplus^2 L_{CC}^\nu}. \quad (36)$$

Equation (36) gives the functional dependence of the Earth-skimming event rate on the BH cross section. This rate is, of course, also proportional to the source neutrino flux Φ^ν at E_0 . Finally, the constant of proportionality is determined by previous studies [29,30], where all the experimental issues entering tau detection have been included. Given these inputs, the number of Earth-skimming neutrino events detected in 5 years is

$$N_{\text{ES}} \approx C_{\text{ES}} \frac{\Phi^\nu}{\Phi_0^\nu} \frac{\sigma_{CC}^{\nu 2}}{(\sigma_{CC}^\nu + \sigma_{\text{BH}}^\nu)^2}, \quad (37)$$

where C_{ES} is the number of Earth-skimming events expected for the standard cosmogenic flux Φ_0^ν in the absence of BH production. For detection by the Auger ground array, $C_{\text{ES}} \approx 3.0$, assuming maximal neutrino mixing and the β_τ value given above [29]. The fluorescence detectors of HiRes provide additional sensitivity [30], as do those of Auger [85]. We conservatively take $C_{\text{ES}} = 3$ for the combined rate in 5 years expected in the SM. Note, however, that the rate may be greatly suppressed for large BH cross sections, as anticipated.

In contrast to Eq. (37), the rate for quasihorizontal showers follows simply from Eq. (27), and has the form

$$N_{\text{QH}} = C_{\text{QH}} \frac{\Phi^\nu}{\Phi_0^\nu} \frac{\sigma_{CC}^\nu + \sigma_{\text{BH}}^\nu}{\sigma_{CC}^\nu}, \quad (38)$$

where $C_{\text{QH}} = 2.5$ for the Auger ground array, as noted previously.

Given a flux Φ^ν and BH cross section σ_{BH}^ν , both N_{ES} and N_{QH} are determined. Event contours are given in the left panel of Fig. 12. As can be seen, given a quasihorizontal event rate N_{QH} , it is impossible to differentiate between an enhancement from large BH cross section and large flux. However, in the region where significant event rates are expected, the N_{QH} and N_{ES} contours are more or less orthogo-

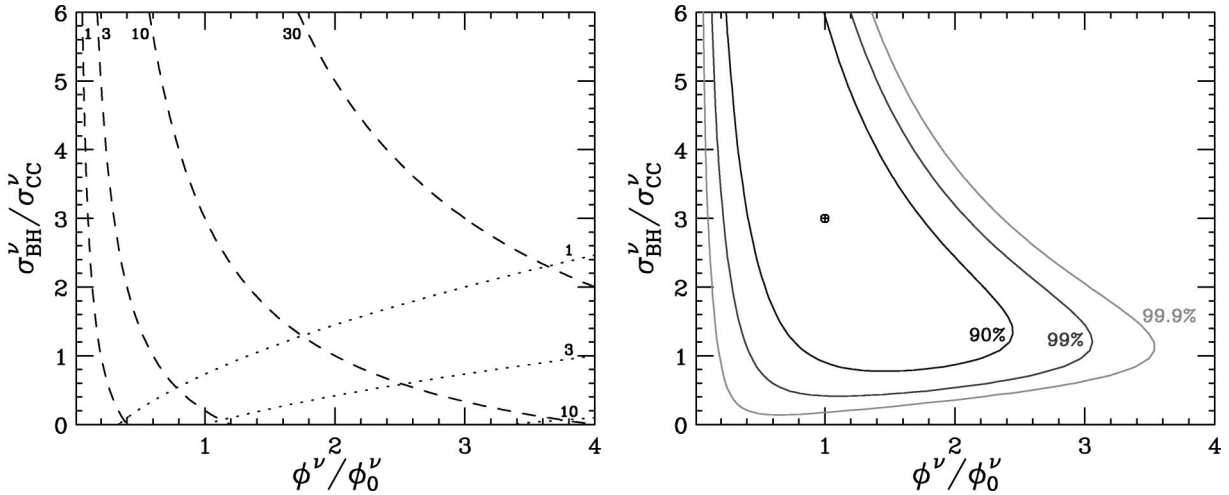


FIG. 12. Left: Contours of constant number of quasihorizontal showers N_{QH} (dashed) and Earth-skimming neutrino events N_{ES} (dotted) as functions of source flux Φ^ν and BH production cross section σ_{BH}^ν . 5 year running times for Auger and HiRes are assumed. Right: Confidence level contours, assuming $\Phi^\nu = \Phi_0^\nu$ and $\sigma_{\text{BH}}^\nu = 3\sigma_{\text{CC}}^\nu$, corresponding to $(N_{\text{QH}}, N_{\text{ES}}) \approx (10, 0.2)$.

nal, and provide complementary information. With measurements of N_{QH} and N_{ES} , both σ_{BH}^ν and Φ^ν may be determined independently, and neutrino interactions beyond the SM may be unambiguously identified. (See also Ref. [31].)

As an example, consider the case in which $\sigma_{\text{BH}}^\nu / \sigma_{\text{CC}}^\nu = 3$, and $\Phi^\nu / \Phi_0^\nu = 1$. On average, one would then observe a total of $N_{\text{QH}} = 10$ deep quasihorizontal showers, an excess of 8 above SM expectations. On average, one also expects $N_{\text{ES}} \approx 0.2$ Earth-skimming events. A SM explanation (with $\sigma_{\text{BH}}^\nu = 0$) of the deeply penetrating event rate would require $\Phi^\nu / \Phi_0^\nu = 4$ and predict 12 Earth-skimming events, a possibility that would be clearly excluded at high confidence level.

More generally, one might try to salvage a SM explanation by attributing the observed rates to statistical fluctuations in both N_{QH} and N_{ES} . Using a maximum likelihood method for Poisson-distributed data [86], we give contours of constant χ^2 in the right panel of Fig. 12. The possibility of a SM interpretation along the $\sigma_{\text{BH}}^\nu = 0$ axis would be excluded at greater than 99.9% C.L. for any assumed flux. The power of the Earth-skimming information is such that the best fit is in fact found for $\Phi^\nu < \Phi_0^\nu$. We find, then, that if even an excess of a handful of quasihorizontal events is observed, by comparing to the Earth-skimming neutrino rate, attempts to explain the excess by SM interactions alone may be excluded. These arguments require only counting experiments, and do not rely on measurements of shower properties.

BH production will most likely be accompanied by more model-independent sub-Planckian effects. In particular, neutral current neutrino cross sections may be enhanced in extra-dimensional scenarios through the exchange of KK gravitons. This will raise the quasihorizontal rate, but will have very little effect on the Earth-skimming event rate, since neutrinos suffer very little energy loss during this process [14]. We expect such effects, then, to further enhance the ratio $N_{\text{QH}} / N_{\text{ES}}$, making a SM explanation even more untenable.

So far, we have not explicitly considered the question of

distinguishing BH events from other types of new physics. However, the prediction of enhanced quasihorizontal event rates *and* diminished Earth-skimming rates is incisive. For example, new physics that increases quasihorizontal rates by enhancing cross sections for $\nu N \rightarrow \ell X$ will also increase Earth-skimming rates. The prediction of suppressed Earth-skimming rates relies on the efficient conversion of neutrino energy directly to hadronic energy, that is, a process with large cross section and large inelasticity. This is a peculiar property of BHs that separates BH production from other possible forms of new physics. The comparison between deep quasi-horizontal shower and Earth-skimming neutrino rates therefore not only effectively excludes a SM interpretation of BH events, but goes a long way toward excluding other new physics explanations.

IX. SUMMARY OF RESULTS AND CONCLUSIONS

In this work we have shown that cosmic ray observations in the recent past (AGASA) and in the near future (Auger) provide extremely sensitive probes of low-scale gravity and extra dimensions. We have focused on the production of TeV-scale BHs resulting from collisions of ultrahigh-energy cosmic neutrinos in the Earth's atmosphere, and have considered the impact of various theoretical issues in the determination of the BH production cross section. In particular, mass shedding, the production of BHs with nonzero angular momentum, and a possible enhancement of the BH cross section can be expected to give minor perturbations. The exponential suppression proposed by Voloshin is more significant, but large and observable BH event rates are still possible.

More specifically, in the case of n extra spatial dimensions compactified on an n -torus with a common radius, we have found the following:

Present bounds on atmospheric BH production imply 95% C.L. lower limits on the fundamental Planck mass of $M_D \geq 1.3\text{--}1.5$ TeV for $n=4$, rising to $M_D \geq 1.6\text{--}1.8$ TeV for

$n=7$. These bounds follow from the nonobservation of a significant excess of deep, quasihorizontal showers in 1710.5 days of running recently reported by the AGASA Collaboration [25].

The absence of a deeply penetrating signal in the Fly's Eye data [26] also implies lower bounds on M_D . These are consistently weaker, however. For example, for $n=6$, $x_{\min}=1$, and the same (PJ) flux we have used, Ringwald and Tu find $M_D > 900$ GeV [16]. We find this difference to be significant: the AGASA and Fly's Eye constraints rely on identical theoretical assumptions, and given the scaling in Eq. (10), a factor of 2 difference in M_D bounds corresponds to a factor of more than 4 in acceptance or, equivalently, running time.

The AGASA limits derived here exceed the $D\emptyset$ bound $M_D \gtrsim 0.6\text{--}1.2$ TeV, where the variation reflects uncertainty from the choice of ultraviolet cutoff for graviton momenta transverse to the brane. The cosmic ray limits are subject to a separate set of uncertainties, discussed at length above, but follow from conservative evaluations of the neutrino flux and experimental aperture, and $x_{\min}=1$. For $x_{\min}=3$, these limits are somewhat reduced, but still generally exceed the Tevatron bounds.

The cosmic ray bounds from AGASA therefore represent the best existing limits on the scale of TeV gravity for $n \geq 4$ extra spatial dimensions. A summary of the most stringent present bounds on M_D for $n \geq 2$ extra dimensions is given in Fig. 13.

The reach of AGASA will be extended significantly by the Auger Observatory. If no quasihorizontal extended air shower events are observed in 5 years (beyond the expected two SM neutrino events supplemented by as many as 10 hadronic background events), Auger will set a limit of $M_D \gtrsim 3$ TeV, at 95% C.L., for $n \geq 4$. Even in the case where the cross section is decreased by the exponential suppression factor in Eq. (21), a bound $M_D \gtrsim 2$ TeV may be found under the same background assumptions.

Conversely, given the large reach of Auger, tens of BH events may be observed per year. We have discussed in some detail how combined measurements of quasihorizontal air showers and Earth-skimming $\nu_\tau \rightarrow \tau$ events may be used to identify new neutrino interactions beyond the SM, even with complete uncertainty about the incident neutrino flux. In the case of BH production, the quasihorizontal event rate is enhanced, while the Earth-skimming rate is suppressed, since BH production in the Earth acts as an absorptive channel, depleting the SM rate. With counting experiments alone, one can therefore exclude a SM interpretation of BH events, and

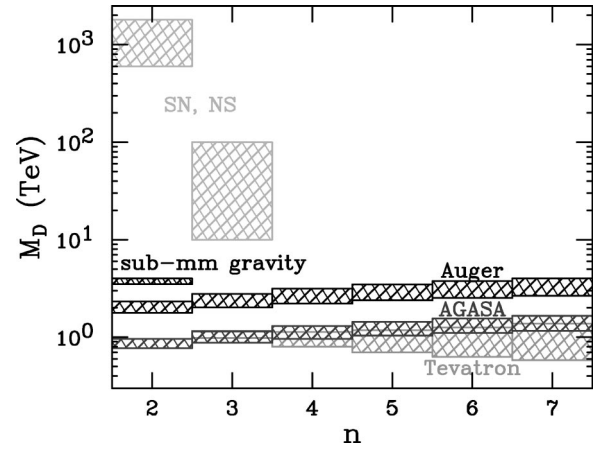


FIG. 13. Bounds on the fundamental Planck scale M_D from tests of Newton's law on sub-millimeter scales, bounds on supernova cooling and neutron star heating, dielectron and diphoton production at the Tevatron, and nonobservation of BH production at AGASA. Future limits from the Auger ground array, assuming 5 years of data and no excess above the SM neutrino background, are also shown. The range in Tevatron bounds corresponds to the range of brane softening parameter $\Lambda/M_D=0.5\text{--}1$. The range in cosmic ray bounds is for $x_{\min}=1\text{--}3$. See text for discussion.

may distinguish BH events from almost all other possible forms of new physics.

In conclusion, in the next several years prior to the analysis of data from the LHC, super-Planckian BH production from cosmic rays provides a promising probe of extra dimensions. Searches for BH-initiated quasihorizontal showers in the Earth's atmosphere at AGASA provide the most stringent bounds on low-scale gravity at present, and the Auger Observatory will extend this sensitivity to fundamental Planck scales well above the TeV scale.

ACKNOWLEDGMENTS

We thank R. Emparan, S. Giddings, A. Ringwald, and G. Sigl for useful communications. J.L.F. thanks A. Guth and A. Vilenkin for discussions regarding Kerr BHs. A.D.S. thanks D. Eardley for a discussion about BH collisions. The work of L.A.A. and H.G. has been partially supported by the US National Science Foundation (NSF), under Grants No. PHY-9972170 and No. PHY-0073034, respectively. The work of J.L.F. was supported in part by the Department of Energy (DOE) under cooperative research agreement DF-FC02-94ER40818. The work of A.D.S. is supported in part by DOE Grant No. DE-FG01-00ER45832 and NSF Grant No. PHY-0071312.

- [1] D. Amati, M. Ciafaloni, and G. Veneziano, Phys. Lett. B **197**, 81 (1987); **289**, 87 (1992).
 [2] G. 't Hooft, Phys. Lett. B **198**, 61 (1987); Commun. Math. Phys. **117**, 685 (1988).
 [3] S.W. Hawking, Commun. Math. Phys. **43**, 199 (1975).
 [4] I. Antoniadis, Phys. Lett. B **246**, 377 (1990); N. Arkani-

- Hamed, S. Dimopoulos, and G.R. Dvali, *ibid.* **429**, 263 (1998).
 [5] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999).
 [6] T. Banks and W. Fischler, hep-th/9906038.
 [7] R. Emparan, G.T. Horowitz, and R.C. Myers, Phys. Rev. Lett. **85**, 499 (2000).
 [8] S.B. Giddings and E. Katz, J. Math. Phys. **42**, 3082 (2001).

- [9] S.B. Giddings and S. Thomas, *Phys. Rev. D* **65**, 056010 (2002).
- [10] S. Dimopoulos and G. Landsberg, *Phys. Rev. Lett.* **87**, 161602 (2001).
- [11] J.L. Feng and A.D. Shapere, *Phys. Rev. Lett.* **88**, 021303 (2001).
- [12] M. Nagano and A.A. Watson, *Rev. Mod. Phys.* **72**, 689 (2000).
- [13] L. Anchordoqui and H. Goldberg, *Phys. Rev. D* **65**, 047502 (2002).
- [14] R. Emparan, M. Masip, and R. Rattazzi, *Phys. Rev. D* **65**, 064023 (2002).
- [15] Y. Uehara, hep-ph/0110382.
- [16] A. Ringwald and H. Tu, *Phys. Lett. B* **525**, 135 (2002).
- [17] S. Nussinov and R. Shrock, *Phys. Rev. D* **59**, 105002 (1999); G. Domokos and S. Kovesi-Domokos, *Phys. Rev. Lett.* **82**, 1366 (1999).
- [18] P. Jain, D.W. McKay, S. Panda, and J.P. Ralston, *Phys. Lett. B* **484**, 267 (2000).
- [19] C. Tyler, A.V. Olinto, and G. Sigl, *Phys. Rev. D* **63**, 055001 (2001); M. Kachelriess and M. Plumacher, *ibid.* **62**, 103006 (2000); L. Anchordoqui, H. Goldberg, T. McCauley, T. Paul, S. Reucroft, and J. Swain, *ibid.* **63**, 124009 (2001).
- [20] J. Alvarez-Muniz, F. Halzen, T. Han, and D. Hooper, *Phys. Rev. Lett.* **88**, 021301 (2002).
- [21] For a review, see G. Sigl, hep-ph/0109202.
- [22] L.A. Anchordoqui, J.L. Feng, H. Goldberg, and A.D. Shapere, Auger Internal Note, GAP-2001-053.
- [23] M.B. Voloshin, *Phys. Lett. B* **518**, 137 (2001).
- [24] M.B. Voloshin, *Phys. Lett. B* **524**, 376 (2002).
- [25] AGASA Collaboration, S. Yoshida *et al.*, in Proceedings of the 27th International Cosmic Ray Conference, Hamburg, Germany, 2001, Vol. 3, p. 1142.
- [26] R.M. Baltrusaitis *et al.*, *Phys. Rev. D* **31**, 2192 (1985).
- [27] For a summary of LEP bounds, see C. Pagliarone, hep-ex/0111063.
- [28] D0 Collaboration, B. Abbott *et al.*, *Phys. Rev. Lett.* **86**, 1156 (2001).
- [29] X. Bertou, P. Billoir, O. Deligny, C. Lachaud, and A. Letessier-Selvon, astro-ph/0104452.
- [30] J.L. Feng, P. Fisher, F. Wilczek, and T.M. Yu, *Phys. Rev. Lett.* **88**, 161102 (2002).
- [31] A. Kusenko and T. Weiler, *Phys. Rev. Lett.* **88**, 161101 (2002).
- [32] G. Domokos and S. Kovesi-Domokos, hep-ph/9801362; hep-ph/9805221; D. Fargion, astro-ph/0002453; astro-ph/0101565.
- [33] K. Cheung, hep-ph/0110163.
- [34] S.C. Park and H.S. Song, hep-ph/0111069.
- [35] T.G. Rizzo, hep-ph/0111230.
- [36] S. Dimopoulos and R. Emparan, hep-ph/0108060.
- [37] S. Hossenfelder, S. Hofmann, M. Bleicher, and H. Stocker, hep-ph/0109085; R. Casadio and B. Harms, hep-th/0110255; S. Hofmann, M. Bleicher, L. Gerland, S. Hossenfelder, S. Schwabe, and H. Stocker, hep-ph/0111052; R. Casadio, A. Fabbri, and L. Mazzacurati, gr-qc/0111072; G. Pasztor and T.G. Rizzo, hep-ph/0112054; G.F. Giudice, R. Rattazzi, and J.D. Wells, hep-ph/0112161; M. Bleicher, S. Hofmann, S. Hossenfelder, and H. Stocker, hep-ph/0112186.
- [38] G.F. Giudice, R. Rattazzi, and J.D. Wells, *Nucl. Phys.* **B544**, 3 (1999).
- [39] J. Lykken and S. Nandi, *Phys. Lett. B* **485**, 224 (2000).
- [40] E.J. Ahn, M. Cavaglia, and A.V. Olinto, hep-th/0201042.
- [41] C.D. Hoyle, U. Schmidt, B.R. Heckel, E.G. Adelberger, J.H. Gundlach, D.J. Kapner, and H.E. Swanson, *Phys. Rev. Lett.* **86**, 1418 (2001).
- [42] T. Kaluza, *Sitzungsber. K. Preuss. Akad. Wiss.* **K1**, 966 (1921); O. Klein, *Z. Phys.* **37**, 895 (1926); *Surv. High Energy Phys.* **5**, 241 (1926); *Nature (London)* **118**, 516 (1926).
- [43] S. Cullen and M. Perelstein, *Phys. Rev. Lett.* **83**, 268 (1999); V. Barger, T. Han, C. Kao, and R.J. Zhang, *Phys. Lett. B* **461**, 34 (1999); L.J. Hall and D.R. Smith, *Phys. Rev. D* **60**, 085008 (1999); C. Hanhart, J.A. Pons, D.R. Phillips, and S. Reddy, *Phys. Lett. B* **509**, 1 (2001); S. Hannestad and G. Raffelt, *Phys. Rev. Lett.* **87**, 051301 (2001); hep-ph/0110067.
- [44] L3 Collaboration, M. Acciarri *et al.*, *Phys. Lett. B* **464**, 135 (1999); L3 Collaboration, M. Acciarri *et al.*, *ibid.* **470**, 281 (1999); see also M. Gataullin, hep-ex/0108008.
- [45] T. Han, J.D. Lykken, and R.J. Zhang, *Phys. Rev. D* **59**, 105006 (1999).
- [46] J.L. Hewett, *Phys. Rev. Lett.* **82**, 4765 (1999).
- [47] M. Bando, T. Kugo, T. Noguchi, and K. Yoshioka, *Phys. Rev. Lett.* **83**, 3601 (1999).
- [48] M. Bando and T. Noguchi, hep-ph/0011374.
- [49] H. Murayama and J.D. Wells, hep-ph/0109004.
- [50] R.C. Myers and M.J. Perry, *Ann. Phys. (N.Y.)* **172**, 304 (1986).
- [51] P.C. Argyres, S. Dimopoulos, and J. March-Russell, *Phys. Lett. B* **441**, 96 (1998).
- [52] CTEQ Collaboration, H.L. Lai *et al.*, *Eur. Phys. J. C* **12**, 375 (2000).
- [53] J. Preskill, P. Schwarz, A.D. Shapere, S. Trivedi, and F. Wilczek, *Mod. Phys. Lett. A* **6**, 2353 (1991).
- [54] P.D. D'Eath and P.N. Payne, *Phys. Rev. D* **46**, 658 (1992); **46**, 675 (1992); **46**, 694 (1992).
- [55] D.M. Eardley and S.B. Giddings, gr-qc/0201034.
- [56] R. Emparan (private communication).
- [57] R.M. Wald, *General Relativity* (University of Chicago Press, Chicago, 1984), p. 491.
- [58] S.N. Solodukhin, hep-ph/0201248.
- [59] J.M. Cornwall, *Phys. Lett. B* **243**, 271 (1990).
- [60] H. Goldberg, *Phys. Lett. B* **246**, 445 (1990).
- [61] Y.Y. Charng and R.S. Willey, *Phys. Rev. D* **65**, 105018 (2002).
- [62] S.B. Giddings, in Proceedings of the APS/DPF/DPB Summer Study on the Future of Particle Physics (Snowmass 2001), edited by R. Davidson and C. Quigg, hep-ph/0110127.
- [63] M. Ave *et al.*, astro-ph/0112071.
- [64] M. Takeda *et al.*, astro-ph/9902239; N. Hayashida *et al.*, *Astrophys. J.* **522**, 225 (1999); M. Hillas, *Nature (London)* **395**, 15 (1998).
- [65] K. Greisen, *Phys. Rev. Lett.* **16**, 748 (1966); G.T. Zatsepin and V.A. Kuzmin, *Zh. Éksp. Teor. Fiz., Pis'ma Red.* **4**, 114 (1966) [*JETP Lett.* **4**, 78 (1966)].
- [66] E.J. Ahn, G. Medina-Tanco, P.L. Biermann, and T. Stanev, astro-ph/9911123; G.R. Farrar and T. Piran, astro-ph/0010370; L.A. Anchordoqui, H. Goldberg, and T.J. Weiler, *Phys. Rev. Lett.* **87**, 081101 (2001); L. Anchordoqui, H. Goldberg, S. Reucroft, and J. Swain, *Phys. Rev. D* **64**, 123004 (2001).
- [67] V.S. Beresinsky and G.T. Zatsepin, *Phys. Lett. B* **28**, 423 (1969).
- [68] F.W. Stecker, *Astrophys. J.* **228**, 919 (1979).

- [69] S. Yoshida and M. Teshima, *Prog. Theor. Phys.* **89**, 833 (1993).
- [70] R.J. Protheroe and P.A. Johnson, *Astropart. Phys.* **4**, 253 (1996); R.J. Protheroe, *Nucl. Phys. B (Proc. Suppl.)* **77**, 465 (1999).
- [71] R. Engel, D. Seckel, and T. Stanev, *Phys. Rev. D* **64**, 093010 (2001).
- [72] C.T. Hill and D.N. Schramm, *Phys. Rev. D* **31**, 564 (1985).
- [73] J.P. Rachen and P.L. Biermann, *Astron. Astrophys.* **272**, 161 (1993).
- [74] J.A. Peacock, *Mon. Not. R. Astron. Soc.* **217**, 601 (1985).
- [75] N. Chiba *et al.*, *Nucl. Instrum. Methods Phys. Res. A* **311**, 338 (1992).
- [76] <http://www.auger.org/admin/>
- [77] R.M. Baltrusaitis *et al.*, *Nucl. Instrum. Methods Phys. Res. A* **240**, 410 (1985).
- [78] AUGER Collaboration, I. Allekote *et al.*, in Proceedings of the 27th International Cosmic Ray Conference, Hamburg, Germany, 2001, p. 370.
- [79] P. Billoir, in Venice 1999, Neutrino telescopes, Vol. 2, p. 111.
- [80] K.S. Capelle, J.W. Cronin, G. Parente, and E. Zas, *Astropart. Phys.* **8**, 321 (1998).
- [81] AGASA Collaboration, N. Inoue, in Proceedings of the 26th International Cosmic Ray Conference (ICRC 99), edited by D. Kieda, M. Salamon, and B. Dingus, Salt Lake City, Utah, 1999, Vol. 1, p. 361.
- [82] For Monte Carlo studies of the closely related “effective area” for AGASA, see Ref. [75] and AGASA Collaboration, H. Ohoka, S. Yoshida, and M. Takeda, *Nucl. Instrum. Methods Phys. Res. A* **385**, 268 (1997).
- [83] G.J. Feldman and R.D. Cousins, *Phys. Rev. D* **57**, 3873 (1998).
- [84] S.I. Dutta, M.H. Reno, I. Sarcevic, and D. Seckel, hep-ph/0012350.
- [85] AUGER Collaboration, R. Cester *et al.*, in Proceedings of the 27th International Cosmic Ray Conference, Hamburg, Germany, 2001, p. 711.
- [86] F. James and R. Cousins, *Eur. Phys. J. C* **15**, 195 (2000).