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Black Hole Chromosphere at the LHC

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If the scale of quantum gravity is near a TeV, black holes will be copiously produced at the LHC. In this work we study the main properties of the light descendants of these black holes. We show that the emitted partons are closely spaced outside the horizon, and hence they do not fragment into hadrons in vacuum but more likely into a kind of quark-gluon plasma. Consequently, the thermal emission occurs far from the horizon, at a temperature characteristic of the QCD scale. We analyze the energy spectrum of the particles emerging from the "chromosphere", and find that the hard hadronic jets are almost entirely suppressed. They are replaced by an isotropic distribution of soft photons and hadrons, with hundreds of particles in the GeV range. This provides a new distinctive signature for black hole events at LHC.

General idea: In the mid-70's Hawking [1] pointed out that a black hole (BH) emits thermal radiation as if it were a black body, with a temperature T_{BH} inversely proportional to its mass M_{BH} . More recently, Heckler [2] noted that for $T_{\text{BH}} > \Lambda_{\text{QCD}}$, the emitted partons do not immediately fragment into hadrons, but propagate away from the BH through a dense quark-gluon plasma losing energy via QCD bremsstrahlung and pair production, and forming a nearly thermal chromosphere. Assuming that TeV scale gravity is realized in nature [3], in this work we first show that such chromosphere would be shielding the BHs to be produced at the Large Hadron Collider (LHC) [4, 5], and then we estimate the emergent thermal photon spectrum from π^0 decay. The ensuing discussion will be framed in the usual context of TeV-scale gravity scenarios, in which the standard model (SM) fields are confined to the brane and only gravity spills into the n internal dimensions of the universe. We also rely on the probe brane approximation, *viz.*, the only effect of the brane field is to bind the BH to the brane. This will be an adequate approximation provided M_{BH} is well above the brane tension, which is presumably of the order of but smaller than the fundamental Planck scale, M_D . Moreover, we assume the BH can be treated as a flat $4+n$ dimensional object. This assumption is valid for both extra dimensions that are larger than the Schwarzschild radius [6],

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

and for warped extra dimensions where the horizon radius is small compared to the curvature scale of the geometry associated with the warped subspace [7].

Hawking evaporation: According to the Thorne's hoop conjecture [8], an event horizon forms when and only when a mass M_{BH} is compacted into a region whose circumference in every direction is less than $2\pi r_s$. The strong gravitational fields around the BH induce spontaneous creation of pairs near the event horizon [1]. While the particle with positive energy can escape to infinity,

the one with negative energy has to tunnel through the horizon into the BH where there are particle states with negative energy with respect to infinity [9]. As the BHs radiate, they lose mass and so will eventually evaporate completely and disappear. The evaporation is generally regarded as thermal in character, with a temperature

$$T_{\text{BH}} = \frac{n+1}{4\pi r_s}, \quad (2)$$

and an entropy $S = 4\pi M_{\text{BH}} r_s / (n+2)$. However, the BH produces an effective potential barrier in the neighborhood of the horizon that backscatters part of the outgoing radiation, modifying the blackbody spectrum. The BH absorption cross section, σ_s (a.k.a. the greybody factor), depends upon the spin of the emitted particles s , their energy Q , and the mass of the BH M_{BH} [10, 11]. At high frequencies ($Q r_s \gg 1$) the greybody factor for each kind of particle must approach the geometrical optics limit. For illustrative simplicity, we adopt throughout this work the geometric optics approximation, and following [12], we conveniently write the greybody factor as a dimensionless constant, $\Gamma_s = \sigma_s / A_4$, normalized to the BH surface area

$$A_4 = 4\pi \left(\frac{n+3}{2} \right)^{2/(n+1)} \frac{n+3}{n+1} r_s^2 \quad (3)$$

seen by the SM fields ($\Gamma_{s=0} = 1$, $\Gamma_{s=1/2} \approx 2/3$, and $\Gamma_{s=1} \approx 1/4$ [12]). 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. This implies that the BH decays with equal probability to a particle on the brane and in the extra dimensions [13]. Therefore, the evaporation process is dominated by the large number of SM brane modes. Since the s -wave greybody factors for fermions [11] differ from the geometric ones by only about 10-20%, we continue to use the latter.

A $4+n$ BH then emits particles with initial total energy between $(Q, Q + dQ)$ at a rate

$$\frac{d\dot{N}_i}{dQ} = \frac{\sigma_s}{8\pi^2} Q^2 \left[\exp\left(\frac{Q}{T_{\text{BH}}}\right) - (-1)^{2s} \right]^{-1} \quad (4)$$

per degree of particle freedom i . Integration of Eq. (4) leads to

$$\dot{N}_i = f \frac{\Gamma_s}{32\pi^3} \frac{(n+3)^{(n+3)/(n+1)} (n+1)}{2^{2/(n+1)}} \Gamma(3) \zeta(3) T_{\text{BH}}, \quad (5)$$

where $\Gamma(x)$ ($\zeta(x)$) is the Gamma (Riemann zeta) function and $f = 1$ ($f = 3/4$) for bosons (fermions). Therefore, the BH emission rate is found to be

$$\dot{N}_i \approx 3.7 \times 10^{21} \frac{(n+3)^{(n+3)/(n+1)}}{2^{2/(n+1)} (n+1)^{-1}} \left(\frac{T_{\text{BH}}}{\text{GeV}} \right) \text{s}^{-1}, \quad (6)$$

$$\dot{N}_i \approx 1.8 \times 10^{21} \frac{(n+3)^{(n+3)/(n+1)}}{2^{2/(n+1)} (n+1)^{-1}} \left(\frac{T_{\text{BH}}}{\text{GeV}} \right) \text{s}^{-1}, \quad (7)$$

$$\dot{N}_i \approx 9.2 \times 10^{20} \frac{(n+3)^{(n+3)/(n+1)}}{2^{2/(n+1)} (n+1)^{-1}} \left(\frac{T_{\text{BH}}}{\text{GeV}} \right) \text{s}^{-1}, \quad (8)$$

for particles with $s = 0, 1/2, 1$, respectively. Note that the BHs to be produced at the LHC are perfectly well defined resonances with mean lifetimes, $\tau_{\text{BH}} \sim M_D^{-1} (M_{\text{BH}}/M_D)^{(n+3)/(n+1)} \approx 10^{-27}$ s, and $T_{\text{BH}} \approx 200$ GeV.

Since thermal fluctuations due to particle emission are small when the entropy $S \gg 1$ [14] and statistical fluctuations in the microcanonical ensemble are small for $\sqrt{S} \gg 1$ [4], we expect the above formulae to be an adequate approximation for $M_{\text{BH}} \gg M_D$. This condition inevitably breaks down during the last stages of the decay process. However, it is noteworthy that for BHs with initial masses well above M_D most of the evaporation takes place within the semi-classical regime. In what follows we require $M_{\text{BH}}/M_D \gtrsim 5$, corresponding to $S \gtrsim 25$. To obtain some quantitative estimates on the particle spectra, we hereafter set $n = 6$ and $M_D = 1.3$ TeV, a value consistent with current limits on the size of extra-dimensions [15].

Thermalization: The evaporation of the BH creates a radiation shell of radius r_s and thickness τ_{BH} , which propagates outward at the speed of light. Using the preceding equations we find that, already for $M_{\text{BH}}/M_D = 5$, there are about 10 quarks and antiquarks in the shell. Before proceeding further, we examine under what conditions these particles can undergo sufficient dissipative interactions (bremsstrahlung and pair production) to initiate thermalization.

At any time t after the end of evaporation, the shell has radius $r = ct$ and thickness τ_{BH} . A quark or antiquark in the shell will interact at a rate given by

$$\Gamma = n c \sigma_{\text{brem}}^{\text{QCD}}, \quad (9)$$

where n is the $q\bar{q}$ effective density at time t , and $\sigma_{\text{brem}}^{\text{QCD}}$ is the cross section for gluon bremsstrahlung in a qq or $q\bar{q}$ collision. The dominant diagram of this process at

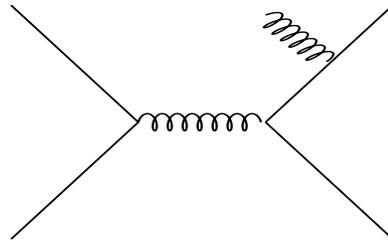


FIG. 1: Feynman diagram for the dominant contribution to gluon bremsstrahlung in a qq or $q\bar{q}$ collision.

low momentum transfer is shown in Fig. 1. The energy-averaged total cross section is given by [2]

$$\sigma_{\text{brem}}^{\text{QCD}} \approx \frac{8\alpha_s^3}{\Lambda^2} \ln\left(\frac{2Q}{\Lambda}\right), \quad (10)$$

where α_s is the QCD coupling constant, and Λ is an infrared cutoff related to the off-shell momentum of the exchanged gluon. Note that the exchanged (virtual) gluon can travel at most a distance $\sim r$ during the interaction, so that the uncertainty principle provides a lower bound $\sim r^{-1}$ on the momentum transfer. Thus, at any time $t = r/c$, the cross section has a maximum value corresponding to $\Lambda \approx r^{-1}$. Of course, Λ is absolutely bounded from below by Λ_{QCD} .

In light of these considerations, we take the effective density to be fully volumetric

$$n = \frac{N_{q\bar{q}}}{\frac{4}{3}\pi r^3} \quad (11)$$

so that a quark or antiquark interacts at a rate

$$\Gamma = \frac{6N_{q\bar{q}}\alpha_s^3}{\pi} \ln\left[\frac{2Q}{\Lambda}\right] \left(\frac{1}{\Lambda r}\right)^2 \frac{1}{r}. \quad (12)$$

For $r \lesssim \Lambda_{\text{QCD}}^{-1}$, we can choose $\Lambda \approx r^{-1}$. In order to remain perturbative, we find the cumulative number of interactions/quark for $\Lambda \sim r^{-1}$ down to some value $\gg \Lambda_{\text{QCD}}$. Integration of (12) with $N_{q\bar{q}} \approx 10$ leads to

$$\mathcal{N}_{\text{int}} \approx 0.15 \left(\frac{\alpha_s(\Lambda)}{0.2}\right)^3 \ln\left[\frac{2Q}{\Lambda}\right] \ln\left[\frac{1}{\Lambda\tau_{\text{BH}}}\right]. \quad (13)$$

The fiducial value of α_s has been chosen appropriate to a momentum transfer scale of 9 GeV [16], corresponding to upilon decay. Even for such a large value of Λ and $Q \simeq 2T_{\text{BH}} \simeq 400$ GeV, Eq. (13) yields ≈ 3 interactions per quark; for $\Lambda = 1.78$ GeV, corresponding to τ -decay [16], $\alpha_s \simeq 0.35$, and Eq. (13) yields ≈ 30 interactions per quark. The creation of a gluon component in the plasma will lead to a significant increase in the bremsstrahlung cascade, because the gluon-gluon cross section is considerably larger than the one of quark-(anti)quark. We take this result as supportive of thermalization, and we proceed on that assumption.

Before interactions have led to a rapid increase in particle numbers, the average separation at distance r of partons emitted by the BH is defined as

$$\langle d \rangle_{\text{parton}} \equiv \left(\sum_{i=q,\bar{q},g} \frac{c_i N_i}{\frac{4}{3} \pi r^3} \right)^{-1/3}, \quad (14)$$

where c_i is the number of internal degrees of freedom of particle species i , i.e., 72 (16) for quarks and antiquarks (gluons). Evaluating this expression for our parameters, we find that $\langle d \rangle_{\text{parton}} \sim 7 \times 10^{-2} \text{GeV}^{-1} \ll \Lambda_{\text{QCD}}^{-1}$. As interactions cascade, the volume occupied by the partons expands as well. The temperature of the chromosphere decreases as the radius increases, and eventually reaches

$$\langle d \rangle_{\text{parton}} \sim \Lambda_{\text{QCD}}^{-1}, \quad (15)$$

at which point quarks and gluons hadronize.

Soft thermal spectrum: The density of the outward-going plasma changes by a large factor within a mean free path, and so the particles never have enough time to fully thermalize. In order to describe the outer edge of the chromosphere we adopt here the heuristic treatment given by Heckler [2]. The quark and gluon spectrum in the observer rest frame is obtained by boosting a thermal spectrum at $T_{\text{ch}} = \Lambda_{\text{QCD}}$ with the Lorentz factor, $\gamma_{\text{ch}} \approx (T_{\text{BH}}/\Lambda_{\text{QCD}})^{1/2}$, of the outer surface of the thermal chromosphere, i.e.,

$$\frac{d\dot{N}_i}{dQ} = c_i \frac{\gamma_{\text{ch}}^2 r_{\text{ch}}^2 Q^2}{2\pi^2} \int_0^1 d\Omega (1 - \beta \cos \theta) \cos \theta \left\{ \exp \left[\frac{\gamma_{\text{ch}} Q (1 - \beta \cos \theta)}{T_{\text{ch}}} \right] - (-1)^{2s} \right\}^{-1}, \quad (16)$$

where the integration over the surface of the chromosphere of radius $r_{\text{ch}} = \gamma_{\text{ch}}/\Lambda_{\text{QCD}}$ has been carried out. It is noteworthy that detailed analyses [17] based on the transport equation yield results in very good agreement with this spectrum.

The observed photon spectrum of the chromosphere is a convolution of the quark-gluon spectrum given in Eq. (16) with the pion fragmentation function and the Lorentz-boosted spectrum from π^0 decay; namely,

$$\frac{d\dot{N}}{dE_\gamma} = \int_{E_0}^{\infty} dE_\pi \frac{dg_{\pi\gamma}(E_\pi)}{dE_\gamma} \frac{d\dot{N}_\pi}{dE_\pi}, \quad (17)$$

where $E_0 = E_\gamma + m_\pi^2/4E_\gamma$. The number of photons with energy E_γ produced by a pion propagating with velocity β and decaying isotropically in its rest frame is

$$\frac{dg_{\pi\gamma}(E_\pi)}{dE_\gamma} = \frac{2}{\gamma m_\pi \beta} = \frac{2}{(E_\pi^2 - m_\pi^2)^{1/2}}, \quad (18)$$

with $\gamma = (1 - \beta^2)^{-1/2}$. Finally, the pion spectrum can be expressed in the form [18]

$$\frac{d\dot{N}_\pi}{dE_\pi} = \sum_j \int_{E_\pi}^{\infty} \frac{d\dot{N}_j(Q, T_{\text{ch}})}{dQ} \frac{dg_{j\pi}(Q, E_\pi)}{dE_\pi} dQ, \quad (19)$$

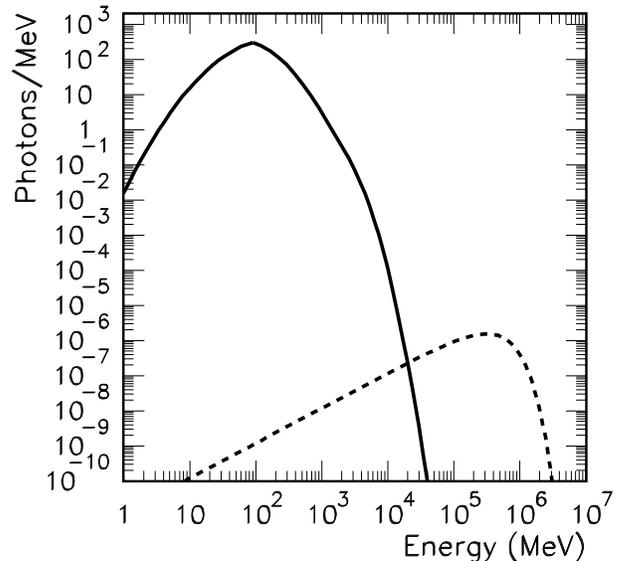


FIG. 2: Photon emission spectra of a 10-dimensional Schwarzschild BH, assuming $M_D \approx 1.3$ TeV and $M_{\text{BH}}/M_D = 5$. This implies, $T_{\text{BH}} \approx 200$ GeV and $\tau_{\text{BH}} \approx 4 \times 10^{-27}$ s. The solid line denotes the spectrum from quark fragmentation and subsequent π^0 decay, whereas the dashed line stands for direct photon emission from the BH.

where the sum is over relevant species in the plasma. The precise nature of the fragmentation process is unknown. As in our previous analysis of BH decay [19], we adopt the quark \rightarrow hadron fragmentation function originally suggested by Hill [20]

$$\frac{dg_{j\pi}}{dE_\pi} \approx \frac{15}{16} z^{-3/2} (1 - z)^2, \quad (20)$$

that is consistent with the so-called “leading-log QCD” behavior and seems to reproduce quite well the multiplicity growth as seen in collider experiments ($z = E_\pi/Q$).

Using Eq. (17) one can easily estimate the mean lifetime of the chromosphere

$$\tau_{\text{ch}} \approx \frac{M_{\text{BH}}}{3} \frac{1}{E_{\gamma_{\text{max}}} \Delta E_\gamma} \left. \frac{d\dot{N}}{dE_\gamma} \right|_{\text{max}}^{-1}, \quad (21)$$

where ΔE_γ is the full width at half maximum (max) of the $d\dot{N}/dE_\gamma$ distribution. The spectrum peaks at an energy of $E_{\gamma_{\text{max}}} \approx 100$ MeV, and thus for our choice of parameters one obtains $\tau_{\text{ch}} \approx 2.2 \times 10^2 \tau_{\text{BH}}$.

In Fig. 2 we show the emergent photon spectrum from the chromosphere hadronization and subsequent π^0 decay, together with that of direct photon emission from the BH. The photon emission rate from the chromosphere is found to be [2]

$$\dot{N}_{\gamma_{\text{ch}}} \approx 2 \times 10^{24} \left(\frac{T_{\text{BH}}}{\text{GeV}} \right)^2 \text{s}^{-1}. \quad (22)$$

Therefore, for a BH with $\tau_{\text{BH}} \approx 4 \times 10^{-27}$ s and $T_{\text{BH}} \approx 200$ GeV, one expects $\approx 7 \times 10^4$ photons as end products of the quark-gluon plasma. The nominal Compact Muon Solenoid (CMS) design anticipates the dynamic range for the electromagnetic calorimeter readout will extend from the noise floor of 25 MeV up to 2 or 2.5 TeV, whereas the hadron calorimeter readout will range from 20 MeV to 2 TeV [21]. Consequently, detection of chromosphere's radiation could become observationally feasible. Moreover, a π^\pm spectrum with roughly the same shape, but shifted in energy by a factor of 2, would trigger the hadron calorimeter.

In general, one expects a statistically significant signal at the LHC from both pp [5] and PbPb collisions [22]. In particular, for the above parameters and assuming a geometric cross section for the colliding partons, about 10^6 BH events with zero net transverse momentum p_T would be produced with an integrated luminosity of 100 fb^{-1} [5]. Proposed modifications, such as BH produced with high p_T [23] and recoil effects from graviton emission into the bulk [24] do not qualitatively change this estimate. Criticisms of the absorptive black disc scattering amplitude, which center on the exponential suppression of transitions involving a (few-particle) quantum state to a (many-particle) semiclassical state [25], have been addressed in [26]. However, it is noteworthy that even if one includes this suppression, the BH production rate will still be quite large at the LHC for over almost all of interesting parameter space [27]. Specifically, for our fiducial values about 10^4 BHs (either at rest or travelling along the beam pipe) would be produced with an integrated luminosity of 100 fb^{-1} .

Note that there is a mild dependence of T_{BH} with the initial BH mass. Consequently, most of the BHs ($M_{\text{BH}}/M_D \gtrsim 5$) to be produced at the LHC have temperatures ≈ 200 GeV. The spectrum from quark fragmentation and subsequent π^0 decay given in Fig. 2 is thus quite general.

Summary and final remarks: The results of this work indicate that the signals for BH production delineated by Giddings–Thomas [4] and Dimopoulos–Landsberg [5] (large transverse momentum, hard electromagnetic component) can be strikingly augmented with the observation of the soft thermal isotropic hadronic and gamma ray component, which includes significant photon and charged pion counting rates in the vicinity of ~ 1 GeV. As can be seen from Fig. 2, one expects about 100 photons (as well as about 200 charged pions) per black hole event with energies in the GeV range. *A hard hadronic component is largely absent.* This is much more dramatic than the absence of hadronic jets with $p_T > r_s^{-1} \sim T_{\text{BH}}$ [28]. According to calorimeter design guidelines for the CMS detector, the observation of this component should be feasible.

TeV-scale BHs can also be produced in cosmic ray collisions [29], and (of course) one expects the quarks and

gluons to be at first freely streaming away from the horizon. However, contrary to collider experiments, the BH is produced with large momentum in the lab system, and its decay products are swept forward with large, $\mathcal{O}(10^6)$, Lorentz factors. This minimizes the effect of the chromosphere on the spectra of the emergent particles analyzed in Ref. [19]. In particular, a 100 MeV pion boosted by 10^6 will still have a mean decay length to muons of ~ 8000 km, much larger than 1 km, which is the estimated mean interaction length of a pion of 10^5 GeV [30]. Thus, an extensive air shower will still occur. The final detectability of the shower will differ little from the situation when the decay particles have energies ~ 100 GeV. If anything, the chromosphere will accelerate the muon production, for reasons similar to the enhancement of the muon component in a shower initiated by a complex nucleus.

Our analysis impacts directly on the identification of the top quark among the BH subproducts [31]. The decay rate of the t -quark in the SM is approximately

$$\Gamma_t \sim \frac{3}{16\pi} m_t m_W^2 (G_F/\sqrt{2}) \sim 0.5 \text{ GeV}, \quad (23)$$

yielding a lifetime $\tau_t \sim 2 \text{ GeV}^{-1}$. Here, m_t (m_W) denotes the mass of the top (W) and G_F is the Fermi coupling constant. Now, since $r_{\text{ch}} \approx T_{\text{BH}}^{1/2}/\Lambda_{\text{QCD}}^{3/2} \approx 150 \text{ GeV}^{-1}$ (and the effects of time dilation are negligible), top quarks emitted as part of the Hawking radiation cannot escape the chromosphere even in the absence of collisions. Diffusive effects will increase the time spent in the chromosphere, and thus strengthen this result. Therefore, top quark identification depends on the leptonic activity resulting from their decay. We also note that both the W and Z decay inside the chromosphere. On the other hand, the lifetime for a Higgs boson decaying into $b\bar{b}$ is

$$\begin{aligned} \tau_{H \rightarrow b\bar{b}} &= \frac{1}{3} (\cos\beta/\sin\alpha)^2 \left(\frac{v/\sqrt{2}}{m_b} \right)^2 \frac{16\pi}{m_H} \\ &\simeq 140 \left(\frac{\cos\beta}{\sin\alpha} \right)^2 \left(\frac{150 \text{ GeV}}{m_H} \right) \text{ GeV}^{-1}, \quad (24) \end{aligned}$$

where m_b (m_H) is the mass of the b -quark (Higgs) and $v = 246$ GeV. In two-Higgs models (such as supersymmetry), $\tan\beta$ is the ratio of the vevs and α is a mixing angle in the neutral Higgs sector. For the SM one-Higgs case, the factor $\cos\beta/\sin\alpha$ is set equal to 1. Thus, in the SM, the Higgs seems likely to escape the chromosphere and generate a hard jetty component in small percentage of the BH decays [32]. In the two-Higgs model, escape from the chromosphere will depend on the mixing angles.

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