SEARCH FOR GAMMA-RAYS FROM BLACK HOLES

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Small black holes with masses much less than the Sun should emit Hawking radiation with a spectrum similar to that of a black body. Searches for such black holes attempt to detect a photon background accumulated over cosmological time or search directly for the final emission stage of individual holes. The ultimate behaviour of a black hole is dictated by particle physics beyond the range of current particle accelerators. We here calculate the experimental signatures of Hawking radiation in the context of the Standard Model of quarks and leptons. The effect of particle degrees of freedom beyond the Standard Model is illustrated. We discuss strategies for improving the existing bound on black hole abundance, obtained from diffuse MeV gamma-ray data,\textsuperscript{1,2} by exploiting air shower telescopes to search for the final $\gamma$-ray emission.

In the last decade a new generation of TeV(10$^{12}$eV) and PeV(10$^{15}$eV) particle telescopes has been constructed, dramatically extending the energies and wavelengths at which the sky can be searched (box A). Their astronomical activity has almost exclusively focused on the search for point sources of ultra-energetic particles.\textsuperscript{3} It has also recently been pointed out that these detectors can probe diffuse high energy photon backgrounds produced by galactic and extragalactic cosmic rays.\textsuperscript{4} In this paper we study another measurement of cosmological relevance, namely the search for Hawking radiation from black holes. We will exclusively discuss primordial black holes (PBHs) produced in the early Universe, since holes produced by astrophysical processes today are too massive to radiate significantly. While the observation of Hawking radiation would be of noblest consequence since it represents a merging of general relativity and quantum mechanics, putting bounds on the PBH abundance is also important since such limits provide direct information on the homogeneity and isotropy of the early Universe.\textsuperscript{5}

The issue of particle emission by PBHs was initially researched over a decade ago\textsuperscript{1,5,6} and has been recently revisited by MacGibbon, Carr and Webber\textsuperscript{2,7–9} in light of the discovery of the quark-lepton structure of matter. On the experimental side the sensitivity and energy-reach of $\gamma$-ray telescopes has improved considerably and such instruments can now probe distances in the TeV- and PeV-energy band larger by about one order of magnitude. We examine the sensitivities of both Čerenkov telescopes and air shower arrays to Hawking radiation in the context of two extreme particle physics models, paying special attention to the capability of the detectors to resolve individual black holes. The relative sensitivity of the TeV- and PeV-detectors and their potential for improving existing bounds on black hole abundance depends strongly on the particle physics model for the final stages of the black hole evaporation. The Standard Model, in which matter has a finite
number of degrees of freedom associated with quarks, leptons and gauge bosons, is
known to be incomplete and additional degrees of freedom are subjects of intense
speculation.

I. PARTICLE MODELS FOR EMISSION

Hawking showed that an uncharged, non-rotating black hole emits particles
with energy in the range \((E, E + dE)\) at a rate\(^6\)

\[
\frac{d^2 N}{dt \, dE} = \frac{\Gamma_s}{2\pi \hbar} \left[ \exp \left( \frac{8\pi GME}{\hbar c^3} \right) - (-1)^{2s} \right]^{-1}
\]

per spin or helicity state. Here \(M\) is the mass of the hole, \(s\) is the particle spin and
\(\Gamma_s\), the absorption probability, is\(^{10}\) in general a function of \(s\), \(E\), and \(M\). In the
\(ME \gg 0\) limit, the emission spectrum is that of a black body and we can associate
a temperature of

\[
T \simeq 1.06 \times 10^{13} \left[ \frac{1}{M} \right] \text{GeV}
\]

with the hole. The spectra as a function of \(ME\) have been calculated by Page\(^{10,11}\)
for spin 1/2 and spin 1 particles and Elster\(^{12}\) and Simpkins\(^{13}\) for spin 0 particles.
The peak flux is

\[
\frac{d^2 N}{dt \, dE} = 9.39 \times 10^{21} \text{GeV}^{-1} \text{s}^{-1}, \quad \text{for } e^\pm
\]

\[
\frac{d^2 N}{dt \, dE} = 1.38 \times 10^{21} \text{GeV}^{-1} \text{s}^{-1}, \quad \text{for } \gamma
\]

after summing over helicity states, at \(ME \simeq 4.26 \times 10^{13} \text{g GeV}\) and \(ME \simeq 6.12 \times
10^{13} \text{g GeV}\), respectively. Integrating over energy, the total instantaneous flux is
\(dN/dt = 3.84 \times 10^{35}/M \text{ s}^{-1}\) for \(e^\pm\) and \(dN/dt = 5.97 \times 10^{34}/M \text{ s}^{-1}\) for \(\gamma\).

The black hole loses mass at a rate\(^{10}\)

\[
\frac{dM}{dt} = -\frac{\alpha(M)}{M^2},
\]

where \(\alpha(M)\) counts the contributions to the energy loss from each emitted species.
As the hole radiates, its temperature rises and \(\alpha\) increases smoothly at each new
particle rest mass threshold (see Fig. 1). Above each threshold, $\alpha$ is approximately

$$\alpha \simeq [7.8 \, d_{s=1/2} + 3.1 \, d_{s=1}] \times 10^{24} \, \text{g}^3 \, \text{s}^{-1},$$

where $d_{s=1/2}$ and $d_{s=1}$ are the number of degrees of freedom (spin, charge, and color) of the emitted particles. For the Standard Model with three fermionic generations, we have $d_{s=1} = 27$ and $d_{s=1/2} = 90$ at $T \sim 100$ GeV and so $\alpha \leq 7.8 \times 10^{26} \, \text{g}^3 \, \text{s}^{-1}$ (ignoring contributions from Higgs scalars and gravitons). In Quantum Chromodynamics, we regard hadrons as composed of quarks and gluons, once the emitted energies exceed the QCD quark-hadron confinement scale, $\Lambda_{\text{qh}} \simeq 100$-300 MeV. Free quarks and gluons are emitted rather than composite hadrons above $\Lambda_{\text{qh}}$ and, of the hadrons, only pions can be directly emitted below $\Lambda_{\text{qh}}$. The relativistic quarks and gluons subsequently fragment into the astrophysically stable photons, neutrinos, electrons, positrons, protons and antiprotons. The timescale between successive emissions is short enough that we can neglect interactions between successive emissions before fragmentation.\footnote{14} Fluctuations in the hole's color charge should also average to zero. The emitted quarks and gluons can therefore be treated as asymptotically free partons hadronizing via jets far beyond the black hole horizon. This is analogous to the production of jets in $e^+e^-$ annihilation.

In this picture, we can then convolute Eq. (1) with empirical parton fragmentation functions to obtain the instantaneous emission spectra, or alternatively use a Monte Carlo $e^+e^-$ event generator.\footnote{7} The resulting $e^\pm$, $\gamma$ and $\nu\bar{\nu}$ flux at $T = 0.3$-100 GeV has a dominant peak around 100 MeV coming from jet pion decays and a slight bump around $5T$ coming from non-decaying direct emission. The $\beta$-decay of jet neutrons produces low energy bumps in the lepton spectra around 1 MeV. These fluxes are shown in Fig. 2 for a $T = 10$ GeV black hole. For $T \gtrsim 100$ GeV, where the correct particle model is not yet known, we can only approximate the fragmentation functions and extrapolate accordingly. The value of $\alpha$ in Eq. (5) is a minimum bound at these energies. The Standard Model may be incomplete and new degrees of freedom could emerge. For example, if Supersymmetry is the underlying theory for elementary particles, $\alpha$ should increase\footnote{8} by a factor of at least three.

In alternative scenarios, Eq. (4) might not represent the correct way to count particle degrees of freedom as the temperature approaches $\Lambda_{\text{qh}}$, even in the Standard Model. For example, if hadrons are made of strings of three quarks or strings of quark-antiquark pairs, the relevant degrees of freedom are those of the string, not of
the underlying quarks. Each degree corresponds to a particle state of definite mass and spin. In such models, the number of states grows exponentially with mass\textsuperscript{15,16}

\[ N \sim m^{-\frac{3}{2}} \exp \left[ \frac{m}{\Lambda} \right]. \]  

(6)

In the Hagedorn picture (which is motivated by the apparent exponential increase in hadronic resonances seen in accelerators), \( \Lambda \) has a typical value of \( \Lambda_H = 160 \text{ MeV} \). The exponentially growing density of states causes the final explosion to be much more violent. Once a limiting temperature is reached, the remaining mass of the hole is suddenly emitted into all available states. This occurs when the exponential increase in the number of states offsets the exponential decrease in the tail of the Hawking radiation. It corresponds to a black hole mass of \( M_c = (1.06 \times 10^{13} \text{ GeV})/\Lambda \text{ g}. \)

While accelerator physics has shown no evidence for Hagedorn type models, and lattice calculations and cosmology do not favour them, such models are not ruled out. The quark hadron phase transition sets in at energies lower than the energies at which perturbative QCD calculations can be applied. Collider experiments, which probe smaller and smaller distances at energies well above \( \Lambda_{qh} \), show evidence for the asymptotic freedom of QCD. The question of what happens around \( \Lambda_{qh} \) however remains uncertain. In the Hagedorn picture of black hole bursts, the large number of states modifies the entropy increase preventing the temperature from rising above \( \Lambda_{qh} \).

Regardless of the behaviour around \( \Lambda_{qh} \), an exponentially growing spectrum of states may also occur at the ultra-high energies associated with superstring theories \((\Lambda = 10^{14} - 10^{16} \text{ GeV})\). In this case, the hole's behaviour is only modified in the very last stages, when the mass and lifetime are small enough that the integrated emission spectrum does not change significantly.

More realistically, equations of the form (6) may indicate that a phase transition, whose details are as yet unknown, occurs at \( \Lambda \). In this case, the quasi-exponential spectrum would not continue until \( m \to \infty \) but would last for, at most, the duration of the finite phase transition. The evolution would then proceed smoothly once a temperature above \( \Lambda \) is attained. Applying this to black hole evaporation, we could expect the hole to emit a burst which carries off a significant fraction of its total energy when it reaches \( \Lambda_H \); but subsequently to evaporate as described by Eq. (4) at temperatures above \( \Lambda_H \). Because the burst would last a very short time (certainly when compared to Galactic timescales), and since \( M = 1/T \)
implies that a hole loses a significant fraction of its remaining mass at any given temperature anyway, this should not change the photon background emitted by a PBH distribution (see Section II). We would expect to observe, though, occasional bursts above the background from holes just reaching \( \Lambda \).

In Sections V and VI, we investigate the chance of detecting the final evaporation stage in both the standard and Hagedorn scenarios. These scenarios represent the two extreme extrapolations of the particle model beyond accelerator energies.

II. THE MeV LIMIT AND ITS IMPLICATIONS

By integrating Eq. (4), one can relate\(^{10}\) the lifetime of a black hole, \( \tau_{\text{evap}} \), to its initial mass \( M_i \). To a good approximation, we have

\[
M_i \simeq [3\alpha(M_i)\tau_{\text{evap}}]^{1/3} \\
\simeq 5.4 \times 10^8 \left[ \frac{\alpha(M_i)}{5.3 \times 10^{25} \text{ g}^3 \text{ s}^{-1}} \right]^{1/3} \left[ \frac{\tau_{\text{evap}}}{1 \text{ s}} \right]^{1/3} \text{ g}.
\]

The mass of a hole which formed in the Early Universe and is just completing its evaporation today is thus \( M_* \simeq [3\alpha(M_*)t_0]^{1/3} \) where \( t_0 \) is the age of the Universe. Because \( M_* \) is close to the thresholds for \( \mu^\pm, \pi^{0, \pm} \), and \( u, d \) quark emission, we integrate Eq. (4) and solve numerically for \( M_* \) by approximating \( \alpha \) linearly at each threshold.\(^{17}\) Figure 3 shows the solutions for \( M_* \) as a function of \( \Lambda_{\text{qh}} \) for extreme cases corresponding to Friedman models with \( \Omega_m = 1 \) and 0.06 and Hubble constant \( h_0 = 0.4 \) and 1 where \( H_0 = h_0/100 \text{ km s}^{-1} \text{ Mpc}^{-1} \). We find that \( M_* \) lies between\(^8\)

\[
4.3 \times 10^{14} < M_* < 7.0 \times 10^{14} \text{ g}
\]

for \( 125 \lesssim \Lambda_{\text{qh}} \lesssim 300 \text{ MeV} \), 0.06 \( \lesssim \Omega_m \lesssim 1.0 \) and 0.4 \( < h_0 < 1 \). Since \( \tau_{\text{evap}} \propto M_i^3 \), black holes with \( M_i \gg M_* \) have not lost a significant amount of mass by today while those with \( M_i < M_* \) have completely evaporated.

PBHs may have formed in the Early Universe in a number of scenarios.\(^{18}\) For example, they may have been produced by initial density inhomogeneities\(^{19}\) or as a result of a phase transition.\(^{20}\) Hawking\(^{21}\) and Zembowicz and Polnarev\(^{22}\) have also recently proposed that a spectrum of black holes may be created by the collapse of cosmic strings. If spherically symmetric PBHs form from scale-invariant initial
density perturbations that have a Gaussian distribution, the number density of holes created with masses in the range \((M_i, M_i + dM_i)\) is

\[
\frac{dn}{dM_i} = (\beta - 2) \, \Omega_{\text{pbh}} \, \frac{\rho_c}{M_*^2} \left( \frac{M_i}{M_*} \right)^{-\beta}.
\]  

(9)

Here \(\rho_c\) denotes the cosmological critical mass density and \(\Omega_{\text{pbh}}\) is the present fraction of \(\rho_c\) in \(M \geq M_*\) holes. We define, for convenience, \(N = (\beta - 2)\Omega_{\text{pbh}} \rho_c/M_*\). (\(N\) corresponds to the initial number density of holes per logarithmic mass interval at \(M = M_*\).) If the equation of state of the Universe is \(p = \gamma p\) at the formation epoch, then \(\beta = (1 + 3\gamma)/(1 + \gamma) + 1\). For PBHs forming in the radiation dominated era, we have \(\beta = 2.5\).

The emission over the lifetime of the Universe will produce diffuse particle backgrounds. These can be compared with\(^1\) observations to obtain an upper limit on \(\Omega_{\text{pbh}}\). The diffuse photon spectrum is shown in Fig. 4 for the initial distribution Eq. (9) with \(\beta = 2.5\). We have included quark and gluon emission and neglected all losses other than redshift. The general form of the integrated emission\(^{17}\) is an \(E^{-1}\) slope below 100-300 MeV for \(\beta < 3\), turning over into an \(E^{-3}\) slope above 300 MeV for all \(\beta\). The \(E^{-1}\) slope is determined by jet fragmentation; the \(E^{-3}\) slope comes from evaporation in the present epoch; and the turnover at 100-300 MeV corresponds to direct photon emission. Matching to the observed gamma-ray spectrum at 100 MeV (where photon losses are negligible), MacGibbon and Carr find that\(^2\)

\[
\Omega_{\text{pbh}} \lesssim 7.6(\pm 2.6) \times 10^{-9} h_0^{-1.95 \pm 0.15}
\]  

(10)

for \(\Omega_m = 1.0\) and \(\beta = 2.5\). The limit is about 60% weaker for \(\Omega_m = 0.06\) and is also \(\lesssim 50\%\) sensitive to the value of \(\Lambda_{\text{qcd}}\). MacGibbon and Carr\(^2\) also obtain bounds on \(\Omega_{\text{pbh}}\) from the antiproton, electron and positron fluxes, again including jet fragmentation. This corresponds to \(N \lesssim 10^4\) pc\(^{-3}\). These bounds are remarkably similar to the gamma-ray limit (10) if the PBHs are clustered to the same degree as other material in the Galactic halo. The bounds may even all overlap, allowing the possibility that PBH emission may be contributing significantly to all observed interstellar cosmic ray and gamma ray spectra between 0.1-1 GeV. If the initial distribution (9) does not apply, or if \(\beta \neq 2.5\), the limit given in Eq. (10) is still approximately the limit on the present fraction of \(\rho_c\) in holes of mass \(M_*\) since it is derived at the peak of the emission from these holes.
From the initial PBH distribution, we can calculate the distribution of holes at any later time \( t \). Invoking Eq. (7), the PBH distribution of Eq. (9) becomes

\[
\frac{dn}{dM} = \frac{dn}{dM_i} \frac{dM_i}{dM} \\
\simeq \frac{N}{M_*} \left[ 1 + \left( \frac{t}{t_0} \right) \left( \frac{\alpha(M)}{\alpha(M_*)} \right) \left( \frac{M}{M_*} \right)^{-3} \right]^{-\frac{\alpha_3^+}{3}} \left[ \frac{M}{M_*} \right]^\gamma \left( \frac{\alpha(M)}{\alpha(M_*)} \right)^{\frac{\beta-1}{3}}
\]

\[
\propto \begin{cases} 
M^2, & M \ll 10^9 \, t^{1/3} \, \text{g} \\
M^{-\beta}, & M \gg 10^9 \, t^{1/3} \, \text{g}
\end{cases}
\]  

(11)

where \( M \) is the mass at time \( t \) of a hole with initial mass \( M_i \). Note that the present distribution for \( M \ll M_* \) is independent of \( \beta \): regardless of their formation mechanism, holes today with masses less than \( M_* \) have the distribution \( dn/dM \propto M^2 \). From Eq. (11), or directly from Eq. (4), we now obtain the current rate of expiring PBHs,

\[
\frac{dn}{dt} = \frac{\alpha(M_*)}{M_*^2} N \simeq \frac{N}{t_0}.
\]  

(12)

Allowing for uncertainties in \( M_* \) and \( \alpha(M_*) \), the present rate of explosions averaged over the Universe should thus lie in the range

\[
1.2 \times 10^{-11} \, \text{yr}^{-1} \, (h_0 = 0.4) < \frac{dn}{dt} \, N^{-1} < 4.0 \times 10^{-11} \, \text{yr}^{-1} \, (h_0 = 1.0).
\]  

(13)

It is also reasonable to expect that PBHs cluster in galactic halos along with other material, since they should not have large peculiar velocities. This would enhance the local PBH density by a factor \( \zeta \). Taking \( ^{24}R_\odot = 8.5(\pm 1.1) \) kpc as the Sun’s distance from the Galactic centre and \( \Theta(R_\odot) = 2.20(\pm 0.20) \times 10^7 \) cm s\(^{-1} \) as the circular velocity at the Sun, the mass of the Galaxy within the Sun’s orbit is \( M_{G\odot} = R_\odot \Theta_\odot^2 / G = 1.94(\pm 0.06) \times 10^{44} \) g. The average density within the sphere defined by the Sun’s orbit is then

\[
\bar{\rho}_\odot = \frac{3M_{G\odot}}{4\pi R_\odot^3} = 2.56(\pm 1.79) \times 10^{-24} \, \text{g cm}^{-3}.
\]  

(14)
Assuming an isothermal halo, this gives a local enhancement factor at $R_\odot$ of

$$\zeta = \frac{\rho_\odot}{3 \Omega_\odot \rho_c} = 1.36(\pm 0.9) \times 10^7 \left( \frac{\Omega_\odot}{0.01} \right)^{-1} h_0^{-2},$$

(15)

where $\Omega_\odot$ is the visible fraction of $\rho_c$ in galaxies out to $R_\odot$—the visible fraction approximately also equals the dark component, up to a factor of 2. The large uncertainties associated with clustering strongly affect the possibilities of detecting PBH evaporations. For $N = 10^4 \text{ pc}^{-3}$, the present rate of explosions could be as low as $10^{-7} \text{ pc}^{-3} \text{ yr}^{-1}$, if holes are unclustered, or as high as $10 \text{ pc}^{-3} \text{ yr}^{-1}$ if holes are clustered. Since gamma-rays (unlike charged particles) are not confined in the Galaxy, though, the PBH photon background is not affected by clustering.

III. FINAL STAGE EMISSION

The spectrum of final radiation from a black hole is dictated by particle physics. Here we consider two extreme models. In the first, the particle degrees of freedom at very high temperatures are those of the Standard Model. A black hole of mass $M_B$ evaporates completely over the time $\Delta \tau$ given by Eq. (7). If $\Delta \tau$ is small compared with the exposure time of the detector, we should consider the flux integrated over $\Delta \tau$. The photons are emitted directly by the Hawking mechanism and also via the decay of quark and gluon jets. The direct quark flux integrated over $\Delta \tau$ peaks at an energy of about

$$Q \simeq 40 \left( \frac{1 \text{ s}}{\Delta \tau} \right)^{1/3} \text{ TeV}.$$ 

(16)

After fragmentation, the photon flux comes mainly from $\pi$ decays and peaks in the vicinity of the pion mass $m_\pi$, several orders of magnitude below $Q$. Not surprisingly, the spectrum is very sensitive to the details of jet fragmentation. As an example, let us assume a fragmentation function of the form,

$$dN_\pi/dz = (15/16)z^{-3/2}(1 - z)^2$$

where $z = E_\pi/E_{\text{jet}}$, which gives an empirically-motivated multiplicity growth of
The decay photons then have the distribution,

\[
\frac{dN_\gamma}{dE_\gamma} = 2 \int_{E_\gamma}^{E_\text{jet}} \frac{1}{E_\pi^0} \frac{dN_\pi^0}{dE_\gamma} \frac{dE_\pi^0}{dE_\gamma}
\]

\[
= 5 \left[ 1 + \frac{1}{8} \left( \frac{E_\gamma}{E_\text{jet}} \right)^{-3/2} - \frac{3}{4} \left( \frac{E_\gamma}{E_\text{jet}} \right)^{-1/2} - \frac{3}{8} \left( \frac{E_\gamma}{E_\text{jet}} \right)^{1/2} \right].
\]

Integrating the photon flux from quark and gluon jets above a detector threshold \( E_D < Q \), we find that

\[
\frac{dN_\gamma}{dt} (> E_D) \simeq 8.0 \times 10^{23} \frac{Q}{1 \text{ GeV}} \left[ \frac{1}{4} \left( \frac{Q}{E_D} \right)^{1/2} \left( 1 + \frac{E^2}{Q^2} \right) - 1 + \frac{3}{2} \left( \frac{E_D}{Q} \right)^{1/2} - \frac{E_D}{Q} \right] \text{s}^{-1}.
\]

Here we have made the approximation that the total instantaneous flux \( dN/dt \) per degree of freedom, as given in Section I, is emitted at the corresponding direct peak.

Alternatively, if we convolute Eq. (17) with the direct jet spectra and integrate over \( \Delta \tau \), the combined direct and jet-decay photon flux decreases as \( E_\gamma^{-3} \) above \( E_\gamma \simeq Q \) (since \( \tau_{\text{evap}} \propto M^3 \)). Below \( E_\gamma \simeq Q \), it traces the slope of the fragmentation function and is eventually cut off at \( E < m_\pi \). The total number of photons \( > E_D \), integrated over \( \Delta \tau \), now becomes

\[
N_\gamma (> E_D) \simeq 2.4 \times 10^{37} \left( \frac{\text{GeV}}{Q} \right)^2,
\]

\[
\times \begin{cases} \left[ \sqrt{\frac{Q}{E_D}} \left( \frac{5}{6} + \frac{3}{14} \frac{E_D}{Q} + \frac{5 E_D^2}{Q^2} \right) - 5 \frac{E_D}{Q} - \frac{5}{2} \right] + \frac{1}{250}, & E_D < Q; \\
\left( \frac{Q}{E_D} \right)^2 \left[ \frac{1}{42} + \frac{1}{150} \left( 1 - \frac{2}{5} \max \left( 0, 2 - \frac{Q}{E_D} \right) \right) \right], & E_D \geq Q. \end{cases}
\]

The direct photons contribute the last terms in these expressions. Decay photons dominate at all energies. This is because there are 72 quark and 16 gluon degrees of freedom, but only two direct photon degrees. The contribution per helicity state is also greater for \( s = 1/2 \) than for \( s = 1 \) states. The enhancement due to the jet decays, however, is critically sensitive to the \( z \rightarrow 1 \) behaviour of the fragmentation function, which is unknown at temperatures above current accelerator energies. At energies close to the black hole temperature, other processes such as prompt photon production may also enhance the flux.
The neutrino and cosmic ray emission from the PBHs is comparable with the photon emission. Since the neutrino and cosmic ray backgrounds are greater than any postulated diffuse photon background, however, photon detectors offer the greatest sensitivity for searching for PBH evaporations.

IV. INDIVIDUAL HOLE EMISSION AND DIFFUSE GAMMA RAYS

Can an individual black hole be detected above the background created by all evaporating PBHs or above an extrapolation of the observed extragalactic background? Satellite measurements of the extragalactic diffuse 35-170 MeV photon background show that the flux falls off roughly as $E^{-2.4}$. Measurements above 170 MeV have not yet been made. The EGRET instrument, due for launch in November 1990, will be capable of detecting photons up to energies of about 20 GeV. However, since the galactic emission falls off less steeply than the extragalactic background, it is not known if the diffuse background can be resolved from the galactic emission above a few hundred MeV, even at high Galactic latitude. Regardless of their source, photons with energies $> 10^6$ GeV are cut off by pair-production on the cosmic microwave background.

Let us first consider an individual PBH. If $n_{bh}$ is the local number density of PBHs with present mass less than $M_{th}$, then integrating Eq. (9), we have

$$n_{bh} \simeq \zeta N \frac{\alpha(M_*)}{3 \alpha(M_{th})} \left[ \frac{M_{th}}{M_*} \right]^3$$

for $M_{th} \ll M_*$. In a volume $n_{bh}^{-1}$, we expect to find the closest hole, with, on average, a mass of $0.75 M_{th}$ according to Eq. (20). The photon flux at the detector from this hole is thus

$$\frac{d^3 N_{\gamma}}{dA \, dt \, dE_\gamma} = \frac{d^2 N_{\gamma}}{dt \, dE_\gamma} \left[ \frac{M_{th}}{M_*} \right]^2 \left[ \frac{\alpha(M_*)}{\alpha(M_{th})} \right]^{\frac{3}{2}} \zeta^{\frac{3}{2}} N^{\frac{3}{2}} \frac{8}{3 \sqrt{12} \pi},$$

where $d^2 N_{\gamma}/dt \, dE_\gamma$ is given by Eq. (3) for direct photons, or is derived from Eq. (17) for decay photons.

On the other hand, the background emission from PBHs can be parameterized
by\textsuperscript{17}

\[ \frac{d^4N_\gamma}{dA\,dt\,d\omega\,dE_\gamma} \simeq 7 - 14 \times 10^{-11} \left[ \frac{N}{1\ \text{pc}^{-3}} \right] E^{-3}\ \text{cm}^{-2}\ \text{s}^{-1}\ \text{sr}^{-1}\ \text{GeV}^{-1} \] \hspace{1cm} (22)

above 300 MeV, for $0.06 \leq \Omega_m \leq 1.0$, $0.3 \leq h_0 \leq 0.7$, and $\beta = 2.5$. The observed photon background between 35-170 MeV is\textsuperscript{8}

\[ \frac{d^4N_\gamma}{dA\,dt\,d\omega\,dE_\gamma} = 2.7(\pm 0.5) \times 10^{-4} \left[ \frac{E_\gamma}{100\ \text{MeV}} \right]^{-2.4(\pm 0.2)} \ \text{cm}^{-2}\ \text{s}^{-1}\ \text{sr}^{-1}\ \text{GeV}^{-1}. \] \hspace{1cm} (23)

In order to estimate the detector resolution necessary to resolve a point source out of the diffuse background, we require that the diffuse photon flux within the resolved solid angle of the detector does not exceed the point source flux, i.e.,

\[ \frac{d^3N_\gamma}{dA\,dt\,dE_\gamma} \gtrsim \omega_D \frac{d^4N_\gamma}{dA\,dt\,d\omega\,dE_\gamma}. \] \hspace{1cm} (24)

As an example, let us substitute Eqs. (21) and (22) into (24) assuming $\omega_D = 10^{-3}$ and $\zeta N \simeq 10^{10}$ and consider only the direct photon emission. In this case, black holes with mass $M < 4 \times 10^{10}$ g will stand out from the PBH background above about 2 TeV. If we compare (21) instead to the extrapolation of the measured background (23), the critical energy is about two orders of magnitude greater. Alternatively, if we consider the decay distribution (17) and look at $E_\gamma/E_{\text{jet}} \simeq 0.1$, we find that both critical energies are about two orders of magnitude smaller. Thus it may be possible to resolve the instantaneous decay emission from the backgrounds at energies above 0.01-1 TeV. This conclusion, however, strongly depends on the correct form of the fragmentation function as $E_\gamma/E_{\text{jet}} \rightarrow 1$ and the degree to which the PBHs cluster.

For an estimate more relevant to detectors, we integrate the individual flux above the detector threshold $E_D$. Figure 5 shows the distances at which a hole of temperature $T$ produces a flux in the detector [Eq. (18)] equal to the maximum diffuse PBH background [Eq. (22)] allowed by the 100 MeV data. Direct photons produce the lower temperature bump and decay photons produce the higher temperature bump. If the PBHs are clustered with $N\zeta = 10^{12}$, then detectors with $E_D \gtrsim 30$ GeV may be able to resolve $T \gtrsim 20$ GeV holes from the background.
However if $N\zeta = 10^{10}$, the energy threshold for the detector must exceed 300 GeV and the black hole temperature 700 GeV. For smaller temperatures, the detector should see a sky uniformly bright in black holes and we have another version of Olber's paradox for black holes. In this case, the black hole horizon is extremely small and the radiation density of the background sky is much smaller than the black hole surface radiation density. No practical detector can single out a point source because a large number of sources appear in any field of view.

Although the results are model dependent, the general features illustrated in Fig. 5 should be correct. The distance to the nearest hole of a given temperature should scale as $\zeta^{-1/3}$. The curves should shift upward as $\sqrt{10^{-3}/\omega_F}$ with improved detector resolution. Optimum signal/noise is obtained at $E_D \sim T$. In Fig. 5 we use an $E^{-3}$ diffuse background. If the background falls as $E^{-2.4}$, as extrapolated from the MeV data, the conclusions are more restrictive.

At high energies, air shower arrays also have the capability to search for a diffuse photon background, e.g. by exploiting the small muon abundance in photon-induced air showers to reject cosmic rays. From Eq. (22) we expect a flux of

$$\frac{d^2N_\gamma}{dA dt} (> 100 \text{ TeV}) < 3 \times 10^{-16} \text{ cm}^{-2} \text{ s}^{-1}. \quad (25)$$

This flux however is below that from known TeV $\gamma$-ray sources such as galactic cosmic rays interacting with the interstellar medium. Note, however, that the lifetime of a 10 TeV black hole is of the order of 1 s [see Eq. (7)], we now consider alternative search techniques based on individual hole detection.

V. DETECTION POSSIBILITIES IN THE STANDARD MODEL

Ground-based gamma-ray telescopes, which are sensitive in the TeV and PeV energy ranges (see box B) can be used to detect, or to set limits on, the density of PBHs expiring in the present epoch. Experimentally there are two fundamental considerations in such searches: (i) how large a volume can the telescope search efficiently and for what exposure time; and (ii) what distinguishes individual PBH emissions from a chance coincidence of background cosmic ray showers. If the detector parameters are well known it is straightforward to set an upper limit on the PBH density; it is much more difficult to prove that any statistically significant excess of candidate events is of PBH origin. All searches so far have proved unsuccessful.
A renewed search for individual PBHs is prompted by the development of new atmospheric Čerenkov and air shower array telescopes which can identify gamma-ray showers amid the much greater background of charged cosmic rays; and by the revised Standard Model prediction for final evaporation (see Section IV). The proposed detection techniques are based generally on the same principles as before.

At 100 TeV energies, experiments are under construction with collection areas in excess of $10^{10}$ cm$^2$ and angular resolutions less than $4 \times 10^{-4}$ sr. Hadron rejection ratios (based on the muon-to-electron ratio) of $10^{-4}$-$10^{-5}$ are anticipated although this has yet to be demonstrated in the detection of a discrete source. Air shower arrays are now being discussed with energy thresholds in the 10-100 TeV range. Unfortunately near threshold, the angular resolution, collection areas and rejection ratios are not optimal. As an illustration, let us consider a somewhat idealized air shower array telescope with the following parameters: collection area $A_D = 10^8$ cm$^2$; angular resolution $\omega_D = 4 \times 10^{-4}$ sr; energy threshold $E_D = 10$ TeV; operation time = 1 year; and hadron rejection ratio = $10^{-3}$. We will assume a typical zenith angle cutoff of 30 degrees, and hence an acceptance angle of $\omega_A = 0.75$ sr. The depth of sky covered by the telescope can be calculated from the minimum flux sensitivity for detection of a burst. If $n_\gamma$ is the minimum number of gamma-rays of energy $> E_D$ required to register as a burst, then

$$n_\gamma = \frac{A}{4 \pi r^2} N_\gamma(> E_D),$$

(26)

where $r$ is the distance to the PBH.

From Eq. (7), both Extensive Air Shower arrays and Čerenkov telescopes become sensitive to the direct photon flux from a black hole at a time

$$\Delta \tau \approx \left[ \frac{5 \times 10^{13} \text{ g GeV}}{E_D} \right]^3 \frac{1}{3 \alpha(M)},$$

(27)

before the hole completely evaporates. The relevant emission period, $\Delta \tau$, is very sensitive to the detector threshold $E_D$. If $E_D = 10$ TeV, then $\Delta \tau \approx 64$ s, which in turn fixes the black hole temperature and $Q$ for this detector. The hole emits a total flux of about $5.7 \times 10^{27}$ photons above threshold in the last 64 s [see Eq. (19)], although the exact number will depend on the correct form of the fragmentation function as $E_\gamma/E_{\text{jet}} \rightarrow 1$. Of this flux 77% is emitted in the last second. If we
require a signal of $n_\gamma$ photons within $\Delta \tau = 1$ s, and within the detector resolution, the depth of sky scanned by the detector is

$$
    r_\gamma \simeq 0.035 \sqrt[3]{\frac{A}{10^8 \text{ cm}^2}} \frac{3}{n_\gamma} \text{ pc}.
$$

(28)

For the above values, $r_\gamma = 0.035$ pc giving a scanned volume of $V_\gamma = \omega A r^3 / 3 = 1.1 \times 10^{-5}$ pc$^3$.

The relevant background for our search comes from hadronic cosmic rays with energy close to $E_D$. The background rate of cosmic ray showers is

$$
    \frac{d^2 N_{cr}}{dA \, d\omega \, d\nu} (> E_D) = 0.21 \left[ \frac{E}{1 \text{ GeV}} \right]^{-1.5} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}.
$$

(29)

For $E_D = 10$ TeV and $\omega_D = 10^{-3}$ sr, the telescope would see a background of $0.02$ s$^{-1}$ per resolution element. If the rejection ratio is $10^{-3}$, the minimum detectable photon event rate is then $2 \times 10^{-5}$ s$^{-1}$ per resolution element. For $n_\gamma = 3$ and $\Delta \tau = 1$ s, the chance of cosmic rays mimicking individual PBH emission from any direction is $< 10^{-4}$ yr$^{-1}$. If no bursts are detected in one year of operation we could set a limit of 4.3 events yr$^{-1}$ at the 99% confidence level in $V_\gamma$. This would lead to an upper limit on the PBH explosion density of $dn/dt \lesssim 3.9 \times 10^3$ yr$^{-1}$ pc$^{-3}$. Practical considerations may be the deadtime of the telescope and modifications needed to efficiently register candidate events.

An alternative approach is to use atmospheric Čerenkov telescopes. Unlike air shower arrays, these have considerably lower $E_D$, much smaller fields of view and must be operated under clear dark skies. Recently it has been shown that sophisticated detectors can reject background hadronic cosmic rays with 99% efficiency using differences in the Čerenkov images from hadronic and photonic air showers. To evaluate the PBH detection efficiency we consider the GRANITE telescope, now under construction at the Whipple Observatory. It consists of two 10 m aperture optical concentrators, each fitted with a 109 pixel electronic camera, and separated by 120 m. The telescope is triggered by coincidences between corresponding pixels in the two concentrators. The energy threshold for gamma-ray shower detection is 100 GeV; the full field is 2.5 degrees ($\omega_A = 1.4 \times 10^{-3}$ sr); the angular resolution is of order the pixel size 0.25 degrees; and the collection area is $A_D = 5 \times 10^8$ cm$^2$. 

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Lower thresholds imply higher detectable fluxes due to fragmentation, and hence larger probed depths. This compensates for the reduced field of view so that their sensitivity is comparable to extensive air shower arrays. Table 1 summarizes the flux and time-integrated flux, above \( E_D \), as estimated from Eqs. (17) and (19). For low zenith angles and \( \Delta \tau = 1 \text{ s} \), the probed depth is \( r_\gamma \approx 0.48 \text{ pc} \) and the sensitive volume is \( V_\gamma \approx 5.2 \times 10^{-5} \text{ pc}^3 \). The sensitivity can be further improved by exploiting the unusual zenith angle variation of Čerenkov telescopes. Simulations show that as zenith angle increases, \( E_D \) and \( A_D \) increase by a factor \( f \). At a zenith angle of 75 degrees, we have \( f = 100 \). In this case, \( A_D = 5 \times 10^{10} \text{ cm}^2 \) and \( E_D = 10 \text{ TeV} \), which gives for \( n_\gamma = 4 \), \( r_\gamma \approx 0.72 \text{ pc} \) and \( V_\gamma \approx 1.7 \times 10^{-4} \text{ pc}^3 \).

<table>
<thead>
<tr>
<th>( \Delta \tau )</th>
<th>( 10^6 \text{ s} )</th>
<th>( 10^3 \text{ s} )</th>
<th>( 1 \text{ s} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( N_\gamma[&gt; E_D, \tau &lt; \Delta \tau] )</td>
<td>( 4.4 \times 10^{31} )</td>
<td>( 4.8 \times 10^{30} )</td>
<td>( 2.2 \times 10^{29} )</td>
</tr>
<tr>
<td>( \frac{dN_\gamma[&gt; E_D]}{dt} )</td>
<td>( 1.0 \times 10^{25} \text{ s}^{-1} )</td>
<td>( 2.5 \times 10^{27} \text{ s}^{-1} )</td>
<td>( 1.3 \times 10^{29} \text{ s}^{-1} )</td>
</tr>
</tbody>
</table>

The background of cosmic ray events comes from (i) distant proton showers of similar energies which arrive nearly parallel to the telescope’s optical axis and (ii) more local showers which traverse both telescopes’ fields of view and radiate into the cone of the cameras. The former is estimated at 0.0015 per second per pixel for a rejection ratio of \( 10^{-2} \). The latter can be rejected by examining the images. Observing time is limited since the low zenith angle precludes the accumulation of data bases on discrete sources, which usually must be as close as possible to transit. It is possible, though, to conduct PBH searches under less than optimum conditions. An observing time of 1000 hours in a dedicated experiment and \( n_\gamma = 3 \) will give a chance of mimicking a PBH as \( < 0.01 \). If no events are found in this time, then an upper limit of \( \lesssim 1.3 \times 10^5 \text{ PBH explosions yr}^{-1} \text{ pc}^3 \) for \( V_\gamma = 3 \times 10^{-4} \text{ pc}^3 \) can be inferred. This is weaker than the limit obtained from the MeV bound assuming maximum clustering [Eq. (15)]. Hence no improvement is offered over the
100 MeV diffuse bound by the present generation of atmospheric Čerenkov or air shower array experiments, assuming the Standard Model for PBH evaporation.

VI. MODELS WITH DIFFERENT PARTICLE SPECTRA

As mentioned in the introduction, we now illustrate the dependence of these results by taking a model with an exponentially increasing number of degrees of freedom. In the Hagedorn model, $6.0 \times 10^{34}$ erg of energy are evaporated in a final burst lasting $10^{-7}$ s. A significant fraction of the emission will be $\gamma$ rays of average energy 250 MeV. In this case the volume scanned by a dedicated telescope is large enough to significantly decrease the MeV bound on $\Omega_{pbh}$. The most sensitive detectors are atmospheric Čerenkov telescopes which detect bursts as if they were large air showers. A burst of 250 MeV gamma rays strikes the top of the atmosphere like a plane wave and interacts to produce secondary electrons. These in turn cause the atmosphere to radiate Čerenkov light which reaches the detector. Several independent estimates suggest that each vertically incident 250 MeV gamma ray produces on average about 2,300 optical photons reaching detector level.\textsuperscript{30,35} Although individual photons cannot be detected, the combined effect is equivalent to a giant shower of 100 ns duration. The challenge is to design a detector which can detect the weakest possible gamma ray flux (and hence be sensitive to the greatest possible depth) while at the same time maximizing the field of view (and hence the volume $V_\gamma$). The accidental trigger rate from background events (local air showers, night-sky fluctuations, man-made sources, etc.) must be minimized to get maximum exposure time under clear dark night-sky conditions. In general these conditions are not found in conventional atmospheric Čerenkov experiments and the most sensitive measurements require specially designed experiments.

The most sensitive direct limit on the explosion rate (Table 2) was set by the operation of two large optical reflectors in coincidence in 1976.\textsuperscript{35} In this experiment a cluster of phototubes was operated in coincidence on a 9 m optical reflector at the White Sands Missile Range and on a 10 m reflector at the Whipple Observatory 400 km away. Both reflectors were pointed at the same direction under clear skies for 22.5 hours; no coincidences were seen. The threshold for the detection of an event was 26 optical photons m$^{-2}$, giving a burst sensitivity of $5.6 \times 10^{-10}$ erg cm$^{-2}$. For $E_B \simeq 10^{34}$ ergs, this gave a maximum detectable distance of 390 pc, a sensitive volume of $V_\gamma = 4.4 \times 10^4$ pc$^3$ and a limit on Hagedorn-type PBH bursts of $< 0.04$ explosions pc$^{-3}$ yr$^{-1}$. 

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Greater sensitivity could be achieved using the two 11 m aperture Solar concentrators at Sandia Labs, Albuquerque \cite{36} together with the GRANITE telescopes (2 × 10 m) \cite{33} now under construction at the Whipple Observatory in Arizona. Each reflector would be equipped with an array of phototubes to give a full field of view of 3°. A local coincidence would be demanded between each set of telescopes which would then be put in coincidence (via telephone link) with each other. The 100 ns duration of the pulses would be verified using waveform digitizers at each telescope (eliminating possible spurious events). With the increased collection area and a more sophisticated trigger, the threshold per burst could be reduced to 1.8 × 10^{-10} \text{ erg cm}^{-2}, increasing \( V_\gamma \) by a factor of eight. The exposure time could also be increased to 450 hours in six months of operation. A Hagedorn PBH limit could then be as low as 2.5 × 10^{-4} explosions pc^{-3} yr^{-1} might then be obtained.

**VII. RESULTS**

Within the context of the Standard Model of quarks and leptons, we have discussed the signatures of black hole evaporations and compared them to those of the Hagedorn model. We reviewed the bound on the PBH abundance from MeV data; related it to bounds on the present rate of PBH explosions; and subsequently argued the possibility for improving the MeV bound by exploiting the new generation of TeV and PeV telescopes. A search for high energy (\( \sim 10 \text{ TeV} \)) emission from exploding holes may improve the bound only if arrays of imaging atmospheric Čerenkov telescopes or large arrays with muon detection capabilities and low energy thresholds are built (at high elevations). No existing array is presently suitable. The large uncertainty in the degree to which the PBHs are clustered in the Galaxy, however, will always make the comparison between bounds from the diffuse background and direct searches ambiguous.

In Hagedorn-type models where the particle spectrum rises exponentially with the mass of the particle, existing Čerenkov telescopes with good time resolution and low energy threshold can scan a sensitive volume three orders of magnitude larger than that excluded by the MeV bound in the Standard Model with maximum clustering. The Hagedorn model of PBH evaporation, however, seems unlikely given our present understanding of QCD physics.
ACKNOWLEDGEMENTS

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REFERENCES

FIGURE CAPTIONS

1) The parameter $\alpha$ counting the degrees of freedom of the Hawking mass, radiation as a function of the black hole temperature and lifetime. $\Lambda_{qh}$ is the quark-hadron deconfinement scale.

2) Particle fluxes radiated by a black hole of temperature $T = 10$ TeV.

3) Critical mass as a function of deconfinement scale $\Lambda_{qh}$.

4) Hawking radiation of MeV-photons for the standard model of quarks and leptons and four fundamental forces. Shown as a solid line is the calculation for $\Omega_{p\bar{p}h} = 7.6 \times 10^{-9}$ from bound in Eq. (10). Any concentration in excess of this limit would exceed the observed diffuse $\gamma$-ray background. The hatched areas correspond to $1\sigma$ errors.

5) The nearest distance at which a black hole of temperature $> T$ is expected for the density allowed by the MeV bound in Fig. 4. This distance is shown by the straight lines for three values of the clustering factor $\zeta$. The other lines are the distance at which the flux (integrated above detector energy threshold shown) from the nearest black hole can be resolved from the diffuse background flux. The detector is assumed to have an angular resolution $\omega_D = 10^{-3}$ sr and the energy thresholds shown.


The predicted fluxes of TeV and PeV cosmic photons are smaller than the cosmic ray flux in the corresponding energy range. The photons are absorbed by the atmosphere and detectors of order 1 m² carried by rockets or satellites are barely sensitive to the cosmic ray background. Fortunately the same atmosphere which shields the cosmic rays from Earth constitutes a calorimeter medium that can be exploited to detect the very low flux of high energy cosmic photons. The atmosphere represents a layer of material with a column density of 1000 g cm⁻² corresponding to roughly 25 radiation lengths. The primary photons initiate showers which dissipate their energy into a vast number of electrons and photons which are absorbed by the atmosphere at a height of about 10 km. Only the energy dissipated in weakly interacting neutrinos and muons, which are the decay products of pions photoproduced in the cascades, penetrate to observation level on a mountain top or at sea level. Muons lose 2 MeV of their energy through ionization for every g cm⁻² of matter traversed. To penetrate 1000 g of atmosphere they must have started with at least 2 GeV energy and a parent primary must have an energy of at least 10 GeV in order to be detected via sea level muons.

The showers initiated by cosmic photons with energies exceeding 100 TeV have a lateral spread of order 0.1-1 km because the secondary particles diverge as a result of their transverse momentum and secondary interactions in the atmosphere. An "extensive air shower array" of say 100 m² can detect showers over an effective area of 10 km², a gain in sensitivity of 10⁵, because only a fraction of the shower must traverse the array of detectors in order to be detected. The trajectory of the primary particle must not be contained inside the array. The array itself can consist of relatively widely separated particle detectors (usually a scintillator viewed by a phototube) sampling the 100 m² area. Enough particles survive to mountain (sea level) altitudes to thus detect showers of 10 (100) TeV primary energy.

Another technique consists of using mirrors viewed by phototubes collecting the Čerenkov light produced by shower particles in the atmosphere. Showers with energies as low as 0.1 TeV can be detected using this technique. A 100 GeV shower still produces several hundred electrons at an altitude of 10 km placed on a mountain. A mirror with 2 degrees aperture views showers over an effective area of 10⁵ m² at an altitude of 10 km. Recent experiments have been successful in distinguishing the gamma ray showers from the much more numerous proton
showers on the basis of the size of the Čerenkov light image.

At higher energies a popular technique is to simultaneously measure the electromagnetic and muon component in a single shower. The muon detector can be a GeV-muon detector at or near the surface or deep underground TeV-muon detector triggered by the air shower array. The simultaneous determination of the number of electrons and muons in the cascade could give a better handle on the determination of the nature of the primary particle, i.e., whether it is a proton or a heavy nucleus. Another possibility to be explored with such a hybrid detector is to do gamma ray astronomy in the presence of cosmic ray backgrounds by selecting showers poor in muons. Theoretical considerations lead us to expect that the number of muons in a photon shower is less than 5% of the number in a proton shower with similar energy or similar number of electrons. Muon-poor astronomy therefore requires measurement of the muon and electron number in each individual shower.
BOX B
Previous searches for PBH gamma-ray bursts

Several direct experimental searches for PBH explosions were made in the late seventies following the theoretical discovery of PBH evaporation. In general these involved the use of ground-based gamma-ray telescopes and used both archival data and observations made with specially configured detectors. No statistically significant detection was reported but upper limits were derived that placed limits on the PBH density.

Two extreme models for the final stages of the evaporation were considered, each predicting quite different parameters for the gamma-ray burst. The experiments were confined to searches for bursts with these characteristics. The Standard Model (also known as the Elementary Particle Model or EPM) predicted an outburst of $10^{30}$ ergs in gamma rays of energy 5 TeV, lasting 0.1 s. The Hagedorn Model, which corresponds to the exponential model discussed in the text, predicted an outburst of $10^{34}$ ergs in 250 MeV gamma-rays, lasting 100 ns.

In addition it was postulated that under certain conditions the exploding PBH might produce detectable radio or optical emission. In this scenario, the final $e^+e^-$ emission from the PBH is braked by an ambient magnetic field as the particle shell expands relativistically away from the hole, thus producing an electromagnetic pulse. The pulse would only be appreciable if the explosion is governed by the Hagedorn model — a conclusion strengthened by the most recent analysis including quark and gluon emission. Because of the greater sensitivity of radio and optical techniques, it was possible to probe greater depths with these experiments. However, a null result in these searches had less force since the conditions for pulse production are highly model dependent.

In Table 2 we list the principal searches that have been made, the techniques and assumptions used and the limits derived on the PBH density. The latter show a spread of $10^{14}$. 

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Table 1
Observations limits on PBH densities.

<table>
<thead>
<tr>
<th>Technique</th>
<th>Group</th>
<th>Assumptions</th>
<th>Density event- pc^{-3} . yr^{-1}</th>
<th>Reference</th>
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<tbody>
<tr>
<td>ACT</td>
<td>3 x 1.5 m</td>
<td>SAO-UCD</td>
<td>HM</td>
<td>&lt; 0.47</td>
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<tr>
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<td>HM</td>
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</tr>
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<td>HM</td>
<td>&lt; 2.1</td>
</tr>
<tr>
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<td>SAO-UCD</td>
<td>HM</td>
<td>&lt; 0.04</td>
</tr>
<tr>
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<td>SAO-UCD</td>
<td>EPM</td>
<td>&lt; 7 x 10^5</td>
</tr>
<tr>
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<td>10 m</td>
<td>SAO-UCD</td>
<td>EPM</td>
<td>&lt; 8 x 10^5</td>
</tr>
<tr>
<td>ACT</td>
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<td>SAO-UCD</td>
<td>EPM</td>
<td>&lt; 3 x 10^4</td>
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<tr>
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<td>SAO-UCD</td>
<td>EPM</td>
<td>&lt; 2 x 10^4</td>
</tr>
<tr>
<td>EAS</td>
<td>2 arrays</td>
<td>UCC-UCD</td>
<td>EPM + UHE emission</td>
<td>&lt; 2 x 10^2</td>
</tr>
<tr>
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<td>1 array</td>
<td>Tata</td>
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<td>&lt; 3 x 10^3</td>
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<td>EPM + optical emiss.</td>
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<tr>
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<td>U. Mass</td>
<td>EPM + radio emission</td>
<td>&lt; 7 x 10^{-10}</td>
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References


Also: