Electron-positron outflow from black holes

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Gamma-ray bursts (GRBs) appear as the brightest transient phenomena in the Universe. The nature of the central engine in GRBs is a missing link in the theory of fireballs to their stellar mass progenitors. Here it is shown that rotating black holes produce electron-positron outflow when brought into contact with a strong magnetic field. The outflow is produced by a coupling of the spin of the black hole to the orbit of the particles. For a nearly extreme Kerr black hole, particle outflow from an initial state of electrostatic equilibrium has a normalized isotropic emission of

\[ \sim 5 \times 10^{48} \left( \frac{B}{B_c} \right)^2 \left( \frac{M}{7M_\odot} \right)^2 \sin^2 \theta \text{ erg/s}, \]

where \( B \) is the external magnetic field strength, \( B_c = 4.4 \times 10^{13} \text{G} \), and \( M \) is the mass of the black hole. This initial outflow has a half-opening angle \( \theta \geq \sqrt{B_c/3B}. \) A connection with fireballs in γ-ray bursts is given.

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Gamma-ray bursts (GRBs) are now believed to be of cosmological origin, as inferred from an isotropic distribution in the sky [24], redshifts of order unity when detected [8,27,22] and \(< V/V_{\text{max}} > = 0.334 \pm 0.008\) distinctly less than the Euclidean value 1/2 (for long bursts) [38]. The observed emission is well described by the internal-external shock fireball model [36,37,33,34,26,11]. Their engines are compact, as indicated by the ratio \(\alpha = G\Delta E/c^5\delta t\), where \(\Delta E\) is the energy released, \(G\) is Newton’s constant, \(c\) is the velocity of light and \(\delta t \sim 1\text{ms}\) is the observed time-scale of variability. The typical values \(\alpha = 10^{-4} - 10^{-2}\) for GRBs is extremely large compared to other burst-type phenomena, such as supernovae including the GRB/SN 1998bw event [14] and outbursts in accreting systems such as GRS 1915+105 [25,23], and is close to the maximal value of unity.

The cosmological distances of GRBs indicate a typical energy release of the order of \(1M_\odot\), while the value of \(\alpha\) suggests an association with a black hole of a few solar masses. Black holes are a natural outcome of the evolution of binaries of young massive stars. Their evolution branches out over a diverse spectrum of intermediate states, such as accreting neutron stars or black holes with a red giant companion. These branches collectively produce either supernovae, their “failed” counterparts [50] or hypernovae [31,3], or coalescence of binaries made of neutron stars, black holes, white dwarfs or compact He-cores (see, e.g. [12]). The outcome of these events is probably associated with the formation of a new black hole, and most likely through an intermediate black hole/torus or black hole/disk state [30,50]. The energetic output from these configurations derives from tapping a fraction of the gravitational binding energy of the surrounding matter upon accretion onto the black hole, or tapping of the rotational energy of the black hole, which can reach about one third of its total rest mass. The associated efficiencies and beaming ultimately determine the apparent energies. The GRB990123 event [13] is important in this respect, in that its apparent energy was in excess of \(1M_\odot\) assuming isotropic emission.

This Letter focuses on the nature of the central engine. A theory is described for electron-positron pair-creation powered by a rapidly spinning black hole when brought into contact with a strong magnetic field. The magnetic field is supplied by the surrounding matter as in forementioned black hole/torus or disk systems. A rapidly spinning black hole couples to the surrounding matter by Maxwell stresses [43,46]. This coupling suppresses accretion, and in the case of a torus formed from the break-up of a neutron star, is expected to be intermittent on the time-scales of 0.15-1.5s [13]. In this case, the magnetic field is the remnant field of its progenitor, and if the torus remains at nuclear density, it may reach magnetic field strengths of up to \(10^{17}\text{G}\) by linear amplification [21]. The theory is described by perturbative calculations about a Wald field in electrostatic equilibrium. Particle creation towards electrostatic equilibrium has been considered in a previous analysis [13]. Here we
pay particular attention to the coupling of the black hole spin to the orbit of the charged particles, and find it is important in a continuation of the pair creation process.

Pair-creation can be calculated from the evolution of wave-fronts in curved spacetime, which is well-defined between asymptotically flat in- or out-vacua. By this device, any inequivalence between them becomes apparent, and generally gives rise to particle production \([7,1]\). It is perhaps best known from the Schwinger process \([28,10,4]\), and in dynamical spacetimes in cosmological scenarios \([4]\). Such particle production process is driven primarily by the jump in the zero-energy levels of the asymptotic vacua, and to a lesser degree depends on the nature of the transition between them. The energy spectrum of the particles is ordinarily nonthermal, with the notable exception of the thermal spectrum in Hawking radiation from a horizon surface formed in gravitational collapse to a black hole \([19]\). There are natural choices of the asymptotic vacua in asymptotically flat Minkowski spacetimes, where a time-like Killing vector can be used to select a preferred set of observers. This leaves the in- and out-vacua determined up to Lorentz transformations on the observers and gauge transformations on the wave-function of interest. These ambiguities can be circumvented by making reference to Hilbert spaces on null trajectories - the past and future null infinities \(J^\pm\) in Hawking’s proposal - and by working with gauge-covariant frequencies. The latter received some mention in Hawking’s original treatise \([19]\), and is briefly as follows.

Hawking radiation derives from tracing wave-fronts from \(J^+\) to \(J^-\), past any potential barrier and through the collapsing matter, with subsequent Bogoliubov projections on the Hilbert space of radiative states on \(J^-\). This procedure assumes gauge covariance, by tracing wave-fronts associated with gauge-covariant frequencies in the presence of a background vector potential \(A_a\). The generalization to a rotating black hole obtains by taking these frequencies relative to real, zero-angular momentum observers (cf. \([13,43]\)), whose worldlines are orthogonal to the azimuthal Killing vector as given by \(\xi^a \partial_a = \partial_t - (g_{t\phi}/g_{\phi\phi})\partial_\phi\). Then \(\xi^a \sim \partial_t\) at infinity and \(\xi^a \partial_a\) assumes corotation upon approaching the horizon, where \(g_{ab}\) denotes the Kerr metric. This obtains consistent particle-antiparticle conjugation by complex conjugation among all observers, except for the interpretation of a particle or an antiparticle. Consequently, Hawking emission from the horizon of a rotating black hole gives rise to a flux to infinity

\[
\frac{d^2n}{d\omega dt} = \frac{1}{2\pi} \frac{\Gamma}{e^{2\pi(\omega-V_F)/\kappa} + 1},
\]

for a particle of energy \(\omega\) at infinity. Here, \(\kappa = 1/4M\) and \(\Omega_H\) are the surface gravity and angular velocity of the black hole of mass \(M\), \(\Gamma\) is the relevant absorption factor. The Fermi-level \(V_F\) derives from the (normalized) gauge-covariant frequency as observed by a zero-angular momentum observer close to the horizon, namely, \(\omega - V_F = \omega_{ZAMO} + eV = \omega - \nu\Omega_H + eV\) for a particle of charge \(-e\) and azimuthal quantum number \(\nu\), where \(V\) is the
potential of the horizon relative to infinity. The results for antiparticles (as seen at infinity) follow with a change of sign in the charge, which may be seen to be equivalent to the usual transformation rule \( \omega \to -\omega \) and \( \nu \to -\nu \).

In case of \( V = 0 \) Hawking radiation is symmetric under particle-antiparticle conjugation, whereby Schwarzschild or Kerr black holes in vacuum show equal emission in particles and antiparticles. For a Schwarzschild black hole, then, the resulting luminosity of (1) is thermal with Hawking temperature \( T \sim 10^{-7}(M_{\odot}/M)K \), which is negligible for black holes of astrophysical size \([32,41]\). The charged case forms an interesting exception, however, where the Fermi-level \(-eV\) gives rise to spontaneous emission by which the black hole equilibrates on a dynamical time-scale \([15,42,6]\). In contrast, the Fermi-level \( \nu \Omega_H \) of a rotating black hole acting on neutrinos is extremely inefficient in producing spontaneous emission at infinity \([44]\).

This is due to an exponential cut-off due to a surrounding angular momentum barrier, which acts universally on neutrinos independent of the sign of their orbital angular momentum. This illustrates that (1) should be viewed with two different processes in mind: nonthermal spontaneous emission in response to a non-zero Fermi-level, and thermal radiation beyond \([7]\).

Upon exposing a rotating black hole to an external magnetic field this radiation picture is expected to change, particularly in regards to \( V_F \) and the absorption coefficient \( \Gamma \). The radiative states are now characterized by conservation of magnetic flux rather than conservation of particle angular momentum, which has some interesting consequences.

Consider a black hole of mass \( M \) and specific angular momentum \( a \), which is brought into contact with an external magnetic field \( B \) aligned with its axis of rotation. Its lowest electrostatic energy state is reached upon accretion of a Wald charge \([47]\) \( q = 2BJ \), where \( J = aM \) is the angular momentum of the black hole \([48]\). Hereby the source-free Wald solution of the vector potential,

\[
A_a = \frac{1}{2} Bk_a \left( \frac{q}{2M} - aB \right) \eta_a, \tag{2}
\]

reduces to

\[
A_a = \frac{1}{2} Bk_a, \tag{3}
\]

where \( k^a \) is the azimuthal Killing vector and \( \eta^a \) is the asymptotically time-like Killing vector. In electrostatic equilibrium, then, \( \xi^a A_a = 0 \). The charge \( q \) can be understood by noting that it restores the flux through a hemisphere of the horizon of a rapidly spinning black hole \([8]\). On dimensional grounds, it contributes a flux \( \Phi' = kqM\Omega_H = 2kB M^2 \sin^2(\lambda/2) \), where \( k \) is a dimensionless constant of proportionality and \( \Omega_H = \tan(\lambda/2)/2M \) using \( \sin \lambda = a/M \) \([45]\). With \( k = 4\pi \), \( \Phi' \) plus the flux \( 4\pi BM^2 \cos \lambda \) through a charge-free horizon recovers the
Schwarzschild value $\Phi = 4\pi BM^2$, as given by the exact expression (2). It should be noted that the equilibrium value $q = 2BJ$ obtains the uncharged Kerr metric upon neglecting gravitational contributions of the stress-energy tensor of the electromagnetic field [9]. The state of electrostatic equilibrium of a black hole in vacuum is therefore described by (3) corresponding to a maximal horizon flux $2\pi A_\phi$ [47,9].

In an axisymmetric magnetic field $B$ parallel to the axis of rotation, the wave-functions of charged particles can be expanded locally in coordinates $(\rho, \phi, s, t)$ as $e^{-i\omega t}e^{i\nu \phi}e^{ip_s s}\psi(\rho)$, where $s$ denotes arclength along the magnetic field. Comparison with the theory of plane-wave solutions [20] gives a localization on the $\nu$-th flux surface at which $g_{\phi\phi}^{1/2} = \sqrt{2\nu/eB}$ with Landau levels $E_{n\alpha} = \{m_e^2 + p_s^2 + |eB|(2n + 1 - \alpha)\}^{1/2}$, where $m_e$ is the electron mass and $\alpha = \pm 1$ refers to spin orientation along $B$. These states enclose a flux $A_\phi = \frac{1}{2} Bk^2 = \nu/e$. Note that the latter have effective cross sections $\Sigma_\nu = 2\pi/|eB|$. The gauge-covariant frequency of the Landau states near the horizon follows from

$$-\xi^a (i^{-1} \partial_a + eA_a)\psi = (\omega - \nu \Omega_H)\psi.$$  

The jump

$$V_F = [-\xi^a (i^{-1} \partial_a + eA_a)]^H_{\infty} \psi = \nu \Omega_H$$

between the horizon and infinity defines the Fermi-level of the particles at the horizon. In contrast, the Wald field about an uncharged black hole has $V_F = \nu \Omega_H - eaB_0$, which shows that it is out of electrostatic equilibrium. Note that the canonical angular momentum of the Landau states vanishes: $k^a \pi_a \psi = (i^{-1} \partial_\phi - eA_\phi)\psi = 0$. The Fermi-level (4) combines the spin coupling of the black hole to the vector potential $A_\alpha$ and the particle wave-function $\psi$. The equilibrium state in the sense of $\partial_t q \sim 0$, or at most $q/\partial_t q \sim a/\partial_t a$, derives from this complete $V_F$. For this reason, we shall study the state of electrostatic equilibrium as an initial condition, to infer aspects of the late time evolution.

The strength of the spin-orbit coupling which drives a Schwinger-type process on the surfaces of constant flux may be compared with the spin coupling to the vector potential $A_a$. The latter can be expressed in terms of the $EMF_\nu$ over a loop which closes at infinity and extends over the axis of rotation, the horizon and the $\nu$–th flux surface with flux $\Psi_\nu$. Thus, we have $EMF_\nu = \Omega_H \Psi_\nu/2\pi$ [2,43], which gives rise to the new identity

$$eEMF_\nu = \nu \Omega_H.$$  

It should be mentioned that (6) continues to hold away from electrostatic equilibrium (i.e.: $q \neq 2BJ$), since $\xi^a A_a = 0$ and hence $\nu - eA_\phi = 0$ on the horizon. Since the latter is a conserved quantity, it, in fact, continues to hold everywhere in the Wald-field approximation.
In the assumed electrostatic equilibrium state, $\xi^a A_a = 0$, and the generalization of (3) to points $(s, \nu)$ away from the horizon is

$$\left[-\xi^a (i^{-1} \partial_a + e A_a)\right]_{\infty}^{(s, \nu)} \psi = -\nu \frac{g_{a\rho}}{g_{\phi\phi}} (s, \nu) = -\frac{1}{2} e B g_{t\phi} (s, \nu) = -e A_t (s, \nu)$$

(7)

for particles of charge $-e$. Thus, (7) localizes (3) by expressing the coupling of the black hole spin to the wave-functions in terms of the electrostatic potential $V = A_t$ in Boyer-Linquist coordinates. Note that the real zero-angular momentum observers detect no electric potential.

The particle outflow derives from the distribution function (1) by calculation of the transmission coefficient through a barrier in the so-called level-crossing picture (4). The WKB approximation (e.g.: as derived by zero-angular momentum observers) gives the inhomogeneous dispersion relation

$$(\omega - V_F)^2 = m^2_e + |eB|(2n + 1 - \alpha) + p^2_s,$$

(8)

where $V_F = V_F(s, \nu)$ is the $s$-dependent Fermi-level on the $\nu$-th flux surface. The classical limit of (8) is illustrative, noting that the energy $\epsilon$ of the particle is always the same relative to the local ZAMOs that it passes. Indeed, since $w^a (m a_a - e A_a)$ is conserved when $w^a$ is a Killing vector (19), $\eta^a (m u_a - e A_a) = \pi_t$ and $k^a (m u_a - e A_a) = \pi_\phi$ are constants of motion, where $u^a$ is the four-velocity of the guiding center of the particle, and $\pi_t = E_{na}, \pi_\phi = 0$ in a Landau state. With $\xi^a A_a = 0, \epsilon = -\xi^a m u_a = -\xi^a (m u_a - e A_a) = -\eta^a (m u_a - e A_a) = \pi_t$. This conservation law circumvents discussions on the role of $E \cdot B$ (generally nonzero in a Wald field). The energy of the particle relative to infinity is $\omega$. This relates to the energy $\epsilon$ as measured by the ZAMOs following a shift $V_F(s, \nu)$ due to their angular velocity. Thus (8) pertains to observations in ZAMO frames, but is expressed in terms of the energy at infinity $\omega$. It follows that particle/antiparticle pair creation (as in pair creation of neutrinos (44)) is set by

$$\eta = |\partial V_F / \partial s| \sim \left| \partial_r \left( \frac{1}{2} e B g_{t\phi} \right) \right| = |\partial_r (e A_t)| = e B a M \frac{r^2 - a^2 \cos^2 \theta}{(r^2 + a^2 \cos^2 \theta)^2} \sin^2 \theta,$$

(9)

using $\partial_s \sim \partial_r$. Radiation states at infinity are separated from those near the horizon by a barrier where $p^2_s < 0$ about $V_F(s_0) = \omega$. The WKB approximation gives the transmission coefficient

$$|T_{na}|^2 = e^{-\pi m^2_e + |eB|(2n + 1 - \alpha)} / \eta.$$  

(10)

Since the Wald field $B$ is approximately uniform, any additional magnetic mirror effects can be neglected. Also, $\eta \leq \frac{1}{8} e B (M/a) \tan^2 \theta \leq \frac{1}{4} e B$ and $|eB|(2n + 1 - \alpha) / \eta \geq 4(2n + 1 - \alpha)$, so that $T$ is dominated by $n = 0$ and $\alpha = 1$. 

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By (7), the pair production rate by the forcing $\eta$ in (4) can be derived from the analogous results for the pair production rate produced by an electric field $E$ along $B$. The results from the latter [6, 4] imply a production rate $\dot{N}$ of particles given by

$$
\dot{N} = \frac{e}{4\pi^2} \int \frac{\eta Be^{-\pi m^2/\eta}}{\tan(\pi Be/\eta)} \sqrt{-g} d^3 x \sim \frac{e^2 B^2 M a}{2\pi} \int \frac{r^2 - a^2 \cos^2 \theta}{r^2 + a^2 \cos^2 \theta} e^{-\pi m^2/\eta} \sin^3 \theta dr d\theta.
$$

Here $1/\eta \sim (e BaM \sin^2 \theta)^{-1}(8a^2 + 12(r - \sqrt{3}a \cos \theta)^2)$ about $r = \sqrt{3}a \cos \theta$. For a rapidly spinning black hole, $\sqrt{3}a \cos \theta$ is outside the horizon in the small angle approximation, whereby after $r$-integration of (11) we are left with

$$
\dot{N} \sim \frac{e^2 B^2 a^2 M}{8\pi \sqrt{3} c} \int e^{-8\pi c/\sin^2 \theta} \sin^4 \theta d\theta \sim \frac{N_H^2}{128\sqrt{3} \pi^2 M} \left(\frac{a}{M}\right)^4 c^{-7/2} e^{-8\pi c/\theta^2} \theta^7
$$

asymptotically as $8\pi c/\theta^2 \gg 1$. Here $c = m_e^2 a/e BM$, $N_H = m_e^2 M^2$ is a characteristic number of particles on the horizon, and $\theta$ is the half-opening angle of the outflow. The right hand-side of (12) forms a lower limit in case of $8\pi c/\theta \leq 1$. When $a \sim M$, $N_H/c$ is characteristic for the total number of flux surfaces $\nu$, which penetrate the horizon and $c \sim Bc/B$, where $Bc = 4.4 \times 10^{13}$G is the field strength which sets the first Landau level at the rest mass energy. By (3) and (11), a similar calculation obtains for the luminosity in particles $L_p$ normalized to isotropic emission the estimate

$$
L'_p \sim \frac{L_p}{\theta^2 / 4} \sim \sqrt{3} eBM \dot{N}.
$$

This analysis shows that the black hole departs from electrostatic equilibrium by outflow of $e^-$ (with the sign convention $B \Omega_H > 0$) towards $q > 2BJ$. The calculations therefore pertain to the initial jet formation. Upon accumulating $q > 2BJ$, the jet evolves towards an inner jet of $e^+$ outflow produced by $V_F$ which is dominated by a positive electrostatic potential near the polar caps, surrounded by $e^-$ outflow produced by $V_F$ which is dominated by the spin-orbit coupling as in the initial jet. A full calculation of the resulting equilibrium outflow (in the sense as described before) falls outside the present scope. Nevertheless, it is expected that the luminosity (13) remains characteristic for the combined particle-antiparticle outflow in the evolved jet.

The outflow (13) as follows from (12) saturates when the back reaction of the outflow is taken into account. This is non-dissipative in the form of an attenuation of the magnetic field by the collective contribution of the charged particles to azimuthal currents, and dissipative due to a finite surface conductivity of $4\pi$ (see [13]) in current closure over the horizon. The four-current of the charged particles, and the associated a Poynting flux, forms a perturbation away from the source-free Wald field in our derivation. While these two types of back reactions are similar in order of magnitude, in angular dependence the dissipative back reaction is dominant. It follows that
for the outer, spin-orbit driven $e^-$ outflow to proceed, up to a logarithmic factor of order
\[ \ln \left( \frac{\pi}{2\theta} \right), \]
where $\nu$ is considered at the half-opening angle $\theta$ of the outflow. Consistency with
\[(12)\] gives for the minimum opening angle $\theta_0$ for the outflow in $e^-$
\[ \theta_0 \sim \sqrt{\frac{B_c}{3B}}, \]
up to logarithmic corrections. For $\theta > \theta_0$, the outflow is effectively set by the saturation
limit \[(14),\] whereby the particle luminosity \[(13)\] becomes
\[ L'_p \sim \left( 5 \times 10^{48} \frac{\text{erg}}{\text{sec}} \right) \left( \frac{B}{B_c} \right)^2 \left( \frac{M}{7M_\odot} \right)^2 \sin^2 \theta, \]
where $\theta > \theta_0$ is the half-opening angle of the outflow. This calculation applies formally to
the initial jet. A full calculation of the evolved jet, which consists of combined $e^\pm$-outflow
saturated against dissipative losses in the horizon \(\text{in the sense as described above) falls \outside\ the present scope. Nevertheless, it is expected that the luminosity \[(16)\] remains characteristic for particle-antiparticle outflow in the evolved jet, whose opening angle will
be bounded below by the initial value \[(15).\]

A connection to fireballs \cite{36,37,11} in the theory of $\gamma$-ray bursts \cite{33} is at hand, as \[(16)\]
represents a GRB type of luminosity for $B \sim 10^{16}$G, whose energetics are consistent with
the GRB990123 event, with ultra-relativistic electron-positron outflow along open field-lines
\(\text{supported by the horizon charge } q) in agreement with the input of current fireball models.

The theory of fireballs accounts for the high luminosity in nonthermal emission either by
way of baryon loading or intermittency, to circumvent the thermal emission in the preceding
steady-state models of electron-positron fireballs \cite{17,18}. A small amount of baryonic matter
enables efficient conversion of the energy in the fireball into kinetic energy, with subsequent
shocked emission in interaction with the interstellar medium or wind \cite{39,36}. Some contamina-
tion should be expected, for example, when the interstellar medium or a wind from the
surrounding torus is entrained, and, in case of hypernovae, by interaction with the hydrogen
envelope. On the other hand, intermittency at the source \(\text{see e.g. } \cite{33,35,45}\) will give rise
to unsteadiness in the flow as discussed in the compact fireball model of Eichler & Levinson
\cite{11} with angular variations \(\text{both in } \theta_0 \text{ and in orientation) and internal shocks even when baryon free. Shocks may also result from interactions with collimating baryonic material, perhaps subject to radiative viscosity, whereby a nonthermal component in the emission is produced from a broad range of radii } \cite{11}. \text{These considerations imply an emission spectrum substantially different from that in forementioned steady-state models.}
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