

# Quest for Black Holes and Superstring Excitations in Cosmic Ray Data

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## Abstract

In this talk we discuss aspects of TeV-scale gravitational collapse to black holes and string balls and their subsequent evaporation. Special emphasis is placed on the interplay of the string  $\rightleftharpoons$  black hole correspondence principle. These ideas are then explored in the context of cosmic ray physics. First, the potential for observing showers mediated by black holes or superstring excitations is examined. Next, existing data from neutrino telescopes are used to constrain the parameter space for the unseen dimensions of the universe. Finally, we close with a discussion of future prospects.

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## I. INTRODUCTION

A promising route towards reconciling the apparent mismatch of the fundamental scales of particle physics and gravity is to modify the short distance behavior of gravity at scales much larger than the Planck length. This can be accomplished in a straightforward manner [1, 2] if one assumes that the standard model (SM) fields are confined to a 4-dimensional world (corresponding to our apparent universe), while gravity lives in a higher dimensional space. One virtue of this assumption is that very large extra dimensions are allowed without conflicting with current experimental bounds [3], leading to a fundamental Planck mass much lower than its effective 4-dimensional value, even as low as the electroweak scale. In particular, if the spacetime is taken as a direct product of a 4-dimensional spacetime and a flat spatial torus  $T^{D-4}$  (of common linear size  $2\pi r_c$ ), one obtains a definite representation of this picture in which the effective 4-dimensional Planck scale,  $M_{\text{Pl}} \sim 10^{19}$  GeV, is related to the fundamental scale of gravity,  $M_D$ , according to  $M_{\text{Pl}}^2 = 8\pi M_D^{D-2} r_c^{D-4}$  [1].

One of the most startling predictions of scenarios of this sort is that microscopic black holes (BHs) could be observed in forthcoming colliders [4] and at existing cosmic ray detectors [5]. In fact, the non-observation of cosmic neutrinos has already been used to set lower bounds on the fundamental scale of  $D$ -dimensional gravity [6]. This work relied on classical models of BH production, which are valid only for sufficiently massive BHs,  $M_{\text{BH}} \gg M_D$ . For masses close to  $M_D$ , gravity becomes strong and the classical description can no longer be trusted. 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 [7]. In principle embedding TeV-scale gravity models in realistic string models might facilitate the calculation of cross sections for BHs (and string excitations) of masses close to  $M_D$ .

Our purpose here is to explore this possibility within a simple string theory. We begin by reviewing the standard classical and semiclassical pictures of BH production and decay, and the string  $\Leftrightarrow$  BH correspondence principle. We then calculate the  $\nu N$  cross section for production of excited string states within superstring theory, and find it to be larger than the BH production cross section [8]. This makes it plausible that the BH cross section can be used as a lower bound down to  $M_{\text{BH}} = M_D$ , which would significantly strengthen existing lower bounds on  $M_D$ .

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 into the internal dimensions of the universe [2]. 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,  $D = 10$ . Generalization to arbitrary number of dimensions is straightforward.

## II. BLACK HOLES AND STRING BALLS WITH LOW-SCALE GRAVITY

Analytic and numerical studies have revealed that gravitational collapse takes place at sufficiently high energies and small impact parameters, as conjectured years ago by Thorne [9]. In the case of 4-dimensional head-on collisions [10], as well as those with non-zero impact parameter [11], a horizon forms when and only when a mass is compacted into a hoop whose circumference in every direction is less than  $2\pi$  times its Schwarzschild radius up to a factor

of order 1. In the 10-dimensional scenario the Schwarzschild radius still characterizes the maximum impact parameter for horizon formation [12]. In the course of collapse, a certain amount of energy is radiated in gravitational waves by the multipole moments of the incoming shock waves [10], leaving a fraction  $y \equiv M_{\text{BH}}/\sqrt{\hat{s}}$  available to Hawking evaporate [13]. 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 as partons. This ratio depends on the impact parameter of the collision, as well as on the dimensionality of space-time [14].

Subsequent to formation, the BH proceeds to decay [15]. 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 [16]

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

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

$$r_s(M_{\text{BH}}) = \frac{1}{M_{10}} \left[ \frac{M_{\text{BH}}}{M_{10}} 8\pi^{3/2} \Gamma(9/2) \right]^{1/7} \quad (2)$$

is the Schwarzschild radius [17],

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

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  $\sim T_{\text{BH}} \sim 1/r_s$ , 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 [18]. In the low energy limit,  $Q r_s \ll 1$ , higher order terms get suppressed by a factor of  $3(Q r_s)^{-2}$  for fermions and  $25(Q r_s)^{-2}$  for gauge bosons. For an average particle energy,  $\langle Q \rangle \approx r_s^{-1}$ , higher partial waves also get suppressed, though the suppression is not as large. This strongly suggests that the BH is only sensitive to the radial coordinate and does not make use of the extra angular modes available in the internal space [19]. A recent numerical study [20] 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 [19], the bulk emission of gravitons would need to show 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 [21] 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 total number of particles emitted is roughly equal to the BH entropy,

$$S_{\text{BH}} = \frac{\pi}{2} M_{\text{BH}} r_s. \quad (4)$$

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

$$\frac{d\dot{M}_{\text{BH}}}{dQ} = - \sum_i c_i \frac{\sigma_s(Q, r_s)}{8\pi^2} Q^3 \left[ \exp\left(\frac{Q}{T_{\text{BH}}}\right) - (-1)^{2s} \right]^{-1}, \quad (5)$$

where  $c_i$  is the number of internal degrees of freedom of particle species  $i$ . Integration of Eq. (5) leads to

$$\dot{M}_{\text{BH}} = - \sum_i c_i f \frac{\Gamma_s}{8 \pi^2} \Gamma(4) \zeta(4) T_{\text{BH}}^4 A_4, \quad (6)$$

where  $f = 1$  ( $f = 7/8$ ) 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 [19]

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

seen by the SM fields ( $\Gamma_{s=1/2} \approx 0.33$  and  $\Gamma_{s=1} \approx 0.34$  [22]). Now, since the ratio of degrees of freedom for gauge bosons, quarks and leptons is 29:72:18 (the Higgs boson is not included), from Eq. (6) one obtains a rough estimate of the mean lifetime,

$$\tau_{\text{BH}} \approx 1.67 \times 10^{-27} \left(\frac{M_{\text{BH}}}{M_{10}}\right)^{9/7} \left(\frac{\text{TeV}}{M_{10}}\right) \text{ s}, \quad (8)$$

which indicates that BHs evaporate instantaneously into visible quanta. This is within a factor of 2 of the heuristic estimate made in Ref. [23],

$$\tau_{\text{BH}} \sim 6.58 \times 10^{-28} \left(\frac{M_{\text{BH}}}{M_{10}}\right)^{9/7} \left(\frac{\text{TeV}}{M_{10}}\right) \text{ s}, \quad (9)$$

which is generally used in the literature.

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

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

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 be safely neglected [24]. For BH with initial masses well above  $M_{10}$ , most of the decay process can be well described within the semiclassical approximation. However, the condition stated in Eq. (10) 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  $\rho(M) = e^{\beta_H M}/M$ , because the partition function,

$$Z(\beta) = \int_0^\infty dM \rho(M) e^{-M\beta}, \quad (11)$$

would fail to converge [25]. Indeed, the partition function converges only if the temperature is less than the Hagedorn temperature,  $\beta_H^{-1}$ , which is expected to be  $\sim M_s$ . As  $\beta$  decreases to the transition point  $\beta_H$ , 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 [26]. In the limit, the total probability diverges, indicating that the canonical ensemble is inadequate for the treatment of the systems. However, one can still employ a microcanonical ensemble of a large number of similar insulated system each with a given fixed energy  $E$ . With the center-of-mass at rest,  $E = M$ . This means that the density of

states is just  $\rho(M)$  and the entropy  $S = \ln \rho(M)$ . In this picture the equilibrium among systems is governed by the equality of the temperatures, defined for each system as

$$T \equiv \left( \frac{\partial S}{\partial M} \right)^{-1} = \frac{M}{\beta_H M - 1}. \quad (12)$$

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,

$$C \equiv -\frac{1}{T^2} \left( \frac{\partial^2 S}{\partial M^2} \right)^{-1} = -\left( \frac{M}{T} \right)^2, \quad (13)$$

is negative, as is the case for BHs [27]. 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 [28]. 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 a minimal radius for the Schwarzschild radius of any BH [29]; (ii)  $T_{\text{BH}} \sim \beta_H^{-1}$  for  $r_s \sim M_s^{-1}$  [30]. (iii) There is a seeming correlation between the greybody factors in BH decay and the level structure of excited strings [31]. The string  $\rightleftharpoons$  BH “correspondence principle” [32] 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 [33]. (The probability that a BH will radiate a large string, or else that a large string would undergo fluctuation to become a BH is very small [34].)

The continuity of the cross section at the correspondence point, at least parametrically in energy and string coupling, provides independent supportive argument for this picture [7]. 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  $g_s^2 \hat{s} / M_s^2 \sim 1$  [35], 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 [7]

$$\hat{\sigma}(\sqrt{\hat{s}}) \sim \begin{cases} \frac{g_s^2 \hat{s}}{M_s^4} & M_s \ll \sqrt{\hat{s}} \leq M_s / g_s, \\ \frac{1}{M_s^2} & M_s / g_s < \sqrt{\hat{s}} \leq M_s / g_s^2, \\ \frac{1}{M_{10}^2} \left[ \frac{\sqrt{\hat{s}}}{M_{10}} \right]^{2/7} & M_s / g_s^2 < \sqrt{\hat{s}}, \end{cases} \quad (14)$$

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

Colliders have not yet attained the energies required to probe TeV scale gravitational collapse; nevertheless cosmic rays may already provide some clues. To date a handful of cosmic rays have been observed with energy in excess of  $10^{11}$  GeV [36]. When these particles impinge on a stationary nucleon in the upper atmosphere they probe center-of-mass

energies as high as  $\sqrt{s} \sim 400$  TeV. In what follows we estimate the sensitivity of cosmic ray observatories to SB/BH production.

### III. PROBES OF LARGE EXTRA DIMENSIONS WITH COSMIC RAYS

The SB/BH production cross section,  $\mathcal{O}(M_{\text{EW}}^{-1})$ , is about 5 orders of magnitude smaller than QCD cross sections,  $\mathcal{O}(\Lambda_{\text{QCD}}^{-1})$ , thus making it futile to hunt for SB/BHs in hadronic cosmic rays. On the other hand, SM neutrinos interact so weakly that any quantum gravitational enhancement of the cross section can be experimentally distinguished from background. A substantial neutrino component may accompany the observed cosmic ray flux because of neutrino production in cosmic beam dumps (such as gamma ray burst fireballs [37], X-ray binaries [38], or active galactic nuclei [39]). In addition, the inelastic collision of ultra-high energy ( $> 10^{10.6}$  GeV) nucleons with the relic photons permeating the universe should deplete the hadronic cosmic ray intensity at the end of the spectrum and breed neutrinos from decay products of charged pions [40]. The reaction generating these cosmogenic neutrinos is well known physics and the flux depends only on the existence of ultra-high energy nucleon sources more distant than  $\approx 8$  Mpc. In this case, the expected  $\nu$ -flux can be accurately predicted by fitting the observed spectrum with a homogeneous population of nucleon sources [41]. The existence of the cosmogenic neutrinos is a safe bet, implying a minimum guaranteed flux for testing any new physics that promotes the neutrino cross section to a sub-hadronic size.<sup>1</sup>

The inclusive production of BHs proceeds through different final states for different classical impact parameters  $b$  [14]. These final states are characterized by the fraction  $y(z)$  of the initial parton center-of-mass energy,  $\sqrt{\hat{s}} = \sqrt{x s}$ , which is trapped within the horizon. Here,  $z = b/b_{\text{max}}$ , and  $b_{\text{max}} = 1.3 r_s(\sqrt{\hat{s}})$  [14]. With a lower cutoff  $M_{\text{BH,min}}$  on the BH mass required for the validity of the semi-classical description, this implies a joint constraint

$$y(z) \sqrt{x s} \geq M_{\text{BH,min}} \quad (15)$$

on the parameters  $x$  and  $z$ . Because of the monotonically decreasing nature of  $y(z)$ , Eq. (15) 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  $\bar{b} = \bar{z} b_{\text{max}}$ . The total BH production cross section is then [6]

$$\sigma_{\text{BH}}(E_\nu, M_{\text{BH,min}}, M_{10}) \equiv \int_{\frac{M_{\text{BH,min}}^2}{y^2(0)s}}^1 dx \sum_i f_i(x, Q) \hat{\sigma}_{\text{BH}}(x), \quad (16)$$

where  $i$  labels parton species and the  $f_i(x, Q)$  are parton distribution functions (pdfs) [44]. The momentum scale  $Q$  is taken as  $r_s^{-1}$ , which is a typical momentum transfer during the gravitational collapse. A useful criterion for a BH description is that the entropy is sizable, so that Eq. (10) is satisfied. For  $M_{\text{BH,min}} = 3 M_{10}$ ,  $S_{\text{BH}} \approx 13 \gg 1$ .

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<sup>1</sup> In the simple  $T^6$  compactification, there are some unitarization procedures for the multi Kaluza-Klein graviton exchange amplitudes which can yield hadronic cross sections at  $E_\nu \gtrsim 10^{10}$  GeV [42]. However, these cross sections have quadratic energy dependence, and so will generate moderately penetrating showers with definite profiles at lower energies [43]. These have not been reported to date, and their absence places serious constraints on the enhancement of the neutrino-nucleon cross section to sub-hadronic size.

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

$$\sigma_{\text{SB}}(E_\nu, M_{\text{SB},\text{min}}, M_{10}) = \int_{\frac{M_{\text{SB},\text{min}}^2}{s}}^1 dx \sum_i f_i(x, Q) \hat{\sigma}_{\text{SB}}(\hat{s}), \quad (17)$$

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, yielding one or more arbitrary constants. The analysis in the  $\nu q \rightarrow \nu q$  channel illustrates this point [45]. The  $\nu g$  scattering, relevant for  $\nu N$  interactions at ultra-high energies, introduces additional ambiguities. For simplicity, in the spirit of [7], we take  $\hat{\sigma}_{\text{SB}} = (\pi^2 g_s^2 \hat{s}) / (8M_s^4)$ , derived from the Virasoro-Shapiro amplitude. The momentum scale in this case varies from  $Q \simeq \sqrt{\hat{s}}$  at parton center-of-mass energies corresponding to low-lying string excitations, to the BH value  $Q \approx r_s^{-1}$  for  $\sqrt{\hat{s}} \approx M_s/g_s^2$ , the transition point energy. This corresponds to very high string excitations energies. In practice, this may be approximated as a two stage designation:

$$Q(\sqrt{\hat{s}}) \approx \begin{cases} \sqrt{\hat{s}} & \text{if } \sqrt{\hat{s}} \leq M_s/g_s^2, \\ r_s^{-1} & \text{if } M_s/g_s^2 < \sqrt{\hat{s}}, \end{cases} \quad (18)$$

with small ( $\approx 10\%$ ) uncertainty [46].

In Fig. 1 we show a comparison between the total BH production cross section and the particular model we are considering for SBs. As can be seen in Fig. 1, for  $E_\nu \gtrsim 10^7$  GeV,  $\sigma_{\text{SB}} > \sigma_{\text{BH}} > \sigma_{\text{SM}}$ . This makes plausible that in the quantum range,  $M_{10} < M_{\text{BH},\text{min}} < 3M_{10}$ , the BH cross can be thought of as a lower bound on the anomalous  $\nu N$  cross section. Since the BH production cross section is suppressed by  $M_{10}^2$ , bounds on the  $\nu N$  cross section can be translated into lower limits on  $M_{10}$ . If the inequality above holds true in more realistic string models, the bounds on  $M_{10}$  can be significantly strengthened by taking  $M_{\text{BH},\text{min}} = M_{10}$ .

Several techniques have been used to search for the cosmic neutrino component [47]. Showers in the atmosphere can be detected either by observing the fluorescence or Čerenkov light induced by the ionizing tracks in the air, or by directly detecting the charged particles in the shower tail with scintillation counters scattered on the Earth's surface. Neutrino showers mediated by SB/BHs may be distinguished from hadronic primary particles by looking for very inclined showers. This is because for a horizontal path ( $90^\circ$  zenith angle) the neutrino interaction length for SB/BH production is larger than the thickness of the atmosphere ( $36000 \text{ g/cm}^2$ ) and thus, unlike hadrons that shower high in the atmosphere, neutrinos could travel through most of this matter before gravitational collapse is triggered. An alternative technique exploits naturally occurring large volume Čerenkov radiators such as the Antarctic ice-cap. Several radio antennae monitor cold ice for radio frequency Čerenkov radiation resulting from neutrino in-ice cascades. Potential signal events are distinguished from background using vertex location and signal shape.

The experimental situation is that, in spite of extensive searching, only one neutrino-like event has been reported, with an expected background from hadronic cosmic rays of 1.72 events. This implies, at 95% C.L., a maximum of 3.5 events, which combined with the sum of the exposures of the various experiments, and convoluted with the guaranteed cosmogenic neutrino flux, allows the extraction of a lower bound on the fundamental Planck scale,  $M_{10} > 1.0 - 1.4$  TeV, for  $M_{10} < M_{\text{BH},\text{min}} < 3 M_{10}$  [6]. This bound is competitive with that obtained from Tevatron data [48], and these represent the most stringent lower limits to date for the size of 6 large extra dimensions.

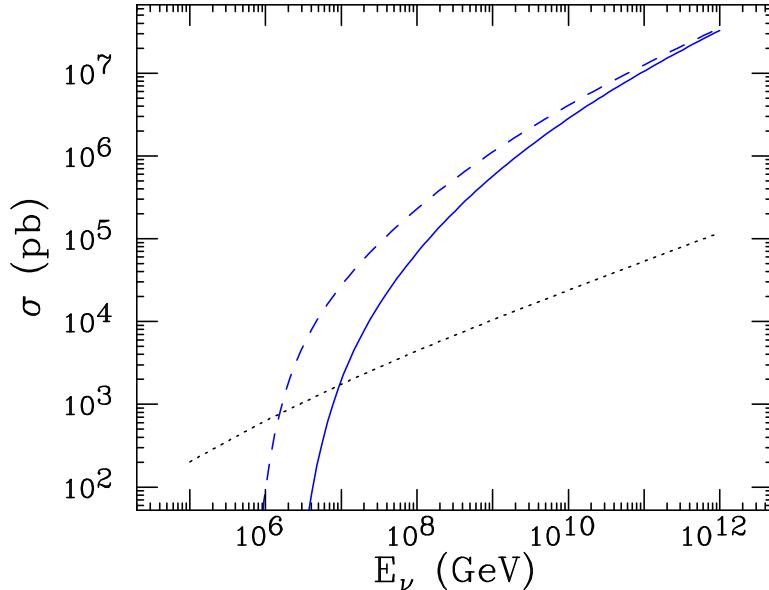


FIG. 1: The solid line indicates the cross section  $\sigma(\nu N \rightarrow \text{BH})$  for  $M_{10} = 1$  TeV and  $M_{\text{BH},\text{min}} = M_{10}$ . Energy loss has been included according to Eq. (16). The dashed line indicates the SB production cross section for  $g_s = 0.2$ . The very lowest threshold was set to  $M_{\text{SB},\text{min}} = 2M_s$ . The SM cross section  $\sigma(\nu N \rightarrow \ell X)$  is indicated by the dotted line.

#### IV. CONCLUDING REMARKS

We have reviewed the possibility of searching for black holes and superstring excitations in cosmic rays using neutrino interactions in the Earth's atmosphere and the Antarctic ice. The fact that neutrino showers have not yet been observed allows us to put a lower bound on the fundamental Planck scale  $M_{10} > 1.0 - 1.4$  TeV. As discussed previously, the string  $\rightleftharpoons$  BH correspondence principle provides support for the higher bound. The next generation of cosmic ray experiments [49, 50] and dedicated neutrino telescopes [51] should probe values of  $M_{10}$  up to about 5 TeV.

Conversely, if neutrino showers are observed above predicted SM rate, they can be ascribed to either new physics or a larger than expected cosmogenic flux. The most powerful technique to distinguish between these two possibilities exploits the absorption of SB/BH secondaries by the Earth, by separately binning Earth skimming events [52] which arrive at very small angles to the horizontal. An enhanced flux will increase both quasi-horizontal and Earth-skimming event rates, whereas a large SB/BH cross section suppresses the latter, because the hadronic decay products of SB/BH evaporation do not escape the Earth's crust [49]. If such a suppression is observed then string balls and black holes will provide a compelling explanation.

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