THE ASTROPHYSICAL JOURNAL, 365: L55–L58, 1990 December 20

THE APPEARANCE OF COSMIC FIREBALLS

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Received 1990 June 26; accepted 1990 October 2

ABSTRACT

We examine the influence of matter on a radiation fireball. Even when a small amount of baryonic material is present, most of the radiation energy of the fireball will be converted to a kinetic energy of the matter, and almost no electromagnetic signal will be observed. We discuss the implication of this result for several processes in which such a fireball might occur.

Subject headings: gamma rays: bursts — stars: neutron

I. INTRODUCTION

When a large amount of radiative energy is suddenly released into a compact region, an opaque "fireball" of photon-lepton gas is created due to the prolific creation of electron-positron pairs. This radiative sphere expands and cools rapidly until the energy of the photons degrades below the pair-production threshold and the sphere becomes transparent. If some amount of baryonic matter is mixed with the fireball, some of the radiative energy will be converted to a kinetic energy, and the signal from such a burst will be weakened. If the baryonic mass is large enough, almost all the radiation energy will be converted to kinetic energy of the baryonic matter which will be accelerated to relativistic velocities and there will be no observed γ-ray signal.

The pioneering qualitative study of cosmic fireballs was done by Cavallo and Rees (1978), who were motivated by the observed phenomenon of gamma-ray bursts (GRBs). These authors explored the parameter space of a matter-dominated fireball. Huter and Lingenfelter (1983) introduced a fireball model for the 1978 March 25 GRB and concluded that the model is applicable to a large class of GRB, since hard emission components are a common feature of the bursts spectra. However, they assumed that the average photon and lepton energy ($\gamma m_e c^2$) remains constant throughout the fireball evolution, and they ignored the acceleration and conversion of radiative energy to kinetic energy that takes place. This feature was pointed out later by Goodman (1986), who considered a fireball of pure radiation. He noticed that while the photons cool, their bulk motion increases and reaches an extremely large γ factor. The escaping photons are blueshifted, and the final overall spectrum is a modified blackbody spectrum at the initial temperature. If optically thick GRB models like the fireball are applicable, then one might conclude that GRBs are at cosmological distances, and the amount of energy released is some fraction of a solar mass (Paczyński 1986).

Several mechanisms of dense energy injection were recently proposed. We mention three of them here: Michel (1988) suggested that accretion onto a neutron star can cause a supernova-like energy injection when the central density grows beyond that of nuclear matter and a core of tightly bound quark complexes (i.e., new class of particles) is created. Eichler *et al.* (1989) considered the coalescence of neutron stars, and Dar and Ramaty (1990) discussed the creation of a naked neutron star from a white dwarf's core collapse in an accreting binary system. Both phenomena are accompanied by a strong neutrino-antineutrino burst that produces, by annihilation, a burst of electron-positron pairs and gamma rays.

In this *Letter* we explore the influence of baryonic matter on the evolution of an expanding homogeneous fireball, and we study some applications of this model to neutron star coalescence, naked neutron star formation and the GRB phenomenon.

II. THE MODEL

We consider a homogeneous fireball of pure energy E_0 and ionized hydrogen plasma of total mass M initially confined to a sphere of radius R_0. A subscript 0 denotes initial values. The crucial parameter for the fireball evolution is the initial ratio of radiation energy to baryonic rest mass energy: $\eta \equiv E_0/Mc^2$. Since the physical processes are dominated by leptons, we will use the dimensionless $\mathscr{E}$, $\mathscr{T}$, $\mathscr{R}$ for energy, temperature, and radius, measured in units of $m_e c^2$, $m_e c^2/k$, and $\lambda_e \equiv \hbar/m_e c$, and we define ϵ, the average energy of a photon, in these units.

If $\epsilon \geq 1$ and the photon density, n_γ, is sufficiently high, the opacity[2] due to Compton scattering and pair production processes (essentially via $\gamma\gamma \rightarrow e^+e^-$) will lead to a large optical depth. For $\epsilon > 1$ we have

$$\tau = \tau_{e\gamma} + \tau_{\gamma\gamma} \simeq \left(\frac{\sigma_T}{\epsilon^2}\right) n_\gamma R \simeq \frac{8\pi}{3} \frac{\alpha^2 \mathscr{E}_0}{\epsilon^3 \mathscr{R}^2} \quad (1)$$

($\alpha = 1/137$). Clearly, $\tau > 1$ is required in order that the initial fireball will be an opaque sphere in thermal equilibrium

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[2] Abramowicz, Novikov, and Paczyński (1990) point out that the opacity depends on the velocity of the matter. This is true only if we hold two points fixed in the lab frame, and we ask how does the optical depth between them vary as a function of the velocity of the matter. Here we are interested in the optical depth between the center of the fireball to its surface, and in this case we hold the total mass of the matter fixed. In this case, the optical depth does not depend at all on the velocity of the matter.

(Cavallo and Rees 1978). If $\mathscr{T} < 1$, the density of pairs is given by

$$n_{\pm} \approx \frac{\sqrt{2}}{\pi^{3/2}} \lambda_e^{-3} \mathscr{T}^{3/2} \exp\left(\frac{-1}{\mathscr{T}}\right)$$
$$= 4.41 \times 10^{30} \mathscr{T}^{3/2} \exp\left(\frac{-1}{\mathscr{T}}\right) \text{cm}^{-3} . \quad (2)$$

If the initial temperature exceeds $\mathscr{T} \approx 0.032$, the opacity due to pairs is important and generally dominant. Below it, the pair density is negligible. If the fireball is optically thick and thermal equilibrium is established we have

$$\mathscr{E}_0 = g \frac{4\pi^3}{45} \mathscr{T}_0^4 \mathscr{R}_0^3 , \quad (3)$$

where $g = g_0 = 11/4$ for $\mathscr{T} > 1$ (photons and pairs), and it drops to 1 when $\mathscr{T} \ll 1$ (only photons).

If $\tau \gg 1$, the fireball can be approximated as a homogeneous sphere in a thermal equilibrium, characterized by a single temperature at each step. This assumption is valid as long as the mean free path is much smaller than R and is much smaller than the length scale on which any physical quantity (e.g., temperature, pressure, velocity) changes. The first condition holds as long as $\tau > 1$. The second condition breaks down earlier, but it will mainly influence the details of the emerging spectrum which we do not attempt to calculate here. A photon moving forward relative to a given matter shell is redshifted, and a photon moving backward is blueshifted. This results in a distortion of the blackbody spectrum. If clumps are present, the opacity in different directions will vary, and the escape of the radiation (time scale, mean energy, anisotropy, etc.) will be more complicated. Also, if matter concentrates ahead of the radiation, Rayleigh-Taylor instability might occur. We ignore these complications, and we continue our calculations until $\tau = 1$, assuming homogeneity and thermal equilibrium.

In an opaque sphere, the photon escape probability is reduced by the optical depth factor $\tau \gg 1$. When the radiation energy dominates the evolution (i.e., $\eta \gg 1$), the expansion speed dR/dt is of the order of the speed of light, and up to a factor τ^{-1}, the radiative energy losses L are negligible relative to the $P\,dV$ work and one can assume an adiabatic flow:

$$\frac{L}{P(dV/dt)} = \frac{c}{(dR/dt)} \frac{1}{\tau} \simeq \frac{1}{\tau} . \quad (4)$$

If $E_0 > GM^2/R$, i.e., if $\eta > GM/c^2 R_0$, self-gravitational effects can be neglected[3] and the homogeneous fireball evolution is similar to the expansion of the early (Friedmann) universe at the stage when its curvature is negligible (see, e.g., Weinberg 1972). The specific entropy per gas particle is $\sigma \equiv (16\pi^3/135)(\mathscr{R}\mathscr{T})^3(m_p/m_e)(\eta/\mathscr{E}_0)$. If $\sigma \gg 1$, the radiation pressure dominates the fireball evolution. In this case, the temperature-radius relation at each stage is given by[4]

$$(g/g_0)^{1/3} \mathscr{R}\mathscr{T} = \mathscr{R}_0 \mathscr{T}_0 = \text{const} , \quad (5)$$

and $\sigma \approx (4/3)(m_p/m_e)(\eta/g\mathscr{T}_0) \approx 2500\eta/g\mathscr{T}_0$. Thus $\sigma > 1$ when $\eta > 10^{-3}\mathscr{T}_0$. (Note that for $\mathscr{T}_0 \simeq 1$ a fireball with even $\eta \simeq 10^{-2}$ is dominated by the radiation!) The total number of photons, $N_{\text{ph}} \propto (\mathscr{R}\mathscr{T})^3$, is conserved and σ is related to the ratio between photon density and the baryon density: $\sigma = 0.37 \times n_{\text{ph}}/n_{\text{bar}}$. The radiative energy $\mathscr{E}_R$ decreases proportionally to R^{-1}, i.e.,

$$\mathscr{E}_R = \mathscr{E}_0 \frac{\mathscr{T}}{\mathscr{T}_0} = \mathscr{E}_0 \frac{\mathscr{R}_0}{\mathscr{R}} . \quad (6)$$

When the temperature drops, Compton scattering dominates the opacity, and when $\mathscr{T} \leq 0.5$, the scattering cross section is well estimated by the Thompson cross section. The optical depth contains two contributions due to pair-produced electrons and positrons and due to the electrons from the ionized ambient gas

$$\tau = \sigma_{\text{T}} R(n_{\text{pair}} + n_{\text{gas}}) . \quad (7)$$

At $\tau \simeq 1$, the fireball becomes transparent and the radiation escapes in one crossing time. If at this stage the opacity is still dominated by the pairs ($n_{\text{pairs}} \gg n_{\text{gas}}$), the escape temperature, $\mathscr{T}_{\text{esc}}$, will be equal to $\mathscr{T}_p$ which we define (using eq. [2]) from the requirement

$$\tau \simeq \tau_p = \frac{\sqrt{128}}{3\sqrt{\pi}} \alpha^2 g_0^{1/3} \exp\left(-\frac{1}{\mathscr{T}_p}\right)(\mathscr{T}_0 \mathscr{R}_0)$$
$$= 1.6 \times 10^{-4} \mathscr{T}_p^{1/2} \exp\left(-\frac{1}{\mathscr{T}_p}\right)(\mathscr{T}_0 \mathscr{R}_0) = 1 . \quad (8)$$

Figure 1 shows that $\mathscr{T}_p \simeq 0.032$ for the range of parameters which is of interest to us. The condition for $\mathscr{T}_{\text{esc}} = \mathscr{T}_p$ is, therefore,

$$\eta > \frac{8\pi^3}{45} \frac{m_e}{m_p} \alpha^2 g_0^{1/3} \mathscr{T}_p^2 \mathscr{T}_0^2 \mathscr{R}_0 \approx 2.3 \times 10^{-10} \mathscr{T}_0^2 \mathscr{R}_0 . \quad (9)$$

As we will see later, $\mathscr{R} \approx 10^{16}$ and hence equation (9) yields a

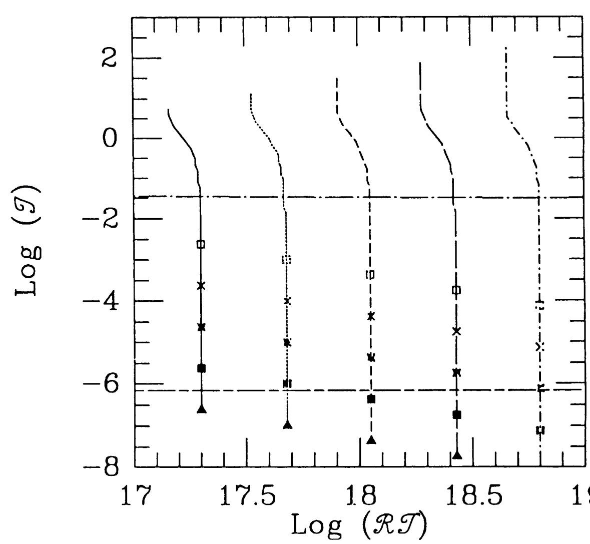

FIG. 1.—Trajectories of the fireball on the $\mathscr{R}\mathscr{T}$-$\mathscr{T}$ plane for $E_0 = 10^{47}$, $10^{48.5}$, 10^{50}, $10^{51.5}$, and 10^{53} ergs, corresponding to the solid, dotted, short-dashed, long-dashed, and dotted-dashed lines, respectively. The initial configuration is always at $R_0 = 10^6$ cm. The various points mark $\mathscr{T}_{\text{esc}} = \mathscr{T}_g$ for $\eta = 10^6$, 10^4, 10^2, 1, and 10^{-2} (with an open square, a cross, a star, a square, and a triangle, respectively). The long-dashed–dotted line marks $\mathscr{T}_p$. Note that $\mathscr{T}_p > \mathscr{T}_g$ for all the cases considered even with $\eta = 10^6$. The long-dashed–short-dashed line denotes $\mathscr{T} = 4000$ K, the recombination temperature.

[3] We also implicitly ignore external gravitational effects.

[4] If $\sigma \ll 1$, namely the gas pressure dominates, the expansion is of an ideal gas: $T \propto R^{-3(\Gamma-1)}$, where Γ is the specific heat ratio of the gas. This is never applicable in cosmology, but it might be relevant in some events of the type that we are discussing.

rather large lower limit on η (and a strong upper limit on the amount of baryonic matter).

The escape temperature may drop far below T_p if condition (9) is not satisfied (but it is always ≥ 4000 K, the recombination temperature[5]). In this case the optical depth $\tau \simeq \tau_g = \sigma_T n_{\rm gas} R \propto R^{-2}$, decreases proportionally to T^2. The escape temperature will reach the value of $\mathscr{T}_g$, defined from $\tau_g = 1$:

$$\mathscr{T}_g \simeq \sqrt{\frac{45}{8\pi^3}} \sqrt{\frac{m_p}{m_e}} \frac{1}{\alpha g_0^{1/6}} \mathscr{T}_0^{-1} \mathscr{R}_0^{-0.5} \eta^{0.5}$$
$$= 2.1 \times 10^3 \mathscr{T}_0^{-1} \mathscr{R}_0^{-0.5} \eta^{0.5} . \tag{10}$$

Figure 1 depicts the fireball evolution on the log $\mathscr{T}$–log $\mathscr{R}\mathscr{T}$ plane. It sketches a vertical downward line, ending at $\mathscr{T} = \mathscr{T}_{\rm esc} = \min(\mathscr{T}_p;\ \mathscr{T}_g)$. A deviation from this straight line occurs at $\mathscr{T} \approx 0.7$, where the photon number increases by a factor 11/4 due to unreversed converson of pairs to photons.

Along the fireball evolution, radiative energy is converted to kinetic energy as the radiation pressure accelerates the gas (photons, pairs, and matter). We approximate the motion of the gas by an average Lorentz factor $\gamma(T)$ that characterizes the bulk motion at each moment. Conservation of energy yields

$$\gamma(\mathscr{T}) = \frac{\mathscr{E}_0 + M}{\mathscr{E}_R + M} = \frac{\eta + 1}{\eta(\mathscr{T}/\mathscr{T}_0) + 1} . \tag{11}$$

The final value of γ, $\gamma_{\rm esc}$, depends on the ratio $\eta(\mathscr{T}_{\rm esc}/\mathscr{T}_0)$. γ approaches $\mathscr{T}_0/\mathscr{T}_{\rm esc}$ if $\eta \gg \mathscr{T}_0/\mathscr{T}_{\rm esc}$ and γ attains a constant value, $(\eta + 1)$, for which the acceleration stops if $\eta \ll \mathscr{T}_0/\mathscr{T}_{\rm esc}$. When we substitute equations (8), (9), and (10) into equation (11) [to obtain $\gamma(\mathscr{T}_{\rm esc})$ as a function of η], we find that γ has a maximum at an intermediate value of η (see Fig. 2).

The emitted energy is blueshifted by the factor γ (Goodman 1986), and the observed temperature is larger than the local temperature by this factor:

$$\mathscr{T}_{\rm obs} = \gamma \mathscr{T}_{\rm esc} = \mathscr{T}_{\rm esc} \mathscr{T}_0 \frac{\eta + 1}{\eta \mathscr{T}_{\rm esc} + \mathscr{T}_0} . \tag{12}$$

The observed spectrum is a blackbody spectrum, modified by relativistic geometrical considerations, by deviation from LTE, and by the radiative transfer processes that takes place just before $\tau = 1$. The ratio between the total radiative energy that is observed from the fireball and the initial radiative energy is equal to the ratio of the observed-to-initial temperature:

$$\frac{E_{\rm obs}}{E_0} = \gamma \frac{\mathscr{T}_{\rm esc}}{\mathscr{T}_0} = \frac{\mathscr{T}_{\rm obs}}{\mathscr{T}_0} . \tag{13}$$

Figure 2 shows the observed temperature and the relative radiation energy that is released when the fireball becomes transparent as a function of η.

[5] Note that at 4000 K $t_{\rm rec} c/R \approx (\sigma_T c)/(\langle \sigma v \rangle_{\rm rec} \tau_g) \approx 2 \times 10^{-2}/\tau_g < 1$, and there is enough time for the gas to recombine if it is optically thick to Thompson scattering.

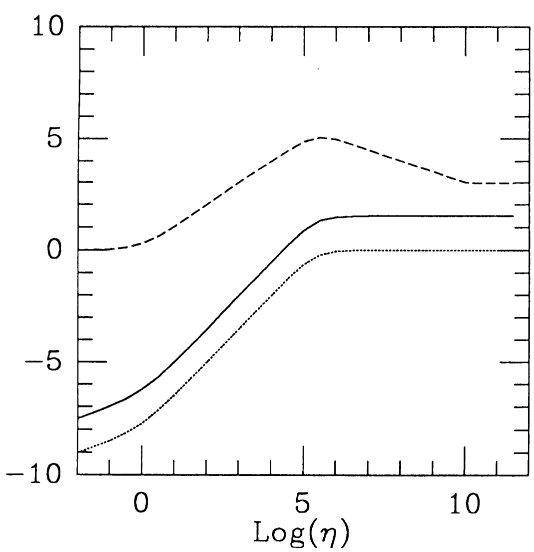

FIG. 2.—The relativistic factor γ (*dashed line*), the observed temperature $\mathscr{T}_{\rm obs}$ (*solid line*), and the fraction of the observed energy of the initial energy (*dotted line*) as a function of $\eta \equiv E_0/mc^2$ for a "canonical" fireball with $E_0 = 10^{50}$ ergs and $R_0 = 10^6$ cm. Note that γ has a maximum $\approx 10^5$ for an intermediate value of η.

III. DISCUSSION

The three mechanisms that we have mentioned in the introduction (Michel 1988; Eichler *et al.* 1989; Dar and Ramaty 1990) deal with situations in which $\approx 10^{50}$ ergs are suddenly released into a region with a typical size 10 km. This corresponds in our units to $\mathscr{E} \approx 1.2 \times 10^{56}$, $\mathscr{R} \approx 2.6 \times 10^{16}$, and $\mathscr{T} \approx 31$. We use these as canonical numbers in the following discussion.

If $\eta > 5.7 \times 10^9$, corresponding to $M < 10^{-14}\ M_\odot$, the escape temperature will be determined by the pairs, $\mathscr{T}_{\rm esc} = \mathscr{T}_p = 0.032$. In this case $\eta \gg \mathscr{T}_0/\mathscr{T}_{\rm esc}$, and the final observed temperature (eq. [12]) will be of order $T_0 \approx 15$ MeV, which is higher than observed in GRBs. This value of T_0 is reduced if one assumes that the energy is released over time $\Delta t \gg R_0/C$.

It is more likely that more mass is mixed with the radiation so the matter will dominate the opacity when the fireball becomes optically thin and we have $\mathscr{T}_{\rm esc} = \mathscr{T}_g$. Substitution of the canonical numbers into equation (10) gives $\mathscr{T}_{\rm esc} \simeq 4.2 \times 10^{-7} \eta^{0.5}$. In this case, more energy will be converted to the kinetic energy of the gas, which will be accelerated until it becomes optically thin, but yet, as is shown in Figure 2, as long as $\eta \geq 10^5$, namely $M \leq 10^{-9}\ M_\odot$, the photons will carry most of the kinetic energy and $E_{\rm obs} \approx E_0$.

If we require that any of these processes will lead, indeed, to some of the observed GRBs, we need $T_{\rm obs} \approx 100$–200 keV ($\mathscr{T}_{\rm obs} \approx 0.2$–0.4). This yields the condition $\eta \approx 8 \times 10^3$, i.e., $M \approx 6 \times 10^{-9}\ M_\odot$. If this condition is satisfied, we have $E_{\rm obs}/E_0 \approx 0.01$, about 1% of the earlier estimate of Eichler *et al.* (1989). Note that there are many bursts having power-law tails extending well above 1 MeV that contain most of the burst energy, and in these cases the observed-to-initial energy ratio increases.

Recently, Paczyński (1990) proposed a different way to approximate the same energy injection events. He considers a steady state, spherical symmetric wind which is driven away from a region surrounding a neutron star by highly super-Eddington rate of energy deposition. This model shows that it is possible to reach a burst of gamma-ray temperature ($T > 10^9$ K) with the optically thick winds, provided that the energy injection rate, $L_{\rm in}$, is sufficiently high, compared to the rest mass injection rate $\dot{M}c^2$: $L_{\rm in} > 10^2 \dot{M}c^2$. When $L_{\rm in}$ is that high, the observed photon luminosity, $L_{\rm out}$, becomes appreciable, with $L_{\rm out} > 10^{-2} L_{\rm in}$. For an event of $\approx 10^{50}$ ergs that occurs at 10 seconds, this translates to $L_{\rm in} = 10^{49}$ ergs s^{-1}, $L_{\rm out} = 10^{47}$ ergs s^{-1}, $T = 10^9$ K, and $\dot{M} \lesssim 10^{26}$ g s^{-1}. It can be immediately translated to our terms: $E_{obs}/E_0 \approx 1\%$ and $M < 5 \times 10^{-7} M_\odot$. In a view of the simplifications that were made in both of these two models, the correspondence of the results is remarkable.

If much more mass is present within the fireballs, almost all the radiative energy will be converted to kinetic energy. The escape temperature has a lower limit, since at $T \approx 4000$ K ($\mathscr{T} \approx 6.7 \times 10^{-7}$) the plasma recombines and the matter becomes transparent. For τ to be > 1 just before this stage, a lower limit for the rest mass is (see eq. [10]) $\eta \leq 1.6$ or $M \geq 3.5 \times 10^{-5} M_\odot$. This corresponds to low values of γ: $\gamma \approx (\eta + 1)$ and the observed temperature, $T_{obs} = T_{esc} \gamma = 4000 \times \gamma$ K. With our canonical parameters, this is a soft UV-optical event of $\leq 10^{43}$ ergs.

It seems that in many cases, perhaps in the majority of them, most of the initial energy will go to acceleration of the gas and the outcomes of the fireball will be the ejection of high-energy particles with typical relativistic gamma factors of order $\leq 10^5$, corresponding to $< 10^8$ MeV protons (γ cannot be much larger than $\approx 10^5$ for reasonable parameters) and a dilute burst, probably not in the γ range. If the processes that we discussed take place at a rate of N per year per galaxy ($N = 10^{-4}$ yr^{-1} per galaxy for the neutron stars coalescence), this will correspond to injection of $< 10^{-9} M_\odot$ protons yr^{-1} at this energy range to the galaxy.

We conclude that the pulses of gravitational radiation and neutrinos coming from a neutron star's coalescence may not be accompanied by a gamma-ray burst or may be accompanied with only a dilute burst, sometimes at observed temperature of the soft UV–optical range. This may pose a difficulty on the identification of such sources, but, at the same time, it might push the upper limit of the rate that such events occur. Similarly, we expect that the limit posed by Dar and Ramaty (1990) on the rate of naked neutron star formation (from the rate of GRB) might not be valid.

There are two possible ways to circumvent our conclusion. Both have to do with anisotropic optical depth that might be related to formation of clumps or to a Rayleigh-Taylor instability (if the matter is confined to a shell supported by radiation). In both cases, "windows" can appear through which radiation can escape before the radiation energy is completely depleted, and we might obtain GRBs even in the presence of baryonic matter. These possibilities and other details of this model are a subject of further study.

We thank B. Kozlovsky and B. Paczyński for helpful discussions. This work was supported by a BSF grant to the Hebrew University.

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