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Gamma-ray bursts from the final stage of primordial black hole evaporation

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ABSTRACT

It is now accepted that, within the Standard Model of particles, evaporating primordial black holes cannot produce detectable gamma-ray bursts because the expected photon flux from black hole explosions is too weak, and consists mainly of GeV photons. Contrary to this verdict, we put forward a scenario in which a large fraction of black hole power is converted into photon luminosity in the MeV spectral range, producing a burst of duration 10^{-1}–10^{3} s. We show that when the black hole temperature exceeds ~ 10 GeV, the charged particle outflow from a black hole forms a well-defined plasma and the magnetohydrodynamical regime of expansion may be realized. In this case, the kinetic energy of particles may be converted into soft gamma-rays due to the synchrotron radiation and the electromagnetic cascade in the close-to-equipartition turbulent magnetic field. We show that some of the gamma-ray bursts detected by BATSE can be associated with evaporating black holes.

Key words: black hole physics – MHD – gamma-rays: bursts.

1 INTRODUCTION

The enigma of gamma-ray bursts (GRBs) has recently stimulated an interest in the fundamental problem of the detection of evaporating primordial black holes (PBHs). The idea of a possible connection between GRBs and the final stage of evaporating black holes was discussed (and rejected) in the pioneering works by Page & Hawking (1976), and recently the possibility of the detection of the final outbursts from evaporating PBHs was considered again in a number of papers (MacGibbon & Webber 1990; MacGibbon & Carr 1991; Halzen et al. 1991; Cline & Hong 1992; Semikoz 1994). The result of these works can be summarized as follows: (i) within the standard quantum chromodynamics (QCD), the individual PBH explosions *cannot* be detected by modern γ-ray telescopes since their expected rate is too small and the photon flux is too weak, even at the last moment of the PBH life (if we take into account the finite time resolution of the detector); (ii) the evaporating PBHs *cannot* be associated with the observed GRBs because the typical photon energies at the final stage of PBH evaporation lie in the range $\gtrsim 100$ MeV (MacGibbon & Webber 1990), in contrast to the observed GRBs in which the main energy release occurs in the 0.1–1 MeV range.

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Contrary to these conclusions, the main statement of this paper can be expressed as follows.

(1) There exists at least one mechanism which allows the conversion of a large fraction of the power of a PBH explosion into photon luminosity in the 'proper' GRB spectral range 0.1–1 MeV during the time interval 10^{-1}–10^{3} s. We show below that when the PBH temperature exceeds ~ 10 GeV, the charged particle outflow from a black hole forms a well-defined plasma. In this case the magnetohydrodynamical (MHD) regime may be realized in the expanding particle wind, and the kinetic energy of particles is converted into soft γ-rays due to the synchrotron radiation and the electromagnetic cascade in the close-to-equipartition turbulent magnetic field.

(2) *If* the significant fraction of the PBH energy at the final stage of its evaporation (when the PBH temperature exceeds ~ 1 TeV) is deposited in the 'proper' GRB range, then individual PBH explosions should already have been detected by BATSE. Therefore, at least some of the GRBs observed by BATSE could originate from evaporating black holes.

In Section 2 we briefly outline the physical picture of the final stage of PBH evaporation on the basis of the standard QCD approach. In Section 3 we discuss the generation of γ-rays in the MHD wind emanating from an evaporating black hole. The role of relativistic shock waves in the PBH evaporation outflow is considered in Section 4. In Section 5 we formulate the main predictions and several observational tests that allow us to check the validity of the model.

2 AN OUTLINE OF PBH EVAPORATION

Black holes of small masses may have been formed in the early Universe by several mechanisms, including initial density and gravitational-wave perturbations, and also by more exotic possibilities like phase transitions or the collapse of cosmic strings; see Polnarev & Khlopov (1985), MacGibbon & Carr (1991), Halzen et al. (1991) and references therein for the review.

According to Hawking's results (Hawking 1974), an isolated black hole radiates particles approximately like a blackbody of temperature

$$T \simeq \frac{10^{10}\,\mathrm{g}}{M}\ \mathrm{TeV}, \tag{1}$$

where M is the PBH mass. We will consider only non-rotating black holes since it was shown by Page (1976) that a rapidly rotating black hole loses its angular momentum long before most of its mass has been evaporated. A black hole of mass $M \gg 10^{17}$ g emits only massless particles. When M drops below this limit, a PBH begins emitting electron–positron (e^-e^+) pairs. When the PBH temperature (equation 1) exceeds the QCD scale, ~ 100 MeV, a PBH also radiates quarks and gluons, which subsequently develop into hadron jets. The evaporation process at this stage was investigated in detail by MacGibbon & Webber (1990). According to the results of their work, when the PBH temperature exceeds a few GeV, the total particle flux from the evaporating black hole is dominated by the products of jet fragmentation. The primary particles of energy $\sim 5\,T$ constitute a very small part of the outcoming flux. After the jet-fragmentation stage (at the comoving-frame distance $\sim 10^{-13}$ cm from the black hole), the particle flux consists mainly of pions, with a small admixture of other hadrons. Finally, pions and other unstable jet products decay into stable particles: photons, e^-e^+ and proton–antiproton ($p\bar{p}$) pairs, and neutrinos.

The PBH loses its mass at a rate (Page & Hawking 1976) of

$$\mathrm{d}M/\mathrm{d}t \simeq -8\times 10^{6}\alpha(T)\tilde{T}^{2}\ \mathrm{g\ s^{-1}}. \tag{2}$$

Here $\tilde{T}$ is temperature in TeV and $\alpha(T) \leq 1$ accounts for the particle degrees of freedom; $\alpha = 1$ at $T \gtrsim 0.1$ TeV. Combining (1) and (2), one obtains the time dependence of the PBH temperature

$$\tilde{T}(t) \simeq \frac{\tilde{T}_0}{[1 - t/\Delta\tau(T_0)]^{1/3}}. \tag{3}$$

Here T_0 is a given initial temperature at $t=0$, $\Delta\tau(T)$ the lifetime of a hole with temperature T before evaporation

$$\Delta\tau \simeq 5\times 10^{2}/\tilde{T}^{3}\ \mathrm{s}. \tag{4}$$

Black holes with temperature $T \sim 1$ TeV live 8 min; those with $T \sim 10$ GeV live 16 years. Black holes having at the moment of birth an initial mass $M \sim 5\times 10^{14}$ g and temperature $T \sim 20$ MeV have a lifetime equal to the age of the Universe and are just evaporating today.

The most reliable restriction on the number of individual PBH explosions comes from the observations of the γ-ray and cosmic ray backgrounds at around 0.05–1 GeV (MacGibbon & Carr 1991; Halzen et al. 1991). It was shown by MacGibbon & Carr (1991) that the PBH emission matches well the γ-ray background spectrum at $\sim$50–170 MeV. If the observed cosmic γ-ray background in this range is due to evaporating PBHs, the present rate of PBH explosions is

$$\mathrm{d}n/\mathrm{d}t \simeq 10\ \mathrm{pc^{-3}\ yr^{-1}}, \tag{5}$$

provided PBHs are clustered to the same degree as the other matter in the Galactic halo. The latter is a reasonable assumption since PBHs should not have large peculiar velocities. If the observed γ-ray background does not originate from PBH emission, the right-hand side (r.h.s) of equation (5) should be considered as an upper limit on the PBH explosion rate. The bounds on the PBH density from the observed charged-particle backgrounds between 0.1 and 1 GeV give an estimation similar to that in equation (5) (MacGibbon & Carr 1991).

3 MHD WIND SCENARIO OF PBH EVAPORATION

In this section we discuss the hypothesis that a large fraction of the PBH energy at the final stage of its evaporation is radiated in the 0.1–1 MeV range due to the electromagnetic cascade and the synchrotron emission of relativistic particles in the outflow. The crucial points in this hypothesis are the validity of the plasma description for the PBH evaporation wind and the existence of the close-to-equipartition turbulent magnetic field.

3.1 Validity of hydrodynamical approximation

As was mentioned above, after the hadronization occurs the outgoing particle wind consists mainly of quark- and gluon-jet products. These are dominated by pions, with practically equal numbers of π^+, π^- and π^0 species. After $\tau_0 \simeq 10^{-16}$ s, neutral pions decay into the γ-quanta of energies $\gtrsim 70$ MeV, and in $\tau_e \simeq 10^{-6}$ s charged pions decay finally (via muons) into e^-e^+ pairs and neutrinos. Of course, these are proper times in the pion and muon rest frames. The relevant particle fluxes and luminosities as functions of the PBH temperature were calculated by MacGibbon & Webber (1990) using the Monte Carlo approach. Though the calculations by MacGibbon & Webber (1990) are performed for PBH temperatures $T \leq 100$ GeV, the results for larger temperatures should not differ significantly. Moreover, within the standard QCD, the pion flux can be estimated analytically for any temperature T by convolving the spectral flux of primary particles with an empirical pion fragmentation function; see Halzen et al. (1991), Semikoz (1994), and estimations below. In any case, the possible appearance of unknown particle species at higher temperatures would only increase the luminosity of the black hole and of the subsequent γ-ray burst.

Using, e.g., the results of MacGibbon & Webber (1990), one can easily find that at all stages of jet fragmentation, particle–particle collisions cannot provide the hydrodynamical approximation. Nevertheless, the hydrodynamical regime *is* possible, if the particle wind forms a well-defined plasma in the usual sense that the local value of the Debye radius r_D at a given distance r from the black hole is less than r. In this case, plasma-wave turbulence may have time to develop during the expansion, and the hydrodynamical regime is supported in the collisionless plasma by the collective wave-particle interactions and/or by the turbulent magnetic field. This assumption usually works well in various astrophysical problems, such as the dynamics of stellar winds, collisionless shocks in the Solar wind, low-rate accretion, etc.

Let us show that an outflow emanating from a black hole becomes a plasma at the final stage of the evaporation, both at the level of the $\pi^-\pi^+$ wind and farther from the black hole, when pions and muons decay finally into e^-e^+ pairs.

Consider first the $\pi^-\pi^+$ wind. The pion flux for all species (π^-, π^+ and π^0) can be calculated by convolving the Hawking emission spectrum with the jet-fragmentation function (MacGibbon & Webber 1990):

$$\frac{d\dot{N}_\pi}{dE}=\sum_j \int_0^\infty \frac{d\dot{N}_j}{dQ}\frac{dM_\pi(Q,E)}{dE}\,dQ. \tag{6}$$

Here j labels the primary particle species,

$$\frac{d\dot{N}_j}{dQ}=\frac{\Gamma_j}{2\pi\hbar}\frac{1}{\exp(Q/T)-(-1)^{2s}} \tag{7}$$

is the emission rate per degree of freedom of primary particles with spin s in the energy interval (Q, $Q+dQ$), and $\Gamma_j(Q,\ T)$ is an absorption probability (MacGibbon & Webber 1990). The fragmentation function for pions gives the number of pions produced in the energy interval (E, $E+dE$) in the quark or gluon jet of energy Q. An empirical expression for the pion fragmentation function, valid above the rest mass threshold, is given by (Hill, Schramm & Walker 1987)

$$\frac{dM_\pi(Q,E)}{dE}=(15/16)x^{-3/2}(1-x)^2. \tag{8}$$

Here $x=E/Q$. An order-of-magnitude estimation of the total pion flux is

$$\dot{N}_\pi\sim k\dot{N}_{\rm prim}M_\pi(Q=5T)\simeq 3\times10^{29}\tilde{T}^{1.5}\ {\rm s}^{-1}, \tag{9}$$

where $\tilde{T}$ is the PBH temperature in TeV, $\dot{N}_{\rm prim}\simeq10^{25}\tilde{T}\ {\rm s}^{-1}$ the total flux of primary particles per degree of particle freedom (its value is very weakly dependent on s), and k is the total number of degrees of freedom (72 for quarks and 16 for gluons at $T\gtrsim100$ GeV). This simple estimation agrees with the Monte Carlo simulations by MacGibbon & Webber (1990) within a factor of 2. The flux of charged pions is roughly 2/3 of the value of equation (9).

The pions produced in the jet have a very broad energy distribution. By numerically integrating equation (6), we find that most of the number flux of pions is contained in particles with energies $E_\pi/m_\pi c^2-1\sim2$, where $m_\pi c^2\sim140$ MeV is the pion rest energy. At the same time, for sufficiently hot black holes the main *power* is carried by ultrarelativistic pions, and the mean energy of jet-produced pions is therefore $\bar{E}_\pi\sim20\sqrt{\tilde{T}}$ GeV (MacGibbon & Webber 1990). Note that we neglect inelastic collisions of pions (strong interaction) that may occur with non-negligible probability close to the radius where pions are created. Such collisions are unlikely to change considerably the energy distribution of the pions.

It can be shown that the Langmuir frequency in such a non-equilibrium relativistic $\pi^-\pi^+$ wind in the rest frame of the black hole (PBH frame) should be estimated using the standard relativistic expression

$$\omega_{L\pi}\simeq\sqrt{4\pi e^2 n_\pi c^2/E_\pi}\simeq2\times10^{12}\tilde{T}^{3/4}r^{-1}\ {\rm s}^{-1}, \tag{10}$$

where $E_\pi\sim3m_\pi c^2$ in our case, and

$$n_\pi\simeq\frac{\dot{N}_{\pi\pm}}{4\pi r^2 c} \tag{11}$$

is the density of charged pions, e the charge of an electron. Using equations (9)–(11), we compare the Debye radius $r_{D\pi}$ at a given distance r from the black hole with the characteristic size of the $\pi^-\pi^+$ plasma, r, and obtain that the inequality

$$r_{D\pi}\sim c/\omega_{L\pi}<r \tag{12}$$

is satisfied in a $\pi^-\pi^+$ wind when the PBH temperature exceeds ~5 GeV. This estimation is not very sensitive to the value of E_π. For example, even if we suppose that all pions have relativistic energies and replace $E_\pi\to\bar{E}_\pi$ in equation (10), the minimum PBH temperature satisfying the inequality (12) will be only ~20 GeV. The estimation of the Debye radius for the charged muon wind gives practically the same temperature threshold. One can also check that the inequality $r_D>n^{-1/3}$ is satisfied for pions and muons immediately after the PBH temperature exceeds the hadronic scale.

The e^-e^+ energy after the decay of pions and muons is roughly 1/5 of the $\pi^-\pi^+$ energy, while the e^-e^+ flux $\dot{N}_e$ is approximately equal to $\dot{N}_{\pi\pm}$. Therefore, the inequality (12) gives approximately the same lower limit on the PBH temperature, $T\sim3$–10 GeV. If the e^-e^+ pairs are efficiently cooled due to synchrotron losses (see Sections 3.2.1 and 3.2.2), the above temperature threshold will be considerably lower.

The resulting plasma outflow should be strongly non-equilibrium and anisotropic. This may naturally give rise to the development of various plasma instabilities and a high level of wave turbulence, thus establishing the fluid-like expansion. The fluid-like behaviour of an expanding wind can be supported also by the turbulent magnetic field which may exist in the plasma as a result of plasma-wave turbulence or because of the propagation of internal shocks (see the next section). In this case the Larmor radius r_B of charged particles plays the role of a mean free path. An equipartition magnetic field in the charged particle wind in the PBH frame is

$$B_{\rm eq}\simeq\left(\frac{8\pi\beta P_{\rm tot}}{4\pi r^2 c}\right)^{1/2}\simeq\frac{7\times10^8(\alpha\beta)^{1/2}\tilde{T}}{r}\ {\rm G}. \tag{13}$$

Here β is a fraction of the total PBH power,

$$P_{\rm tot} = -({\rm d}M/{\rm d}t)c^2 \simeq 7\times10^{27}\alpha(T)\tilde{T}^2 \text{ erg s}^{-1}, \tag{14}$$

that goes into charged particles at different stages of the wind. For a $\pi^-\pi^+$ wind, β is of order 2/3, and for the e^-e^+ stage of the wind $\beta\sim 1/7$ [unless synchrotron cooling takes place or an electromagnetic cascade is developed (see below)]. If the energy density of a magnetic field is a fraction λ of the equipartition value, $B=\lambda^{1/2}B_{\rm eq}$, the inequality $r_{B\pi}=E_\pi/eB<r$ is satisfied for pions with Lorentz factors $\gamma_\pi \lesssim 10^3\lambda^{1/2}\tilde{T}$, i.e., practically for all particles including those with energies $\bar{E}_\pi$ and corresponding Lorentz factors $\bar{\gamma}_\pi\sim 10^2\tilde{T}^{1/2}$ if $\lambda\tilde{T}\gtrsim 2\times10^{-2}$. Here γ_π is the pion Lorentz factor. For muon and e^-e^+ winds the last requirement is even less stringent.

Having in mind the above considerations, we will assume that the hydrodynamical regime of the charged particle outflow is realized, at least, for the PBH temperatures of interest to us (0.5–10 TeV). It is this range that corresponds to the observed GRB duration (see equation 4) and power (see Section 5.1 below).

3.2 Soft gamma-ray luminosity of the evaporation wind

The efficiency of the conversion of particle kinetic energy into radiation in the 0.1–1 MeV range depends crucially on whether or not the energy density of the magnetic field can be close to equipartition with the pressure of a charged particle outflow. Suppose that the energy density of a magnetic field is a fraction λ of the equipartition value determined by equation (13) (with $\alpha=1$). Generally speaking, λ is a slowly varying function of r.

3.2.1 Electromagnetic cascade initiated by charged pions

For $\lambda=$constant, the synchrotron losses, ${\rm d}\gamma_\pi/{\rm d}r=-\gamma_\pi/l_\pi(r)$, by a charged pion propagating in a radial direction in a chaotic magnetic field of strength $B=\lambda^{1/2}B_{\rm eq}\ll m_\pi^2c^3/e\hbar=3\times10^{18}$ G are characterized by the pion free path

$$l_\pi \simeq 5\times10^{26}/\gamma_\pi B^2 \text{ cm} \propto r^2/\lambda. \tag{15}$$

This gives the following radial dependence of a pion Lorentz factor:

$$\gamma_\pi(r)\simeq\frac{\gamma_{\rm i}}{1+(r_{\rm syn}/r_{\rm i})(1-r_{\rm i}/r)}; \qquad r_{\rm syn}\simeq 10^{-9}\lambda\beta\tilde{T}^2\gamma_{\rm i} \text{ cm}. \tag{16}$$

Here $\gamma_{\rm i}\equiv\gamma_\pi(r_{\rm i})$ is a pion Lorentz factor at some initial distance $r_{\rm i}$ from the black hole, $r_0<r_{\rm i}<r_{\rm d}$; r_0 and $r_{\rm d}$ are the distances where the pions are created and decay into muons: $r_0\sim10^{-13}\gamma_\pi(r_0)$ cm, and $r_{\rm d}\simeq10^3\gamma_\pi$ cm. It is clear from (16) that when $r_{\rm i}\ll r_{\rm syn}$, the charged pions should radiate all their energy during a much shorter time than the expansion time $r_{\rm i}/c$. Note that for $\lambda\beta\tilde{T}^2\sim1$ and $\gamma_{\rm i}\sim\gamma_\pi(r_0)$ the value of $r_{\rm syn}$ is still four orders of magnitude greater than r_0. For pions with mean energy $\bar{\gamma}_\pi=\bar{E}_\pi/m_\pi c^2\sim10^2\tilde{T}^{1/2}$, we have

$$\bar{r}_{\rm syn}\simeq10^{-7}\lambda\beta\tilde{T}^{5/2} \text{ cm}. \tag{17}$$

Now let us take into account the radial dependence of $\lambda(r)$. One may expect that the function $\lambda(r)$ grows monotonically with r, i.e. along the flow lines, because of the convective development of plasma-wave turbulence in the expanding wind. The minimum distance $r_{\rm i}$, at which the charged pions can radiate their energy to synchrotron emission during a shorter time than the expansion time $r_{\rm i}/c$, is defined by an implicit relation

$$l_\pi(r_{\rm i})=r_{\rm i}. \tag{18}$$

Since the dependence $\lambda(r)$ is unknown, we cannot find an exact solution for pion energy losses, as in equation (16). However, due to the smoothness of $\lambda(r)$ one may expect that the solution to equation (18) gives the characteristic distance $r_{\rm i}$ similar to that in equation (16): $r_{\rm i}=r_{\rm syn}=10^{-9}\lambda(r_{\rm i})\beta\tilde{T}^2\gamma_\pi$ cm. We will assume hereafter that $\gamma_\pi=\bar{\gamma}_\pi$, so that $r_{\rm i}=\bar{r}_{\rm syn}$. For this distance to be much greater than r_0, the value of λ should be not too small: $\beta\lambda(r_{\rm i})\gg10^{-4}\tilde{T}^{-2}$. In this case the charged pions will radiate nearly all their kinetic energy to synchrotron emission at the distance $r=r_{\rm i}+\Delta r$, where $\Delta r\lesssim r_{\rm i}$.

The characteristic energies of synchrotron photons are (Ginzburg & Syrovatskii 1964)

$$E_{\rm syn}\simeq0.4\hbar\omega_{B\pi}\bar{\gamma}_\pi^2\simeq1[\lambda(r_{\rm i})\beta]^{-1/2}\tilde{T}^{-1/2} \text{ GeV}. \tag{19}$$

Here $\omega_{B\pi}=eB/m_\pi c$ is a pion gyrofrequency. For $r\sim\bar{r}_{\rm syn}$ the magnetic field is very high: $B(r_{\rm i})/B_{\rm cr}\sim10^2(\lambda\beta)^{-1/2}\tilde{T}^{-3/2}$, where the critical field for electrons is $B_{\rm cr}=m_e^2c^3/e\hbar\simeq 4\times10^{13}$ G; m_e is an electron rest mass. Therefore, the synchrotron photons have energies well beyond the one-photon e^-e^+ pair-production threshold and initiate the electromagnetic cascade: photons create e^-e^+ pairs, and electrons (positrons) radiate photons in a strong magnetic field. The kinetic theory of electromagnetic showers in strong magnetic fields of the order of critical value $B_{\rm cr}$ was derived by Akhiezer, Merenkov & Rekalo (1995) in the spirit of the well-known theory of electromagnetic showers in matter. The cascade develops if the energy, E_0, of an initial particle (photon or electron) satisfies the condition $\xi=(B/B_{\rm cr})(E_0/m_ec^2)\gg1$.

It is clear that the magnetic field should have a maximum somewhere at $r_{\rm m}>r_{\rm i}$, because $B\simeq0$ at $r<r_0$, and $B\propto r^{-2}$ for large radii. For the photons of energy (19), and for the maximum magnetic field $B(r_{\rm m})$ the necessary condition for the cascade may be rewritten as

$$\xi\simeq3\times10^5b\lambda^{-1}(r_{\rm i})\beta^{-1}\tilde{T}^{-2}\gg1. \tag{20}$$

Here the factor $b=(r_{\rm i}/r_{\rm m})[\lambda(r_{\rm m})/\lambda(r_{\rm i})]^{1/2}\geq1$, and we do not expect it to be very much greater than unity. The requirement (20) together with inequality $r_0\ll r_{\rm i}$ gives the following interval of parameters,

$$10^{-4}\ll\lambda(r_{\rm i})\beta\tilde{T}^2\ll3\times10^5b. \tag{21}$$

The mean free path l of a photon of energy $\xi\gg1$ with respect to one-photon pair production is given by (Akhiezer et al. 1995)

$$l\simeq6\times10^{-8}(B_{\rm cr}/B)\,\xi^{1/3} \text{ cm}. \tag{22}$$

Under the same condition $\xi\gg1$, the mean free path $l_{\rm e}$ of an electron before emitting a synchrotron photon is given also by (22), with numerical factor approximately two times less than in (22). As a crude approximation, we may assume that at each step of the cascade the energy of the particles is divided by two. The cascade develops until the photon energies degrade to the threshold value

$$E_c = \max[m_e c^2 (B_{cr}/B),\ 2m_e c^2]. \quad (23)$$

Therefore, the number of steps in the cascade is roughly $q \simeq \log(E_0/E_c)/\log 2$, which is of the order of 10 if $E_0 = E_{syn}$ and $E_c \sim 1$ MeV (we will call the latter case a *complete* cascade). A photon of energy (19) will initiate a cascade if the inequality $ql \lesssim r_m$ is fulfilled at the radius r_m. Using (19), (20), the last inequality, can be rewritten as

$$[r_i \lambda(r_i)/r_m \lambda(r_m)]^{2/5} \lesssim \lambda(r_i)\beta \tilde{T}^2, \quad (24)$$

where the quantity on the left-hand side (l.h.s.) of (24) is less than or of the order of unity.

More precise (and weaker) restrictions on the magnetic field and PBH temperature follow from the integration of the inverse free path (22) over the radius, from r_i where $B(r_i) \gg B_{cr}$ to $r_c \simeq 4 \times 10^{-5} [\lambda(r_c)\beta]^{1/2} \tilde{T}$ where $B(r_c) \simeq B_{cr}/2$, assuming $\lambda \simeq \lambda(r_c)$. This gives the necessary condition for the complete cascade,

$$10^{-3} q^{3/2} \tilde{T}^{1/4} \lesssim \lambda(r_c)\beta \tilde{T}^2, \quad (25)$$

which should be fulfilled together with (21). When $\tilde{T} \sim 1$, the inequality (25) requires that the values of the magnetic field should not be too far below equipartition, and for the minimal value $\lambda \sim 1/30$, the cascade starts at $r_i \sim 3 \times 10^{-9}$ cm and is terminated at $r_c \sim 6 \times 10^{-6}$ cm. For lower values of $\tilde{T}$ and λ, when $\lambda \tilde{T}^2 \ll 1/30$, the cascade may be 'incomplete': it terminates at high energies $E_c \gg 1$ MeV. Note that due to the growth of the PBH temperature with time, the final energies of the cascade photons decrease with time from GeV values when $\tilde{T} \ll 1$ to ~ 1 MeV in the complete cascade. Therefore, the cascade gamma-ray spectrum should exhibit a hard-to-soft evolution with a time-scale defined by equation (4).

After the complete cascade, the particle outflow consists mainly of sub-MeV photons and e^-e^+ pairs that share the initial $\pi^-\pi^+$ energy, non-relativistic charged pions, and neutral pions that are left undisturbed by the cascade. The magnetic field in the outflow is supported by e^-e^+ pairs, and its equipartition value is given again by (13) with $\beta \sim 0.2$–0.3. After e^-e^+ pairs have been cooled and annihilated (see below), the magnetic field is supported mainly by non-relativistic charged pions, and its equipartition value (13) decreases, with $\beta \rightarrow 10^{-2}$.

The kinetic theory of the cascade is valid if one may neglect the two-photon pair production, Compton scattering, and also possible collective plasma effects (Debye screening of gamma-rays by over-dense plasma). Neglecting the two-photon processes is justified if the particle distributions are very anisotropic (beamed in radial direction). However, at the end of the cascade the particle distributions are expected to be more isotropic, so the inner region of the outflow, $r \lesssim 4 \times 10^{-3} \beta \tilde{T}^2$ cm, may be very opaque. Furthermore, due to a highly increased number of e^-e^+ pairs, the sub-GeV photons will be screened by the over-dense post-cascade plasma. Therefore, starting from the radius where the cascade is terminated, i.e. from r_c or even earlier, the created cloud of photons and e^-e^+ pairs is expected to expand quasi-adiabatically, as described by Paczyński (1986), until the comoving temperature becomes non-relativistic and e^-e^+ pairs annihilate. The significant difference from the situation considered by Paczyński (1986) and subsequent works on e^-e^+ fireballs is that e^-e^+ pairs should strongly suffer from radiation losses due to synchrotron and magnetic bremsstrahlung emission. Also, the synchrotron reabsorption may be important, even for high-energy photons produced in an incomplete cascade. As a result, an equivalent temperature of the emergent radiation should fall into the 0.1–1 MeV range. The shape of the radiation spectrum is affected by many factors, like radiation losses of e^-e^+ pairs due to magnetic bremsstrahlung emission that give rise to the relatively soft component of the burst spectrum with a high-energy break near 0.5 MeV, or non-equilibrium energy distributions of cascade-produced photons and e^-e^+ pairs. The detailed analysis of this problem will be presented elsewhere.

3.2.2 Electromagnetic cascades from π^0 decays

The second possibility of converting PBH energy into soft γ-rays is related to the electromagnetic cascade initiated by π^0 decays. Suppose that the radius

$$r_{ph} = c\tau_0 \gamma_\pi \simeq 3 \times 10^{-6} \gamma_\pi \text{ cm}, \quad (26)$$

where the neutral pions with Lorentz factor γ_π decay into photons, the energy density of the magnetic field is a fraction λ of the equipartition value (13): $B(r_{ph}) \simeq 2 \times 10^{14} \lambda^{1/2} \tilde{T}/\gamma_\pi$ G. Here we put $\beta = 2/3$ for definiteness, assuming an equipartition with charged pions and neglecting their synchrotron cooling. For $\gamma_\pi = \bar{E}_\pi/m_\pi c^2 \sim 10^2 \tilde{T}^{1/2}$ we have $\bar{r}_{ph} \simeq 4 \times 10^{-4} \tilde{T}^{1/2}$ cm. The magnetic field at this radius is much less than the critical value, even for $\lambda \sim 1$ and $\tilde{T} \sim 10$, e.g.

$$B(\bar{r}_{ph}) \simeq 10^{12} \lambda^{1/2} \tilde{T}^{1/2} \text{ G}. \quad (27)$$

The pion-produced photons have energies of order $E_{ph} \sim E_\pi/2$. Their mean energies are $\bar{E}_\pi/2 \sim 10 \tilde{T}^{1/2}$ GeV (MacGibbon & Webber 1990). The cascade parameter $\xi = (B/B_{cr})(E_{ph}/m_e c^2)$ for pion-produced photons is

$$\xi(r_{ph}) \simeq 7 \times 10^2 \lambda^{1/2} \tilde{T}, \quad (28)$$

and does not depend on the π^0 Lorentz factor. Therefore, in the case $\lambda^{1/2}\tilde{T} \gg 2 \times 10^{-3}$, when $\xi(r_{ph}) \gg 1$, the situation is similar to that described above for the $\pi^-\pi^+$ wind: the photons produced by π_0 decays initiate a cascade. This cascade is incomplete if $B(r_{ph}) \lesssim B_{cr}/2$, i.e. if $\lambda \tilde{T} \lesssim 200$. The maximum possible number of steps, q, in the cascade is reached for $ql \lesssim r_{ph}$, where l is defined by (22). The last inequality is satisfied for $\lambda^{1/2} \tilde{T}^{4/3} \gtrsim q/40$. In this case, for $E_\pi = \bar{E}_\pi$, the number of steps $\bar{q} = \log \xi(\bar{r}_{ph})/\log 2 \sim 10$, and the final energies of particles, produced in the cascade, are given by

$$E_c(\text{MeV}) \sim B_{cr}/2B(\bar{r}_{ph}) \simeq 15 \lambda^{-1/2} \tilde{T}^{-1/2}. \quad (29)$$

The final e^-e^+ pairs of energies $\gamma_c \sim E_c/m_e c^2 \gg 1$ emit all their kinetic energy to the synchrotron radiation, and the corresponding break in the synchrotron spectrum is at the energy

$$E_b \sim 0.4 \hbar \omega_{Be} \gamma_c^2 \sim 5 \lambda^{-1/2} \tilde{T}^{-1/2} \text{ MeV}. \quad (30)$$

Here $\omega_{Be} = eB/mc$ is an electron gyrofrequency at $r = \bar{r}_{ph}$. The resulting spectrum consists of cascade-produced photons of energies $\sim E_c$, and of the synchrotron photons with energies $\lesssim E_b$. Note that the values of E_b and E_c shift to lower energies with increasing temperature. Therefore, the spectrum should exhibit a hard-to-soft evolution with time.

The created cloud of MeV synchrotron photons and e^-e^+ pairs may be marginally optically thick with respect to two-photon pair production and synchrotron reabsorption. If we assume that at the radius $r \sim \bar{r}_{ph}$ roughly half the energy of initial neutral pions is converted into the γ-quanta in the energy range 0.5 MeV $\lesssim E \lesssim E_b$, an optical depth for two-photon pair production is

$$\tau_{\gamma\gamma} \sim (3/8)\,\sigma_T n_{syn} r \sim 0.5\,\lambda^{1/2}\tilde{T}^{3/2}, \tag{31}$$

where n_{syn} is the density of synchrotron photons, σ_T the Thomson cross-section. An opacity of the e^-e^+ pair wind as a result of two-photon pair production and synchrotron reabsorption may lead to the cooling of cascade-produced and synchrotron photons with energies E_b and E_c, and to further softening of the spectrum.

The above consideration means that when the PBH temperature exceeds ~ 0.3 TeV, the power of relativistic pions, which is roughly 30–90 per cent of the total PBH power, is converted into soft γ-rays, while the initial π^0-decay emission flux that dominated the spectrum in the kinetic regime (MacGibbon & Webber 1990) is suppressed by a factor $\sim \exp(-r_{ph}/l) \sim 10^{-2}$–$10^{-3}$ if $\lambda^{1/2}\tilde{T}^{4/3} \gtrsim q/40$. The directly emitted primary photons and e^-e^+ pairs of energies $\sim 5T$, which carry away several per cent of the total PBH power, should have a similar fate, generating an e^-e^+ pair cascade at a point where the cascade parameter ξ becomes greater than 1 and the cascade length is less than the distance r. Thus, a single assumption of a close-to-equipartition magnetic field in a plasma outflow formed by PBH evaporation products changes essentially the observational appearance of a black hole at the final stage of its evaporation, leading to the transformation of nearly all of its total power into soft γ-ray emission.

The question which remains open is the evolution of a magnetic field in the plasma outflow. The characteristic time-scale for the magnetic field generation is probably determined by a turbulent diffusion, with a diffusion coefficient given by $D_m \sim c^2/4\pi\sigma \sim r_D^2\,\nu$. Here σ is an electric conductivity of plasma and ν an effective collision frequency, which is most likely determined by collective processes and not by Coulomb collisions. The local value of $\nu(r)$ cannot be less than c/r. Estimations of D_m for $\pi^-\pi^+$ and e^-e^+ stages of the wind show that the diffusion time needed to establish an equipartition magnetic field in the wind is much greater than the expansion time r/c, but much smaller than the lifetime of a black hole with temperature $T \sim 10$ GeV corresponding to the beginning of MHD expansion. Therefore, there should be a time delay between the beginning of the hydrodynamical expansion regime (when, say, $T \sim 10$ GeV), and a generation of turbulent magnetic field of enough strength to trigger the pair cascade (when $T \sim 0.1$–1 TeV). If the turbulent growth of the magnetic field begins in the outer regions of the outflow and propagates inward, the cascade is 'switched on' at a radius $r \simeq 0.2\lambda^{1/2}\tilde{T}^{3/2}$ cm, where the π_0-produced photons satisfy the condition $\xi \sim 1$. This is possible when T becomes greater than ~ 0.1 TeV. As the magnetic field diffuses inwards, the final energies of cascade-produced and synchrotron photons [equations (23) and (30)] decrease, and the observed spectrum exhibits a hard-to-soft evolution with a time-scale determined by the turbulent growth rate of the magnetic field. With further inward diffusion of the magnetic field, the electromagnetic cascade initiated by charged pions becomes possible. Note that the development of a $\pi^-\pi^+$-initiated cascade could suppress the cascade generated by neutral pions due to the reduced magnetic field at the radius r_{ph}.

4 RELATIVISTIC SHOCKS IN THE PBH EVAPORATION WIND

In this section we will not rely on the scenario of $\pi^-\pi^+$-initiated cascade discussed in the previous section, and assume that charged pions do not radiate their energy before decay.

One possible reason for large fluctuations of the radiation intensity and existence of a close-to-equipartition magnetic field in the outflow is the development of instabilities and shocks in the charged particle wind. The idea is similar to the blast-wave models of the cosmological GRB scenario [see Mészáros & Rees (1993); Mészáros, Laguna & Rees (1993); Rees & Mészáros (1994); Mészáros & Rees (1994); Sari & Piran (1995)]. In these models, an unsteady relativistic wind originating from the coalescence of two compact stars in a binary system leads to the internal and external shocks that dissipate the bulk kinetic energy of the wind in the MHD regime.

The internal shocks can be formed due to the inhomogeneity of the PBH evaporation wind which consists of various particle species that are created or decay in the course of expansion. One possible site of the shock formation is the radius $r = \bar{r}_{ph}$, where, according to the previous section, a persistent injection of a mildly relativistic e^-e^+ plasma occurs. An interaction of the $\pi^-\pi^+$ wind with this e^-e^+ cloud could produce shocks and instabilities in the outflow, leading to fluctuations in the radiation intensity. At the same time, the presence of the shock at radius $r = \bar{r}_{ph}$ could support the turbulent magnetic field, which is important for the scenario in the previous section.

Strong relativistic shocks can also be formed when the PBH evaporation products start decelerating in an interstellar medium. We may try to apply to this situation an approximate self-similar solution for the relativistic, spherically symmetric blast wave (Blandford & McKee 1976). The results are similar to those obtained by Blandford & McKee (1976) since they are derived from simple considerations based on the energy and momentum conservation laws.

Suppose that at the initial momentum of time $t = 0$, when the PBH temperature $T = T_0$, the black hole evaporation turns into the hydrodynamical expansion regime. The contact discontinuity between the jet products and the interstellar medium propagates outward with a bulk Lorentz factor of

$$\Gamma_0 \sim P_c(T_0)/\dot{M}_c(T_0)c^2, \tag{32}$$

where P_c is the power of the charged particle wind and $\dot{M}_c$ the rest-mass ejection rate of charged particles. At the pion stage of the wind we can neglect the contribution of other particle species and obtain

$$\Gamma_0 \sim \bar{\gamma}_\pi = \bar{E}_\pi/m_\pi c^2 \sim 10^2\tilde{T}_0^{1/2}. \tag{33}$$

As interstellar plasma is swept up by the expanding wind, an external shock may be formed which moves ahead of the discontinuity with a Lorentz factor of order Γ_0 (Blandford & McKee 1976). When the pressure of swept-up particles becomes comparable to the wind pressure, the jet products begin decelerating. This should occur at the PBH-frame radius r_s given by

$$\frac{P_c}{4\pi r_s^2 \Gamma_0^2 c} \sim \Gamma_0^2 \rho c^2, \tag{34}$$

where ρ is the density of the ionized fraction of the interstellar medium, assumed to be uniform and cold. The resulting value of r_s turns out to be independent of time (or PBH temperature):

$$r_s \sim \left(\frac{P_c}{4\pi\rho c^2 \Gamma_4^0}\right)^{1/2} \sim 4\times 10^6 (\rho_{-26})^{-1/2}\ \mathrm{cm}, \tag{35}$$

where the scaling $\rho_{-26} = \rho/10^{-26}$ g cm^{-3} is chosen according to the conditions in the Local Fluff (within a few light-years from the Sun) (Cox & Reynolds 1987). In equation (35) the last estimation is made assuming that the wind consists of charged pions or muons (we will not differentiate between the pion and muon stages of the wind, since this practically does not change the estimations for Γ_0 and r_s). However, the radius at which muons decay is approximately

$$r_\mu \sim 3\times 10^4 \bar{\gamma}_\pi \simeq 3\times 10^6 \tilde{T}_0^{1/2}\ \mathrm{cm}. \tag{36}$$

Therefore, for $T_0 \lesssim 2/\rho_{-26}$ TeV the muons decay before the wind starts decelerating. In this case the estimation in (35) for the deceleration radius is no more correct. Due to the decrease in power ($P_{e\pm} \sim 0.2 P_{\pi\pm} \sim 0.1 P_{\mathrm{tot}}$), the pressure of the resulting e^-e^+ wind at a radius r_μ is already *less* than the pressure of an interstellar medium. Therefore, starting from the radius $r \sim r_\mu$, the Lorentz factor of the shock, $\Gamma \gg 1$, decreases according to the following energy conservation law:

$$\frac{1}{\Gamma^2(t)} \int_0^{t_{\mathrm{ret}}(t)} P_{e\pm}(t')\,\mathrm{d}t' \sim (4/3)\pi R^3(t) \rho c^2 \Gamma^2(t), \tag{37}$$

where $R \simeq c(t - t_{\mathrm{ex}}) \simeq ct$ is the radius of the shock in the PBH frame and $t_{\mathrm{ex}} \simeq r_\mu/c$ is an expansion time to radius r_μ, which is very small for time-scales of interest. In equation (37) we take into account that the energy is continuously supplied to the shock in the form of a fluid with a high bulk Lorentz factor $\sim 10^2 \tilde{T}^{1/2}(t)$. Equation (37) expresses the fact that at a given time $t = R/c$, the total energy of particles swept by the shock, $\rho c^2 \Gamma^2 (4/3)\pi R^3$, is approximately equal to the amount of energy in e^-e^+ pairs that has been emitted by a hole and has *reached* the shock front before this moment (to reach the shock at the moment R/c the particles should have been emitted by the hole at the moment $t_{\mathrm{ret}} \simeq R/c\Gamma^2$). Using equations (4) and (14), we can rewrite (37) as

$$\beta M(T_0) c^2 \left\{1 - \left[1 - \frac{t}{\Gamma^2 \Delta\tau(T_0)}\right]^{1/3}\right\} \sim (4/3)\pi R^3 \rho c^2 \Gamma^4. \tag{38}$$

When $t \ll \Gamma^2 \Delta\tau(T_0)$, we obtain

$$\Gamma \sim \left(\frac{P_{e\pm}(T_0)}{4\pi R^2 \rho c^3}\right)^{0.1} \sim 0.3 \tilde{T}_0^{0.8} \left(\frac{\Delta\tau(T_0)}{t}\right)^{0.2}. \tag{39}$$

An external shock becomes non-relativistic at the radius $R_0 \sim 3\times 10^{10} \tilde{T}_0$ cm. It is clear that an external shock is only marginally relativistic, $\Gamma \lesssim 10$, during most of the PBH lifetime. It can be shown that synchrotron and inverse Compton radiation of particles accelerated in the external shock produces only a rather weak optical-to-microwave counterpart to the main gamma-ray burst; cf. Mészáros & Rees (1993), Mészáros, Laguna & Rees (1993) and Sari & Piran (1995), in which an external shock is the primary source of a gamma-ray burst.

At the wind-deceleration stage, the ultrarelativistic flow of charged particles that are continuously supplied by a central black hole should be somewhat decelerated to the velocity of an external shock. This can be realized due to the formation of an internal shock (Blandford & McKee 1976). As far as $\Gamma \gg 1$, an internal shock also propagates outward with a relativistic velocity in the PBH frame (Blandford & McKee 1976). After the moment $t_0 = R_0/c \sim 1\tilde{T}_0$ s, an internal shock may start propagating toward the central black hole with a subrelativistic velocity. Its motion will be, however, ultrarelativistic in the frame comoving with the PBH wind which expands with a bulk Lorentz factor $\Gamma_w \sim \bar{\gamma}_\pi \sim 10^2 \tilde{T}^{1/2}(t)$ at the pion stage, and $\Gamma_w \sim \bar{\gamma}_{e\pm} \sim 10^4 \tilde{T}^{1/2}(t)$ at the e^-e^+ stage of the wind. The reverse shock randomizes the bulk kinetic energy of the outflow, and can be also very important for supporting a close-to-equipartition magnetic field (13).

5 DISCUSSION

In this section we discuss the possible relation of PBH explosions to the detected GRBs and summarize some observational predictions that follow directly from the universal nature of the PBH evaporation process.

5.1 Energetics of PBH-originated bursts

Suppose that the detector is triggered at the moment of time $t = 0$ when the PBH temperature $T = T_0$. The total duration of the burst is determined by the lifetime of a hole with temperature T_0, $\Delta\tau(T_0)$ (see equation 5). Of course, there should be fluctuations on much shorter time-scales, defined by the turbulent diffusion time, the observer-frame synchrotron cooling time, and the expansion time. The mean luminosity of the burst cannot be higher than $\bar{L} \sim \bar{\beta} M(T_0) c^2/\Delta\tau(T_0) \sim \bar{\beta} 10^{28} \tilde{T}_0^2$ erg s^{-1}, where $\bar{\beta}$ is a fraction of the PBH rest mass converted into the radiation in the 0.1–1 MeV interval, expected to be of order 1/3. The source of luminosity $\bar{L}$ will be detected by BATSE from the maximum distance

$$d_{\mathrm{max}} \sim \left(\frac{\bar{L}}{4\pi F_{\mathrm{th}}}\right)^{1/2} \simeq 10^{17} \bar{\beta}^{1/2} \tilde{T}_0\ \mathrm{cm}. \tag{40}$$

Here $F_{\mathrm{th}} \simeq 10^{-7}$ erg cm^{-2} s^{-1} is the threshold peak flux which triggers the BATSE instrument (Fishman et al. 1994). The value of $\bar{L}$ rapidly grows with the PBH temperature. However, BATSE detectors are triggered by the peak flux accumulated during one of the three time intervals, $t_{\mathrm{det}} = 64$, 256 or 1024 ms (Fishman et al. 1994). Therefore, the optimal conditions for triggering BATSE by the faintest (most distant) bursts are realized when $\Delta\tau(T_0) \sim 64$ ms. This gives

$T_{max} \simeq 20$ TeV. The corresponding value of d_{max} is of the order of $0.7\bar{\beta}^{1/2}$ pc, which is at least two orders of magnitude greater than the value of d_{max} obtained under the assumption that the photon flux from the black hole consists mainly of GeV photons produced in π_0 decays (Semikoz 1994). The rate of the PBH explosions in the effective volume scanned by BATSE is

$$\dot{N} \simeq \frac{dn}{dt}\,\frac{4}{3}\,\pi\, d_{max}^2. \tag{41}$$

Accepting an estimation of the PBH explosion rate given by equation (5), we arrived at ~ 10 detectable bursts per year. Of course, this estimate is very sensitive to the precise values of the PBH density and the explosion rate (see equation 5).

To increase the predicted number of bursts detectable by BATSE, we should either assume that a local enhancement of the PBH density is greater than that for other matter in the Galactic Halo, or suppose that the peak intensity during the final stage of the evaporation ($T_0 \gtrsim 1$ TeV) can be much higher than the mean value. In the latter case, to obtain $d_{max} = 3$ pc, and therefore $\dot{N} \simeq 10^3$ yr^{-1}, the peak luminosity should be of the order of 10^{32} erg s^{-1}. This is nearly two orders of magnitude higher than the average luminosity of a black hole with $T = 20$ TeV, and is equivalent to the rest energy of a hole of temperature $\tilde{T}$ TeV released during $0.1/\tilde{T}$ s. If such outbursts are possible, the total energetics of PBH explosions could be consistent with current BATSE statistics for all GRBs. Indeed, the evaporation of a black hole of temperature $\tilde{T}$ TeV at a distance d from the Earth will produce a burst with a MeV fluence of

$$f = \frac{\beta M c^2}{4\pi d^2} \simeq 8 \times 10^{-8} \beta \tilde{T}^{-1} \left(\frac{d}{1\ \mathrm{pc}}\right)^{-2} \mathrm{erg\ cm^{-2}}. \tag{42}$$

The sensitivity threshold of BATSE detectors can be estimated as (Fishman et al. 1994) $f_{min} \sim 10^{-8}$ erg cm^{-2}. Therefore, the maximum distance scanned by BATSE is $d_{max} = (\beta M c^2/4\pi f_{min})^{1/2} \sim 3\beta^{1/2}\tilde{T}^{-1/2}$ pc, leading to $\sim 10^3$ detected bursts per year. This is in agreement with the all-sky GRB rate, estimated as ~ 800 yr^{-1} (Fishman et al. 1994).

5.2 Spatial distribution of bursts

It is natural to expect PBHs to have a homogeneous and isotropic distribution within a few parsecs from the Sun. Therefore, an attempt to relate *all* observed GRBs to evaporating black holes is faced with the problem of an apparent deficiency of weak bursts (deviation of the $\log N$–$\log S$ curve from the $-3/2$ law, see Fishman et al. 1994). Of course, until BATSE results are independently confirmed, there remains the possibility that the deviation from a $-3/2$ slope is merely an instrumental effect. It was argued that the same result may be obtained if the burst distribution is homogeneous, but the probability of the GRB registration falls with decreasing peak intensity, so that the weak bursts are lost (Bisnovatyi-Kogan 1995). If we assume that the deficiency of weak bursts is a real fact, there could be, basically, two reasons for it: first, the heliocentric distribution of PBHs (Horack et al. 1994). This would simultaneously increase the number of bursts due to a larger clustering factor. However, this is possible only if the PBHs had very small relative velocities with respect to the Sun, $\lesssim 1$ km s^{-1}, at the time of the Sun's birth. A second explanation could be the differences in peak luminosities that would originate from different time evolutions of turbulent plasma fireballs, or from differing external conditions. To agree with observations, there should be at least an order-of-magnitude dispersion in the intrinsic peak luminosities of PBH explosions. Also, a sort of fine tuning of PBH distribution over their peak luminosities may be required. Of course, if the PBH-originated bursts constitute only a small fraction of all observed bursts, there is no such problem (until we identify this fraction).

5.3 Spectral peculiarities

(1) Since evaporating black holes are more or less standard candles (at least, with respect to their fluences and mean luminosities $\bar{L}$), the brighter bursts are simply located closer to the Earth than are the fainter bursts. The closely located evaporating black holes (bright bursts) should trigger the BATSE detector at lower initial PBH temperatures than should more distant bursts. According to the expressions for break energies of photons in the cascade (equations 23 and 30), this means that bright bursts should be harder than faint ones. Also, the hard-to-soft evolution of the spectrum follows directly from the analysis in Section 3.2. Both facts are in qualitative agreement with observations (Nemiroff et al. 1994; Bhat et al. 1994; Mitrofanov et al. 1996). Note also that the temporal profile of the PBH-originated burst is expected to be asymmetric, and may consist of an irregular sequence of pulses with hard-to-soft evolution.

(2) As was discussed, e.g. in Baring (1990), pair cascades in a strong magnetic field may produce spectral breaks around 1 MeV, unless two-photon processes smear out this feature. Curiously, several GRBs exhibiting such spectral breaks have been identified (Schaefer et al. 1992). Of course, it is premature to associate just these bursts with evaporating black holes.

(3) The PBH evaporation must be accompanied by neutrino flux originating from pion decays, which may carry a significant fraction of the total emitted power (MacGibbon & Webber 1990; Halzen et al. 1991; Cline & Hong 1992). Also, there is a possibility of detecting the primary particles of energy $\sim 5T$. However, the probability of observing these events with the modern generation of detectors is very low (Halzen et al. 1991; LoSecco 1994).

5.4 Other observational aspects

In view of the discussion in Sections 3 and 4, it makes no sense to search for afterbursts in order to identify some of the observed GRBs as being the result of evaporating black holes, since afterbursts are expected to be very weak if they exist. Instead, one might try to look for the weak quasi-stationary source that existed in the GRB error box *before* the main burst. Such a source appears after the onset of the MHD expansion stage, i.e. when the PBH temperature exceeds ~ 10 GeV. It reveals itself between microwave and UV frequencies due to the synchrotron radiation of e^-e^+ pairs produced in $\pi^-\pi^+$ decays. Assuming an equipartition magnetic field (13), the synchrotron luminosity of pion-produced e^-e^+ pairs has a maximum at frequencies

$$\omega \sim 7 \times 10^{16} \tilde{T}^{3/2} \text{ rad s}^{-1}, \tag{43}$$

and is equal to

$$L \sim 10^{23} \tilde{T}^4 \text{ erg s}^{-1}. \tag{44}$$

The $T=100$ GeV black hole with a lifetime of 6 d before evaporation is an object of 23 mag at a distance of 10^{-2} pc. Therefore, the observational prospects for the quiescent pre-burst sources are not very promising. To increase the chances for identification, optical observations during the main gamma-ray burst would be desirable, since optical luminosity (see equation 44) rapidly grows with PBH temperature.

Independently on any ad hoc assumptions about plasma-like behaviour of evaporation wind and the existence of a close-to-equipartition magnetic field, a TeV-temperature black hole may produce a non-negligible photon flux in the MeV range due to the electromagnetic cascade initiated by collisions of *primary* photons and e^-e^+ pairs with *jet-produced* charged pions. (Remember that most of the pions are only weakly relativistic.) The optical depth for the relevant processes (energy losses of primary photons or electrons with Lorentz factor $\gamma \sim 10^6 \tilde{T}$ resulting from Bethe–Heitler pair production or bremsstrahlung emission), calculated in the usual way as an integral from some initial distance $r_{\rm in}$ to infinity, is greater than unity for $r_{\rm in} < 10^{-8} \tilde{T}^{3/2}$ cm. For TeV temperatures this is 20 times greater than the radius at which pions are created, $r_0 \sim 5 \times 10^{-10} \tilde{T}$ cm. Therefore, an electromagnetic cascade may develop. However, it is questionable whether the energy of primaries can degrade to the MeV range, because the created particle cloud may be very opaque for the two-photon processes. Furthermore, even if the cascade proceeds to MeV energies, the power contained in primary photons and *e–p* pairs is only 2–5 per cent of the total PBH power, and the resulting burst will be overshadowed by the electromagnetic cascade in the MHD wind (discussed in Section 3), *if* the latter really develops.

In conclusion, there exists the possibility to detect γ-radiation from a PBH at the final stage of its evaporation. We suggest that some of the GRBs detected by BATSE can be associated with evaporating black holes.

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