‘Firewall’ phenomenology with astrophysical neutrinos

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Abstract
One of the most fundamental features of a black hole in general relativity is its event horizon: a boundary from which nothing can escape. There has been a recent surge of interest in the nature of these event horizons and their local neighbourhoods. In an attempt to resolve black hole information paradox(es), and more generally, to better understand the path towards quantum gravity, ‘firewalls’ have been proposed as an alternative to black hole event horizons. In this paper, we explore the phenomenological implications of black holes possessing a surface or ‘firewall’, and predict a potentially detectable signature of these firewalls in the form of a high energy astrophysical neutrino flux. We compute the spectrum of this neutrino flux in different models and show that it is a possible candidate for the source of the PeV neutrinos recently detected by IceCube. This opens up a new area of research, bridging the non-perturbative physics of quantum gravity with the observational black hole and high energy astrophysics.

Keywords: phenomenology of quantum gravity, physics of black holes, quantum aspects of black holes, radiative transfer, scattering, neutrino, muon, pion, and other elementary particles, cosmic rays, accretion and accretion disks

(Some figures may appear in colour only in the online journal)
1. Introduction

One of the most celebrated predictions of general relativity has been the possibility of forming black holes: spacetime singularities that are surrounded by event horizons. These event horizons are boundaries in the spacetime from which nothing can escape. In contrast to black holes, other astrophysical compact objects such as neutron stars, possess a physical surface that shows visible signs of radiation to the outside Universe. In this work we explore the possibility that black holes, too, may not be so invisible after all.

There has been much recent interest in studying the nature of black hole event horizons and their local neighbourhoods, in an attempt to resolve black hole information paradox(es), and more generally to better understand the path towards quantum gravity (e.g., Avery et al 2012, Banks and Fischler 2012, Bena et al 2012, Bousso 2012, Brustein 2012, Chowdhury and Puhm 2012, Giddings 2012, Giveon and Itzhaki 2012, Hossenfelder 2012, Jacobson 2012, Mathur and Turton 2012, Nomura and Varela 2012, Ori 2012, Rama 2012, Susskind 2012a, Susskind 2012b, Almheiri et al 2013, Banks and Fischler 2013, Braunstein et al 2013, Giddings 2013, Harlow and Hayden 2013, Hsu 2013, Hutchinson and Stojkovic 2013, Hwang et al 2013, Kawai et al 2013, Kim et al 2013, Larco et al 2013, Lee and Yeom 2013, Lowe and Thorlacius 2013, Nomura et al 2013a, Nomura et al 2013b, Page 2013, Susskind 2013, Pen and Broderick 2014, Vaz 2014, Saravani et al 2015). In particular, alternatives to classical black hole event horizons such as firewalls (Almheiri et al 2013), fuzzballs (e.g., Mathur and Turton 2012), or other exotic byproducts of quantum gravitational collapse (e.g., Giddings 2014, Dodelson and Silverstein 2015) have been discussed in the literature. It is generally believed that, in lieu of an event horizon, accretion onto (or scattering off) such exotic byproducts may have astrophysical observable consequences that distinguish them from classical black holes (Broderick et al 2009, Pen and Broderick 2014, Broderick et al 2015).

We know that a quantum theory of gravity is non-local on the scale of the Planck length/time, if not longer, which means that a radical change in the structure of space–time within a Planck length of the horizon can affect the microscopic vicinity outside the event horizon. In this work, we assume that gravitational collapse leads to a ‘firewall’ (a term we shall use as a placeholder for any exotic alternative to classical event horizons), that can show visible signs of emission to the outside Universe. Such emission may be possible because an accreting firewall is not in global thermal equilibrium, as the timescale for global thermal equilibrium is generically much longer than the timescale for local radiative processes for macroscopic objects. This is analogous to meteorites hitting the Earth’s atmosphere, where light is emitted from the meteor (within seconds), while the long timescale of heating up the entire Earth and global thermal equilibrium are irrelevant.

We construct a phenomenological model to describe and constrain the nature of this radiation and its spectrum. A key feature of our prediction for this radiation is that it will primarily consist of neutrinos. This is because, as in core-collapse supernovae, which nominally share similar radiative properties to firewalls (e.g. surface density and composition), neutrino propagation has a much smaller optical depth and therefore transports energy out more efficiently.

An interesting implication is that accretion onto black holes/firewalls could be a possible source for the recently detected high energy (PeV) neutrinos by IceCube (Aartsen et al 2014a). The source of this detection is as yet unknown, and several candidate sources have already been ruled out or strongly constrained. For example, hadronuclear pp processes (the mechanism of neutrino production in intergalactic shocks, starburst galaxies, and some active galactic nuclei (AGNs) models) require power laws harder ($p \leq 2.1-2.2$) than those
favoured by IceCube’s detection (Murase et al 2013). Inner jets of blazars may be able to account for IceCube’s PeV signal, but they have difficulty explaining the sub-PeV neutrino events (Murase et al 2014). Further constraints on the candidate sources are set by arguments that no extragalactic object that we know of could function as the origin of both the PeV neutrinos and the ultra high energy cosmic rays (Yoshida and Takami 2014).

We start by outlining the main features and assumptions in our model in section 2. We then summarise our estimate of the total flux emitted by the black holes (section 3) followed by modelling their spectra (section 4). While we consider both blackbody and power law spectra, we subsequently focus on the power law spectrum in section 5 (which are more interesting at higher energies), comparing our results to observations and other models. Finally, section 6 concludes the paper.

2. Modelling accretion onto a firewall

The black holes in our model are accretion powered, have Keplerian flows, and have reached approximate steady state.

We assume that the accreting black holes must show evidence of surface radiation. Broderick et al (2009) considered the possibility of photon emission from the surface of Sagittarius A*. As their predicted photon flux exceeds the Sag. A* flux in millimetre and infrared observations, they conclude that there must be an event horizon. In this work, we argue that at the energy scales involved, the neutrino cross-section is much smaller than the photon cross-section, hence (similar to core-collapse supernovae) most of the energy dissipated in an optically thin region will be carried out by the neutrinos. Based on this assumption, we go on to estimate the total flux of neutrinos from all the black holes (or firewalls) in the Universe.

We estimate that a fraction $\epsilon_g$ (where $0 < \epsilon_g \leq 1$) of the gravitational binding energy of particles falling onto the surface of the black hole is radiated away (and reaches infinity), while the rest goes into growing the mass of the black hole. For a continuous accretion flow, as we are assuming here, this gives a total luminosity as measured at infinity of:

$$L_\infty = \epsilon_g \dot{M} c^2.$$  \hspace{1cm} (1)

A fraction of the total luminosity, $L_\gamma$, is converted by the accretion disc into electromagnetic radiation:

$$L_\gamma = \eta L_\infty.$$  \hspace{1cm} (2)

The remainder of $L_\infty$, will be eventually radiated at the surface as neutrinos. We call this surface luminosity as measured at infinity $L_\nu$:

$$L_\nu \equiv L_\infty - L_\gamma = \epsilon_\nu \dot{M} c^2,$$  \hspace{1cm} (3)

where we have introduced a new efficiency $\epsilon_\nu \equiv \epsilon_g - \epsilon_\gamma \eta$ for the surface radiation that we will be interested in for this work. In terms of the observed total integrated flux:

$$F_\nu = \frac{\epsilon_\nu M c^2}{4 \pi d_L^2},$$  \hspace{1cm} (4)

where $d_L$ is the luminosity distance. Now we want to sum over (4) to find the total flux contribution from all the black holes in the observable Universe. We treat the contributions from the stellar mass black holes and supermassive black holes separately. We estimate that 3.3% (Fukugita and Peebles 2004) of the stellar mass budget goes into forming stellar mass
black holes, while $M_{\text{BH}} \sim 10 M_\odot$ and $M_{\text{SMBH}} \sim 10^6 M_\odot$ are typical masses of stellar mass and supermassive black holes respectively. As a typical accretion rate we take $\dot{M} \sim 0.1 M_{\text{Edd}} = 0.1 \times 16 L_{\text{Edd}} c^{-2}$, where $L_{\text{Edd}}$ is the Eddington limit for the black hole mass.

### 3. Total flux

Let us make the following definitions of characteristic radii

$$ r_h = \frac{GM}{c^2}(1 + \sqrt{1 - a_\ast^2}), \quad (5) $$

$$ \alpha_1 = 1 + (1 - a_\ast^2)^{1/3}(1 + a_\ast)^{1/3} + (1 - a_\ast)^{1/3}, \quad (6) $$

$$ \alpha_2 = \sqrt{5a_\ast^2 + \alpha_1}, \quad (7) $$

$$ r_{\text{ISCO}} = \frac{GM}{c^2}[3 + \alpha_2 - \sqrt{(3 - \alpha_1)(3 + \alpha_1 + 2\alpha_2)}], \quad (8) $$

where $r_h$ is the location of the event horizon, and $r_{\text{ISCO}}$ is the radius of the innermost stable circular orbit. For simplicity, we assume that the black holes in our model are typically spinning with a dimensionless spin parameter of $a_\ast \sim 0.7$.

For the disc efficiency we use Frank et al (2002):

$$ \eta = \frac{1}{\epsilon_g} \left[ 1 - \frac{2M}{3r_{\text{ISCO}}} \right]. \quad (9) $$

For example, for $\epsilon_g = 0.5$, we will have efficiencies of $\eta = 0.2$ and $\epsilon_g = 0.4$ for both stellar mass and supermassive black holes.

We primarily use models from Springel and Hernquist (2003) and Di Matteo et al (2003), for the evolution of the star formation rate (SFR) and the supermassive black hole accretion rate (BHAR) comoving densities, although our final results are not sensitive to our choice of these models. Both SFR and BHAR are fitted by a double exponential function:

$$ \dot{\rho}(z) = \rho_0 \frac{b}{b - a + a e^{b(z - z_m)}} e^{a(z - z_m)}, \quad (10) $$

where the best fit parameters for the SFR are $a = 3/5$, $b = 14/15$, $z_m = 5.4$, $\rho_0 = 0.15 M_\odot \text{ yr}^{-1} \text{ Mpc}^{-3}$, and for the BHAR they are $a = 5/4$, $b = 3/2$, $z_m = 4.8$, $\rho_0 = 3 \times 10^{-4} M_\odot \text{ yr}^{-1} \text{ Mpc}^{-3}$.

The total flux then is

$$ F_{\text{BH} \text{ tot}}[z_{\text{max}}] = \frac{\epsilon_g c}{H_0} \int_0^{z_{\text{max}}} (0.033 \dot{\rho}_{\text{SFR}} c^2) \times \frac{1}{4\pi(1 + z)^3 \sqrt{\Omega_k + (1 + z)^3 \Omega_M}} dz, \quad (11) $$
where we use $H_0 = 73 \text{ km s}^{-1} \text{ Mpc}^{-1}$, $\Omega_A = 0.7$, and $\Omega_M = 0.3$. We are working in comoving coordinates, so the volume elements in (11) and (12) are

$$dV = \frac{c}{H_0} \frac{(1+z)^2 d_A^2}{\sqrt{\Omega_A + (1+z)^3 \Omega_M}} dz,$$

where $d_A = \frac{c}{H_0 (1+z)} \int_0^{z_{\text{max}}} \frac{1}{\sqrt{\Omega_A + (1+z)^3 \Omega_M}} dz$ and $d_L = (1+z) d_A$.

For $z_{\text{max}} = 9$, near the redshift of reionization, equations (11) and (12) yield neutrino fluxes

$$F_{\nu,\text{tot}}[z < 9] \approx 3.8 \times 10^{-4} \text{ erg s}^{-1} \text{ cm}^{-2} \text{ sr}^{-1},$$

$$F_{\nu,\text{BH}}[z < 9] \approx 8.9 \times 10^{-6} \text{ erg s}^{-1} \text{ cm}^{-2} \text{ sr}^{-1}.$$

### 4. Spectrum

To compare to observations, we need information about the spectrum of the neutrino flux. We consider two different models: a power law spectrum and a blackbody spectrum.

#### 4.1. Power law with $p \sim 2-3$

For the power law model, the observed neutrino phase space density $f_{\text{tot}}[E_\nu]$ is

$$f_{\text{tot}}[E_\nu] = \int_{E_{\nu}}^{\infty} \left( \frac{\hat{\rho}(z) c^2 dV}{L_\nu(p)} \right) \frac{\pi r_A^2}{4 \pi d_A^2} \left( \frac{E_\nu (1+z)}{E_o} \right)^{-p-2} dE_\nu,$$

where the value of $E_o$ is fixed by the luminosity:

$$L_\nu(p) = 4 \pi r_A^2 \int_{E_o}^{E_\nu} \frac{E_\nu^3}{\pi^2 c^2 h^3} \left( \frac{E_\nu}{E_o} \right)^{-p-2} dE_\nu,$$

$r_A$ is the apparent radius of the black hole, and $E_p = 1.96 \times 10^{16} \text{ erg}$ is the Planck energy (even though the result is insensitive to this choice for $p > 2$). Note that, given that the neutrino phase space density cannot exceed unity due to the Pauli exclusion principle, $E \sim E_o$ is the natural lower cut-off for a power-law spectrum. Details of how to obtain $r_A$ can be found in Broderick et al (2009); for $a_0 \sim 0.7$ we have $r_A \sim 5 \text{ Gm}^{-2}$.

The results depend on the precise choice of index $p$, which is one of our parameters. We will focus on $p \sim 2-3$. The number flux is

$$\frac{dN}{dt dA d\Omega} = B f_{\text{tot}}[E_\nu] \frac{E_\nu^2}{c^2 h^3 \pi^2},$$

where $B$ is a normalisation constant set by

$$\int_{E_o}^{E_\nu} \frac{E_\nu}{dt dA d\Omega} = F_{\text{tot}}[z_{\text{max}}].$$
We make the simplifying assumption that the neutrinos emitted from the set of stellar mass and supermassive black holes each have approximately the same minimum energy \( E_\circ \) and thus the same spectrum.

An example of a power law with \( p \sim 2 \) is first order Fermi acceleration (Fermi 1949). First order Fermi acceleration is a mechanism that describes the acceleration of charged particles, such as electrons, crossing strong shocks. Upon meeting the shock, due to the disturbance in the magnetic field caused by the shock, there is a probability for the incident electrons to get bumped back the way they travelled. There is also a probability for them to pass through the shock. For those charged particles that bounce back, this process occurs repeatedly until they escape. Each bounce causes the particle to gain energy and the statistics of this process yields a spectrum of:

\[
\frac{dN}{dE_\nu} = \frac{N_0}{E_\circ} \left( \frac{E_\nu}{E_\circ} \right)^{p},
\]

where \( p \sim 2 \) in the non-relativistic case (in the relativistic case, the spectrum is still a power law but the index can vary depending on the details of the scattering; see e.g., Vietri 2008). \( N_0 \) and \( E_\circ \) are the initial number and energies of the charged particles.

Fermi acceleration is ubiquitous for charged particles in astrophysical environments, due to their tight coupling to the magnetic field. Even though neutrinos are not charged and do not interact with an electromagnetic field, they do interact with a gravitational field which varies significantly on short time/length scales close to the firewall. Therefore, depending on the radiative properties of the firewall, it is possible for neutrinos to bounce back and forth between the accretion flow and the black hole firewall (through gravitational scattering), which could yield a power law whose precise index would depend on the details of the propagation model. We shall next discuss a possible mechanism for gravitational Fermi acceleration.

### 4.2. Gravitational Fermi acceleration

Fermi acceleration requires presence of converging magnetic mirrors, which can generically occur in a magnetohydrodynamic (MHD) turbulent medium (2nd order), or around strong shock fronts (1st order). For neutral particles such as neutrinos, a similar effect can happen due to gravitational interactions. In particular, relativistic neutrinos can only be significantly deflected by the gravitational field in the vicinity of black holes. However, null geodesics around ordinary Kerr black holes, generically either fall into the horizon or escape to infinity, and thus do not have the opportunity to up-scatter to higher energies.

The picture can be different if we replace the black hole horizon by a firewall. In particular, one possibility that gives a local description of Bekenstein–Hawking entropy for firewalls excises the spacetime inside the stretched horizon, imposing \( Z_2 \) boundary conditions (Saravani et al 2015). This implies that the firewall may act as a mirror for test particles. As motivated in section 1, we shall assume that this picture is indeed valid for neutrinos due to their weak interaction with matter. However, photons can be absorbed by the firewall, due to their electromagnetic interactions.

We now have our two gravitational mirrors, one provided by the gravity of the black hole, and the other by a(n exotic) firewall. We shall next study the evolution of neutrino energy due to the fluctuations of spacetime geometry, caused by e.g. matter accreting onto a black hole. The energy is given by:
where \( \xi^\mu \equiv \delta^\mu_0 \) would be the time-like Killing vector for a stationary metric, while \( p_\mu \) is the four-momentum. For a time-dependent metric, the rate of energy change is given by:

\[
dE = \frac{1}{2} \frac{1}{g_{\alpha\beta}} \left( g^{\mu\nu} \xi_{(\mu;\nu)} \right) E,
\]

i.e. the relative change in particle energy is roughly proportional to the relative change in the metric:

\[
\Delta \ln E = \frac{1}{2} \int d\lambda \left( \frac{g^{\mu\nu} u^\mu u^\nu}{g_{\alpha\beta}} \xi^\alpha \xi^\beta \right) = \mathcal{O}(\Delta \ln \eta_{\eta_{\mu\nu}}),
\]

where \( u^\mu = d x^\mu / d \lambda \), is defined using an affine parameter \( \lambda \).

The change in particle energy can be divided up into a systematic drift, and a random walk. The systematic change can be calculated using adiabatic invariance, i.e. that action variables for bound orbits of integrable systems (away from the resonances):

\[
J_i \equiv \oint p_i d\tau,
\]

remain invariant under slow changes of the Hamiltonian. Geodesics in Kerr spacetime are indeed integrable, as three integrals of motion (energy, azimuthal angular momentum, and Carter constant; Carter 1968) exist. For null geodesics, the action variables simply scale as \( \sim \mu \sim J_{\text{EREM}} \), so adiabatic invariance yields:

\[
\langle \Delta \ln E \rangle \simeq - \Delta \ln M_{\text{BH}},
\]

i.e. bound relativistic particles systematically redshift with the size of the black hole horizon.

The computation of the random contribution to gravitational redshift is more involved, as it depends on the details of metric fluctuations induced by the turbulent accretion disc. However, this could simply be measured by studying geodesics of test particles in existing general relativistic MHD simulations (with live geometry) of accretion discs in their infalling regions. Let us try to estimate this:

Given that an accretion disc of radius \( R \) is turbulent on the scale of its scale-height, \( H \), we estimate the relative change in energy of a particle in one orbit to be:

\[
\Delta \ln E \sim \Delta \ln g_{\mu\nu} \sim \pm \frac{H}{R} \times \frac{m_{\text{disc}}}{M_{\text{BH}}},
\]

Over several orbits, \( \ln E \) would undergo a random walk (plus the systematic drift in (25)), with the variance:

\[
\langle (\Delta \ln E)^2 \rangle \sim \left( \frac{H}{R} \right)^2 \left( \frac{m_{\text{disc}}}{M_{\text{BH}}} \right)^2 \frac{t}{t_{\text{orbit}}}. \tag{27}
\]

On the other hand, the accretion rate (for a radiatively efficient accretion disc) is roughly given by:

\[
M_{\text{BH}} \sim \alpha_{\text{SS}} m_{\text{disc}} \left( \frac{2\pi}{t_{\text{orbit}}} \right) \left( \frac{H}{R} \right)^2,
\]

where \( \alpha_{\text{SS}} \sim 0.01\ldots0.1 \) is the celebrated Shakura–Sunyaev viscosity parameter (e.g., see Balbus and Hawley 1998). This yields:
Equipped with equations (25) and (29), we can write a Fokker–Planck equation for the evolution of the energy distribution of neutrinos, trapped between a firewall and the gravitational potential barrier of an accreting black hole:

$$\frac{\partial n}{\partial t} = \frac{\partial}{\partial E} \left[ -A n + \frac{1}{2} B \frac{\partial n}{\partial E} \right].$$

(30)

where

$$n \equiv \frac{dN}{dE},$$

(31)

$$A \equiv \frac{d}{dt} \langle \Delta \ln E \rangle = -\frac{M_{\text{BH}}}{M_{\text{BH}}},$$

(32)

$$B \equiv \frac{d}{dt} \langle (\Delta \ln E)^2 \rangle = \frac{1}{2\pi^2 \sigma_{\text{SS}}^2} \times \frac{m_{\text{disc}}}{M_{\text{BH}}} \times \frac{M_{\text{BH}}}{M_{\text{BH}}}.$$  

(33)

The equilibrium solution to the Fokker–Planck equation (30), assuming zero flux through (high or low energy) boundaries, gives:

$$p = 1 - \frac{\partial \ln n}{\partial \ln E} = 1 - \frac{2A}{B} \sim 1 + \frac{4\pi\sigma_{\text{SS}}}{m_{\text{disc}}/M_{\text{BH}}}.$$

(34)

or equivalently

$$p - 1 \sim 1.3 \left( \frac{\sigma_{\text{SS}}}{0.01} \right) \left( \frac{m_{\text{disc}}/M_{\text{BH}}}{0.1} \right)^{-1}.$$  

(35)

The requirement of a significant disc to black hole mass fraction ($\gtrsim 10\%$), to fit the IceCube measurement of $p \sim 2.5$, suggests a stellar origin, e.g., the fallback discs of gamma ray bursts. Alternatively, the accretion rate (28) could be significantly smaller for radiatively inefficient flows (e.g., Yuan et al. 2003), but the details will be more model-dependent in this regime.

Note that this calculation provides the steady-state spectrum of trapped neutrinos. However, we expect the spectrum of neutrinos that escape to infinity to be the same, as the escape probability of relativistic particles is independent of energy, and only depends on geometric factors.

Alternatively, protons could be accelerated via first order Fermi mechanism, and then produce neutrinos through hadronic processes (e.g. Waxman and Loeb 2001). We will consider indices $p \sim 2$–3.

In general, we think that a power law with $p \sim 2$ is a natural choice. This is because high energy processes (such as possible Planck-scale physics within firewalls) favour hard spectra, but the requirement of convergence sets the lower limit of $p = 2$ (equal energy per decade).

4.3. Blackbody

If we assume that the black holes are in approximate steady state and are close to thermodynamic equilibrium, we can consider them as blackbodies. Again, we make the simplifying assumptions that the set of stellar mass and supermassive black holes each have approximately the same surface temperature $T$ and thus spectrum. The temperature can be determined
by the neutrino Stefan–Boltzmann law:

$$L_{\nu, BB} \sim 4\pi r_A^2 T^4.$$  \hspace{1cm} (36)

With $\epsilon_\nu = 0.4$, we find $T_{BH} = 1.2 \times 10^7$ K and $T_{SMBH} = 6.7 \times 10^5$ K. The energy distribution is:

$$\frac{dN}{dE_{\nu}} = \frac{E_{\nu}^2}{\pi^2 c^3 h^3} \frac{1}{1 + e^{E_{\nu}/kT}}.$$  \hspace{1cm} (37)

The number flux is again obtained using (18), with the phase space density

$$f_{\nu \nu} [E_{\nu}] = \int^{\infty}_{-\infty} \frac{\rho(\xi) c^2 dV}{L_{\nu, BB}} \left( \frac{\pi r_A^2}{4\pi d_A^2} \right) \frac{1}{1 + e^{E_{\nu} (1+c)/kT}}.$$  \hspace{1cm} (38)

5. Results

The IceCube collaboration has recently observed a flux of PeV ($10^{15}$ eV) neutrinos (Aartsen et al 2014b), and the source of these high energy neutrinos is as yet unknown. There are only upper bounds beyond 2 PeV, and these are consistent with the power laws we consider in this work.

Figure 1 shows our results for the neutrino spectra for the stellar mass black holes and the supermassive black holes, for the two cases of power law (with $p \sim 2$) and blackbody spectra, as discussed in the previous section. (Note that, figure 1 shows only a sample pair of $\epsilon_\nu$, $p$ values that fit the data. For the complete set of parameters that fit the IceCube data, see figure 3.) Also included in the figure are neutrino flux predictions from models of AGN (Mannheim et al 2001) and GRB’s (Hümmer et al 2012), along with a fit to IceCube’s detection from (Aartsen et al 2014b), and atmospheric neutrino detections. The atmospheric
neutrino observations are from Fréjus (Daum et al. 1995) and AMANDA-II (Abbasi et al. 2010). The blackbody models do not have interesting contributions to the high energy spectra, but it would be interesting in future work to compare these results to upper limits on other sources of low energy neutrinos. In this work, we are interested in high energy neutrinos, so we focus on the power law model results. In Figure 2, we have zoomed in on our results and the other predictions, in the vicinity of the IceCube detection. We see the firewall

**Figure 2.** A close-up of figure (1) within the vicinity of the IceCube detection.

**Figure 3.** IceCube likelihood contours for firewall neutrino emission efficiency $\epsilon_\nu$ and spectral index $p$, for supermassive (Upper) and stellar mass (Lower) black holes. The shaded regions from dark to light correspond to the 1$\sigma$, 2$\sigma$, and 3$\sigma$ contours. These contours are obtained from a likelihood-ratio test using the test statistic $-$2$\Delta \ln L = -2(\ln L_1 - \ln L_{\text{max}})$. 

-neutrino observations are from Fréjus (Daum et al 1995) and AMANDA-II (Abbasi et al 2010). The blackbody models do not have interesting contributions to the high energy spectra, but it would be interesting in future work to compare these results to upper limits on other sources of low energy neutrinos. In this work, we are interested in high energy neutrinos, so we focus on the power law model results. In figure 2, we have zoomed in on our results and the other predictions, in the vicinity of the IceCube detection. We see the firewall
neutrino signal can certainly fit the IceCube detection and is a possible candidate for the source of these high energy neutrinos.

Figure 3 shows the parameter combinations for neutrino emission efficiency $\epsilon_\nu$ and spectral index $p$ for which our power law model yields a good fit to the IceCube data. For the power law models, fits to the IceCube’s data are within the physically acceptable range ($0 < \epsilon_\nu \leq 1$) at 96% and 47% probability for stellar mass black holes and supermassive black holes respectively (or 92% and 37% for $0 < \epsilon_\nu \leq 0.5$). These probabilities can change by $\pm 5\%$ across different models for star formation and supermassive black hole growth. At the 1$\sigma$ confidence level the fit parameters for stellar mass (supermassive) black holes are $\epsilon_\nu \geq 10^{-4}(10^{-2})$ and $p = 2.26$–2.68 (2.27–2.45), while at the 2$\sigma$ level they are $\epsilon_\nu \geq 10^{-5}(10^{-3.5})$ and $p = 2.1$–2.7 (2.1–2.47).

6. Conclusions

To summarise, we have argued that accretion onto astrophysical firewalls is likely to predominantly lead to a flux of cosmological neutrinos. If this spectrum has a power-law shape, e.g. due to an analogue of Fermi acceleration for neutrinos, then it could well fit the observed IceCube spectrum of high energy neutrinos at reasonable efficiencies from astrophysical black holes. Future work should focus on more detailed spectral modelling of neutrino emission from putative firewall atmospheres, as well as other possible observational smoking guns for this scenario. It is certainly exciting to entertain the possibility that neutrino astrophysics could open a new window into the physics of quantum gravity and black holes.

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References

Aartsen E et al (IceCube Collaboration) 2014a Phys. Rev. Lett. 113 101101
Aartsen E(IceCube Collaboration) et al 2014b Phys. Rev. D 91 022001
Abbasi R et al (IceCube Collaboration) 2010 Astropart. Phys. 34 48
Almheiri A, Marolf D, Polchinski J and Sully J 2013 J. High Energy Phys. JHEP02(2013)062
Avery S G, Chowdhury B D and Puhm A 2013 J. High Energy Phys. JHEP09(2013)012
Balbus S A and Hawley J F 1998 Rev. Mod. Phys. 70 1
Banks T and Fischler W 2012 arXiv:1208.4757 [hep-th]
Banks T and Fischler W 2013 arXiv:1305.3923 [hep-th]
Bena I, Puhm A and Vercnocke B 2012 J. High Energy Phys. JHEP12(2012)014
Bouso R 2012 Phys. Rev. D 87 124023
Braunstein S, Pirandola S and Zyczkowski K 2013 Phys. Rev. Lett. 110 101301
Broderick A E, Loeb A and Narayan R 2009 Astrophys. J. 701 1357
Broderick A E, Narayan R, Kormendy J, Perlman E S, Rieke M J and Doeleman S S 2015 Astrophys. J. 805 179
Brustein R 2012 arXiv:1209.2686 [hep-th]
