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PENROSE PAIR PRODUCTION AS A POWER SOURCE OF QUASARS AND ACTIVE GALACTIC NUCLEI

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ABSTRACT

In this paper we propose a new mechanism, Penrose pair production in massive, canonical (with $a/M = 0.998$) Kerr black holes, as a power source in quasars, Seyfert galaxies, radio galaxies, and BL Lac objects (i.e., what one usually refers to as active nuclei). As a working hypothesis, we postulate that massive ($10^8\ M_\odot$ or around this value) Kerr black holes reside in the centers of these objects. We also postulate that an accretion disk is formed. In a variety of models, hot, inner disks are expected. If the temperature is sufficiently high—as, for example, in the two-temperature model—then MeV photons enter the ergosphere. The blueshift may boost up the energy of the infalling photons to near GeV values. When this happens, photons can scatter off the tangentially moving protons and produce e^+, e^- pairs. The protons that cross the event horizon give their energy to the ejected pairs (of the order of GeV). This is a Penrose process. Conditions are derived for this mechanism to be important, and it is found that a self-consistent requirement is that these very hot inner disks are also spatially thick. The process can work only if the $r_{\rm mb}$ to $r_{\rm ms}$ target region gets filled up with plasma. This cannot happen in steady-state disks; therefore, Penrose pair production can occur only during periods of instabilities. This may help to explain the variability of the extragalactic compact objects. Moreover, this mechanism may be of profound astrophysical importance in helping to explain the vast energies, synchrotron emission—and variability—of the variable, compact extragalactic sources. Time-dependent calculations of hot, thick disks are needed, but the present work illustrates that such calculations cannot ignore Penrose processes like the one we suggest.

Subject headings: black holes — galaxies: nuclei — quasars

I. INTRODUCTION

Among the most intensely investigated sources of radio emission in the sky are the exceptionally bright and compact radio sources found in the nuclei of active (i.e., Seyfert and radio) galaxies and quasars. In the past few years it has become apparent that the nuclei of Seyfert I galaxies and radio galaxies have X-ray luminosities that are a substantial fraction of their total luminosities (Gursky and Schwartz 1977; Culhane 1978). This is certainly true of 3C 273, and it may also be true for most quasars. It may well be that the primary mechanism for the nuclei of active galaxies and quasars is the X-ray production mechanism (McCray 1978). These compact sources are characterized by flat X-ray spectra, strong optical emission-line spectra, and a compact, variable, radio or millimeter component (Kafatos 1978). Also, in view of the recent observations of M87 (Young *et al.* 1978), it is plausible to postulate the existence of black holes in the centers of active galaxies and quasars. The observations of NGC 6251 (Sargent *et al.* 1978) reinforce this assumption.

In this paper, we study a new physical mechanism, Penrose pair production (PPP), as a way to explain the nature of the vast, fluctuating energy production associated with such active galactic nuclei and quasars. Leiter and Kafatos (1978) (hereafter Paper I) examined PPP for "extreme" Kerr metrics. It was found that in order for this process to occur with the low input energy ratio of $(E_{\rm ph} - \omega l_{\rm ph})/m_p \lesssim 10^{-2}$ (see § II for explanation of the symbols), one must be dealing with an extreme Kerr spinning black hole whose angular momentum density lies in the range $2 \times 10^{-6} \lesssim (1 - a/M) \lesssim 2 \times 10^{-4}$. The reason that such extreme values of angular momentum were needed was because the region where proton targets tend to collect is deep within the ergosphere and must also have a large photon blueshift so that the photon energy E_γ in the "local nonrotating frame" (LNRF) (Bardeen, Press, and Teukolsky 1972) is of the order of the rest mass of the proton, m_p. When this happens, the mass of the Penrose injected protons is converted into the energy of the Penrose ejected electron-positron pairs which subsequently escape from the ergosphere of the black hole. Since it may not be possible for a black hole to spin-up to values of a/M above that of the asymptotic limit $a/M \approx 0.998$ (Thorne 1974), it may be that the extreme values of $1 > a/M > 0.998$ required for this process have to be associated with primordial sources of angular momentum. This is a restriction on the extreme PPP process and suggests that it may be more plausible to study the possibility of PPP occurring for the more commonly expected situation of "canonical" Kerr black holes, with $a/M = 0.998$. Then, in the canonical PPP process, the requirement of $E_\gamma \sim m_p$ can be attained from high-temperature processes in the inner region of the disk itself. This, as we shall see in what follows, is possible in a variety of disk-accretion scenarios. Therefore, the canonical PPP process is of much greater astrophysical interest.

We propose to use this efficient high-energy physical mechanism to offer a unified explanation for the physical properties of radio galaxies, Seyfert galaxies, and quasars. In our picture, the differences between those classes of objects have to do with differences in the masses of the Kerr black holes and differences in accretion rates onto those black holes, as well as different stages in their evolution, rather than different mechanisms for each class of objects.

II. CANONICAL KERR BLACK HOLES AND PENROSE PAIR PRODUCTION

The photon energy E_γ in the LNRF is related to the photon energy $E_{\rm ph}$ and orbital angular momentum $l_{\rm ph}$ at infinity, by the relationship (Paper I)

$$E_\gamma \approx \frac{E_{\rm ph} - \omega l_{\rm ph}}{r - r_+} \approx \left(\frac{8}{2\epsilon^2 - \delta^3}\right)^{1/2}(E_{\rm ph} - \omega l_{\rm ph}) , \qquad (1)$$

where $r_+ = [M + (M^2 - a^2)^{1/2}]^{1/2}$ defines the event horizon and $\omega \sim a/2Mr_+$. Equation (1) deals with a collision occurring at the coordinate $r = M(1 + \epsilon)$, where $\epsilon \ll 1$ close to the event horizon and the Kerr black hole has an angular momentum density $a/M = (1 - \delta^3/4)$ (in Paper I, δ was called α; here, we reserve the letter α for the viscosity parameter). Since we assume canonical black holes (Thorne 1974), $\delta = 0.2$. From now on, even though δ may appear in our formulae, it will be assumed that it has the canonical value.

In Boyer-Lindquist (1967) coordinates, the Kerr metric is given by

$$ds^2 = -e^{2\nu}dt^2 + e^{2\psi}(d\phi - \omega dt)^2 + \exp(2\mu_1)dr^2 + \exp(2\mu_2)d\theta^2 , \qquad (2)$$

where natural units ($c = G = 1$) are assumed. The analytic expression for the metric functions in (2) are given in Bardeen, Press, and Teukolsky (1972). The factor $[8/(2\epsilon^2 - \delta^3)]^{1/2}$ appearing in (1) is nothing more than $e^{-\nu}$, the photon blueshift factor. The targets are located in the $r_{\rm mb}$ to $r_{\rm ms}$ region[1] (Paper I). For these radii, $\epsilon_{\rm ms} \sim \delta$ and $\epsilon_{\rm mb} \sim \delta^{3/2}$. We also define $r_g = GM/c^2 = 1.4822 \times 10^{13} M_8$ cm, where $M_8 = M/10^8\ M_\odot$ and M is the mass of the black hole.

[1] The subscript mb denotes marginally bound; the subscript ms denotes marginally stable.

TABLE 1

METRIC FUNCTIONS FOR A CANONICAL BLACK HOLE ALONG THE EQUATOR

$r_*(GM/c^2)$	$\exp(\mu_1)$	e^ν
1.0632	∞	0
1.0894	17.217	3.159×10^{-2}
1.2	6.324	9.372×10^{-2}
2.	2.004	0.4077
3.	1.501	0.6122
5.	1.250	0.7785
10.	1.111	0.8946
†r_+	1.0632M	
$r_{\rm mb}$	1.0894M	
$r_{\rm ms}$	1.2M	

† Natural units are assumed here, i.e., $c = G = 1$.

In Table 1, we give the values of the factors e^{μ_1}, e^ν of formula (2) for different values of $r_* = r/r_g$ (the radial distance in terms of the gravitational radius r_g). These parameters are needed to compute the blueshift in the ergosphere as well as the proper length (see below). In Table 1, we also give the values of r_+, $r_{\rm mb}$, and $r_{\rm ms}$. All these values are for a canonical Kerr black hole, along the equator of the ergosphere. Note that the ergosphere starts at $r = 2M$ ($r_g = M$ in natural units). According to Shakura and Sunyaev (1973), the relativistic regime starts at $r = 3M$. For radial distances greater than this value, the relativistic corrections are not very large.

In Paper I it was shown that in order to produce a Penrose electron-positron pair in the collision of a proton and a blueshifted photon deep inside the ergosphere (for $r_{\rm mb} \leq r \leq r_{\rm ms}$), the blueshifted photon energy should be $E_\gamma \approx m_p$ (again, natural units are used and m_p is the mass of the proton). This is because the recoiled proton is driven into a "negative energy orbit" (see Bardeen, Press, and Teukolsky 1972); and for this to occur, velocity boosts of about half the speed of light are required in the LNRF. When this happens, the electron-positron pairs escape to infinity with energies of the order of the rest mass of the proton, i.e., about a GeV. Outside the $r_{\rm mb}$ to $r_{\rm ms}$ region, the process becomes inefficient (Bardeen, Press, and Teukolsky 1972). Other Penrose processes might occur; however, for blueshifted photon energies of the order of GeV, Penrose pair production is the dominant process (see Paper I), because the cross section for pair production dominates.

From Table 1, we see that in order for the photon energy in the $r_{\rm mb}$ to $r_{\rm ms}$ region to be ~ 1 GeV, photon input energies $E_{\rm ph}$ in the approximate range 30–100 MeV are required. Such energies are not unreasonable in a variety of accretion disk scenarios. We do not assume a specific source for the accreting mass. It may well be stars that are disrupted by the presence of the hole (see Ipser 1978 and references listed therein).

Here we list those cases which predict hot inner regions from which such photons can be emitted:

a) *The two-temperature model* (Eardley, Lightman, and Shapiro 1975; Shapiro, Lightman, and Eardley 1976). It operates in several regimes. Eardley *et al.* (1978) present two of them: an unsaturated Comptonization regime in which the disk can operate only when there is an external source of copious soft photons; and a self-Comptonized regime in which the disk can operate when it must rely on its own Comptonized bremsstrahlung spectrum. In both regimes of the two-temperature model the ion temperature T_i is much larger than the electron temperature T_e, the latter being in the range of a few billion degrees for the massive black holes considered here. Shapiro, Lightman, and Eardley (1976) find that the ion temperatures are higher than the electron temperatures by factors of 3–300. We emphasize that the models of Shapiro *et al.* are for nonrotating black holes for which relativistic corrections are ignored. Dahlbacka, Chapline, and Weaver (1974) find that spherical accretion onto a nonrotating black hole produces a γ-ray luminosity. The γ-rays arise from the decay of π^0, the latter being produced by ions at $kT_i \gtrsim 100$ MeV. The luminosity computed, however, is too low, due to the very inefficient spherical accretion. Shapiro *et al.* have computed accretion into a rapidly spinning black hole and in a preliminary study find that the γ-ray luminosity L_γ may be as high as the total luminosity in X-rays L_X. Ion temperatures in excess of 100 MeV are obtained. To clarify this, we suggest that it would be extremely important to extend their preliminary results to realistic two-temperature models around Kerr black holes which would include the relativistic corrections.

b) *The optically thin model* (Pringle, Rees, and Pacholczyk 1973; Payne and Eardley 1977). This model operates in two regimes (Eardley *et al.* 1978): a pure bremsstrahlung regime, and a hot relativistic regime. We find that the first predicts temperatures less than $\approx 10^9$ K, too low to produce MeV γ-rays (in the optically thin model $T_i = T_e$). We find that the second regime yields, for the massive black holes considered here, temperatures in the range 2–400 MeV. Modifications are, however, required when relativistic effects are properly taken into account.

c) *The accretion disk corona model* (Liang and Price 1977; Bisnovatyi-Kogan and Blinnikov 1976). A hot corona is assumed to exist around a "standard" disk (Pringle and Rees 1972; Shakura and Sunyaev 1973; Novikov and Thorne 1973). Liang and Price (1977) find temperatures $\sim 10^{11}$ K when the wind dominates in the corona.

These estimates are expected to be higher when a Kerr black hole is considered. Temperatures calculated for nonrotating black holes are expected to become higher when one computes them in the case of rotating black holes due to the higher binding energies attainable from the fact that the accretion disk extends farther inward in the rotating case. This, in fact, is a general argument for higher disk temperatures in cases of rotating black holes. In all three cases listed above we expect the quoted temperatures to be lower limits when realistic accretion models are computed which should include relativistic corrections.

The criteria of applicability of the above three models can be found in Eardley *et al.* (1978). At least one of them applies to the case of massive black holes as long as the ratio $\dot{M}_1/M_8$ is in the approximate range 10^{-3} to 4. Here $\dot{M}_1$ is defined by the relation $\dot{M} = \dot{M}_1 \times (1\ M_\odot\ \mathrm{yr}^{-1})$, where $\dot{M}$ is the mass accretion rate. On the other hand, $\dot{M}_1$ and M_8 are related in such a way that the Eddington limit (Novikov and Thorne 1973) is not violated, i.e.,

$$\dot{M}_1/M_8 \lesssim (6\beta_1)^{-1}, \tag{3}$$

where the luminosity L is related to the accretion rate by the formula

$$L = \beta_1 \dot{M} c^2 \tag{4}$$

and β_1 is the total efficiency of the accretion process which includes the PPP process. We shall call this efficiency the "photon emission efficiency" to distinguish it from the "pair production efficiency," β_2, which describes how efficient the PPP process is., i.e., what fraction of pairs is coming out for each proton falling in. This latter efficiency depends, of course, on the number of infalling γ-rays compared to the number of accreting protons.

From equation (3) it is seen that for reasonable photon emission efficiencies (say $10^{-2} \lesssim \beta_1 \lesssim 0.5$) and as long as the Eddington limit is not violated, it is reasonable to assume that one of the three models (*a*)–(*c*) applies. This possibility and the fact that the gravitational blueshift gives the infalling γ-rays an additional order of magnitude boost in the LNRF—for canonical Kerr black holes—should make the PPP process a likely one in the case of massive black holes that are postulated to exist, and that have been shown to be plausible candidates from the observations of M87, in the centers of active galaxies and quasars.

We emphasize that the PPP process is not tied to any of the models listed above. It is useful, however, to examine this process in the two-temperature model to illustrate the results that would hold in general as long as a significant flux of γ-rays is emitted in the hot inner region close to the ergosphere of a canonical Kerr black hole.

We note that the observations of Cygnus X-1 indicate central temperatures in the disk considerably in excess of 10^9 K (Cunningham 1975), and therefore theory as well as observation suggests high temperatures in the inner regions of accreting disks around Kerr black holes.

III. PPP IN THE INNER ERGOSPHERE OF A MASSIVE, CANONICAL KERR BLACK HOLE

Normally, for a steady-state accretion disk surrounding a canonical Kerr black hole, the plasma moves in slowly in Keplerian fashion, until it reaches r_{ms}. The inner region of the disk ends here and matter quickly free-falls into the black hole; hence, the density of matter builds to a maximum at r_{ms} and then falls quickly to zero inside this radius (Novikov and Thorne 1973).

The situation is very different if instabilities develop in the disk because, in that event, large amounts of plasma may attain temporary Keplerian orbits within the region between r_{ms} and r_{mb}. These instabilities could be secular (Lightman and Eardley 1974; Lightman 1974), thermal (Pringle, Rees, and Pacholczyk 1973), or may be associated with a disruption of the flow pattern of the disk. In any case, we expect that if the angular momentum of the infalling gas is small, but not zero, it will enter the r_{mb} to r_{ms} region with highly eccentric orbits and it will gradually fall in.

As in Paper I, we expect that the density will peak around r_{ms} and then, during instabilities, will remain constant or at least be slowly varying down to the r_{mb} radius (see below). Under these conditions, it is plausible to expect that PPP will occur within the r_{mb} to r_{ms} target region; and, since blueshifted photons with $E \approx m_p$ are also present, the required boost of $c/2$ will be given to the protons.

We have estimated the optical depths for various processes outside and inside the ergosphere. The optical depth of any process is given by

$$\tau = \int_{r_2}^{r_1} \kappa\rho \exp(\mu_1) dr\,, \qquad (5)$$

where κ is the appropriate opacity, ρ is the proper matter density, and $\exp(\mu_1)dr$ is the proper radial distance element, where $\exp(\mu_1)$ can be found in Table 1 for various radial distances. As in Paper I, we assume a radially infalling photon scattering off a tangentially moving proton. The limits are referring to the region(s) under consideration. It is seen from (5) that, in order to estimate the optical depths, one needs the values of the density ρ.

We estimate the density ρ from the relation (Novikov and Thorne 1973)

$$\dot{M} = 2\pi r \Sigma v_r \mathscr{D}^{1/2}\,, \qquad (6)$$

where $\dot{M}$ is a constant independent of r and t, Σ is the surface density of the disk,

$$\Sigma = 2\int_0^h \rho dz\,,$$

and h is the half-thickness of the disk, v_r is the radial drift velocity with which the disk is moving in, and $\mathscr{D}$ is the quantity $1 - 2/r_* + a_*^2/r_*^2$ where $a_* = a/M$, $r_* = r/r_g$. Assuming a constant density with height z, we estimate the number density of matter ($n = \rho/m_p$) from

$$n = \dot{M}/(4\pi r h m_p v_r \mathscr{D}^{1/2})\,. \qquad (7)$$

Both h and v_r are functions of the viscosity parameter α for shear stress $= \alpha P$ ($0.01 \lesssim \alpha \lesssim 1$) (Shakura and Sunyaev 1973) and therefore highly uncertain. Moreover, equations (6) and (7) are strictly correct only for steady-state disks. In the absence of a correct time-dependent theory with hot inner accretion disks around Kerr black holes, equations (6) and (7) can only provide estimates. These estimates should be fairly accurate in between instabilities when the disk is in quasi-steady situation.

a) *Outside the Ergosphere*, $2 \lesssim r_* \lesssim 3$

The half-thickness is estimated to be (cf. Shakura and Sunyaev 1973; Shapiro, Lightman, and Eardley 1976) $h \approx 0.8r$, i.e., we have a spatially thick inner region; the drift velocity v_r is estimated to be $v_r \approx \alpha v_c$, where $v_c = c r_*^{-1/2}$ is the Keplerian velocity (assuming a sound speed appropriate for dominant gas pressure, valid in the two-temperature model). The matter number density is therefore

$$n(3) \approx (\dot{M}_1/M_8{}^2)\alpha^{-1} r_*^{-3/2} \times 2 \times 10^{12}\ \mathrm{cm}^{-3}\,, \quad (8)$$

and the corresponding optical depth for scattering of photons off targets (protons for PPP, electrons for Compton scattering) with a scattering cross section $\sigma = \sigma_{26} \times 10^{-26}\ \mathrm{cm}^2$ is

$$\tau(3) \approx (\dot{M}_1/M_8)(0.1/\alpha)\sigma_{26}\,. \qquad (9)$$

The index (3) in the above equations refers to the region $2 \lesssim r_* \lesssim 3$ which is appropriate for estimates *outside* the ergosphere.

For electron Compton scattering, we find that for photon energies less than 1 MeV ($\sigma_{26} \approx 66$),

$$\tau_{es}(3) \approx \mathscr{R}/(\alpha\beta_1)\,, \qquad (10)$$

where $\mathscr{R} = L/L_{Ed}$ and L_{Ed} is the Eddington luminosity. For reasonable emission efficiencies, $10^{-2} \lesssim \beta_1 \lesssim 0.5$, $\tau_{es} \gg 1$ (assuming $\alpha \approx 0.1$) as long as the luminosity is not much smaller than the Eddington limit. Since, however, the free-free optical depth $\tau_\nu{}^{ff}$ is very small, the effective optical depth τ_* (Novikov and Thorne 1973) $\tau_* \equiv [\tau_{es}\tau_\nu{}^{ff}]^{1/2} \ll 1$, i.e., the region is optically thin as expected.

For pair production off protons at energies $E_{ph} \approx 100$ MeV, we find

$$\tau_{pp}(3) \approx 10^{-2}\mathscr{R}/(\alpha\beta_1)\,. \qquad (11)$$

This estimate shows that any high-energy γ-rays outside the ergosphere will penetrate it without any appreciable pair production; i.e., we find $\tau_{pp}(3) \lesssim 1$, again assuming $\alpha \approx 0.1$, and the optical depth is smaller the farther away the luminosity is from the Eddington limit. These estimates also hold for lower (say down to $\sim$1 MeV) photon energies.

b) *Inside the Ergosphere*

We consider two regions, the *outer* ergosphere ($r_{ms} \lesssim r \lesssim 2r_g$), and the *inner* ergosphere or target region ($r_{mb} \lesssim r \lesssim r_{ms}$). We use the index (2) to refer to the first and (1) to the second.

We estimate the optical depths and number density from (5) and (7) as before. The requirement that PPP is important in the inner ergosphere [i.e., $\tau_{pp}(1) > 1$] implies that the accretion disk should be spatially fairly thick (say $h \approx 0.2r$). We find that since $v_r \to c/2$ as $r \to r_{mb}$, then

$$\langle h/r \rangle \lesssim 0.2\sigma_{26}\mathscr{R}/\beta_1 \,. \qquad (12)$$

This is likely to happen when the inner region gets filled up with plasma as a result of instabilities. The density should be

$$n(1) \gtrsim 5 \times 10^{12}(M_8\sigma_{26})^{-1}\ \mathrm{cm}^{-3} \,. \qquad (13)$$

Note that $\sigma_{26} \approx 1$ for GeV, blueshifted, photons.

On the other hand, to have an efficient PPP process in the target domain, the photons should travel through the outer ergosphere with relatively few pair production scatterings, i.e., $\tau_{pp}(2) \lesssim 1$. Assuming comparable average ratios of the quantity h/r in the inner and outer ergosphere, we find the requirement for this to be

$$v_r(2) > c/8 \,. \qquad (14)$$

Since $v_r \approx c/15$ in the region just outside the ergosphere (see discussion of case [*a*] above) and since $v_r \to c/2$ close to the r_{mb} radius, condition (14) is certainly a very reasonable one.

Inside the ergosphere, electron scattering becomes less important and pair production becomes progressively more important as one approaches the inner ergosphere. Other processes are not important. For instance, pair production through two-photon scattering is never important, even at MeV energies, because the number of γ-rays is less than the number of protons (or electrons) by over an order of magnitude (see below). However, two-photon pair production may be important in the accretion disk outside the ergosphere.

We conclude that pair production becomes progressively more important as the target region between r_{ms} and r_{mb} is approached, finally becoming Penrose pair production in this inner ergosphere. Relations (8) and (11) and conditions (12), (13), and (14) are all certainly reasonable. The fundamental equation (7) is strictly correct only for steady-state disks, but it should provide a useful estimate particularly during periods of rearrangement of the flow pattern in the disk. As we see from the ratio of (13) to (8), the density steadily builds up as $r \to r_{ms}$, while the drift velocity steadily builds up from a fraction of the speed of light outside the ergosphere, to half the speed of light in the inner ergosphere as PPP becomes important. We emphasize that the requirement that the inner disk (or, for that matter, corona) should be of appreciable spatial thickness, is certainly self-consistent for these very hot inner regions of accretion disks.

In Figure 1, we show how PPP operates in the inner ergosphere of a canonical Kerr black hole, assuming the two-temperature model applies. Gamma-rays (γ) produced from the very hot protons ($T_i \sim 5 \times 10^{11}$ K) enter the ergosphere. Radially moving blueshifted γ-rays (γ') scatter off the tangentially moving protons in the r_{mb} to r_{ms} region. When the protons get injected through the horizon r_+ after such scatterings, the resultant pairs are ejected predominantly within 40° of the equator. In this figure, as in the entire paper, we assume that the infalling plasma is in direct orbits (Bardeen, Press, and Teukolsky 1972). The ergosphere is shown to scale. The dashed line indicates the boundary of the inner, very hot, accretion disk. No other scattering processes (e.g., electron scattering) are shown here.

IV. DISCUSSION AND CONCLUSIONS

The Penrose ejection of the PPP pairs is similar to the extreme case examined in Paper I. They are ejected within a 40° angle relative to the equator. Their energy is of the order of 1 GeV per pair. Thus the relativistic electrons seen in active galaxies and quasars through their synchrotron emission are naturally accounted for in our picture. The pressure of the electron-positron relativistic gas is proportional to their number density. The latter is estimated from the condition that away from the ergosphere (where the relativistic effects are not important) the number of outgoing pairs per unit time cannot be greater than the number of infalling protons per unit time through a spherical surface of radius r; equality holds only in special situations and then both emission and production efficiencies are close to 1, i.e., when PPP is the dominant power source of the active nucleus. The relativistic pairs have a pressure 10^2–10^3 times as great as the thermal electrons and comparable to the ions for efficiencies $\beta_2 \sim 1$. It is therefore seen that, for production efficiencies as low as 10^{-3} to 10^{-2}, the ejected pairs can have a strong influence on the extended hot inner disk, or corona around a standard disk. The escaping pairs will tend to disturb the disk as they leave the ergosphere and will enhance the possibility of further instabilities in the disk, thus starting the cycle all over again.

This suggests that PPP may come in bursts, with the duration of the burst being determined by the collapse time of the inner disk. The hydrodynamical time is expected to be longer than the free-fall time (unless spherical accretion persists—in which case PPP would then be inefficient) but shorter than the drift time for steady-state disks, when distances applicable to the inner disk are considered. The exact details of this burst time would require calculations with a detailed hydrodynamic model which would include a specific mechanism for the viscosity of the plasma (Eardley and Lightman 1975). In addition to the characteristic burst time, one might expect smaller fluctuations to be superposed on the burst due to local instabilities associated with the optically thin inner region of the disk. These shorter time scale variations could be associated with thermal instabilities, or with varying conditions, say varying densities, in the inner hot disk. Finally, there would be an overall lifetime for the active nucleus which would depend on the time required to exhaust the disk of its matter content. This time depends on the particular scheme to provide the accreting mass and has to do with the environment of the black hole. This last lifetime is hardest to estimate. However, the temporal scenario seems to qualitatively fit observed data from variable radio sources (Dent and Hobbs 1973; Hobbs and Dent 1977).

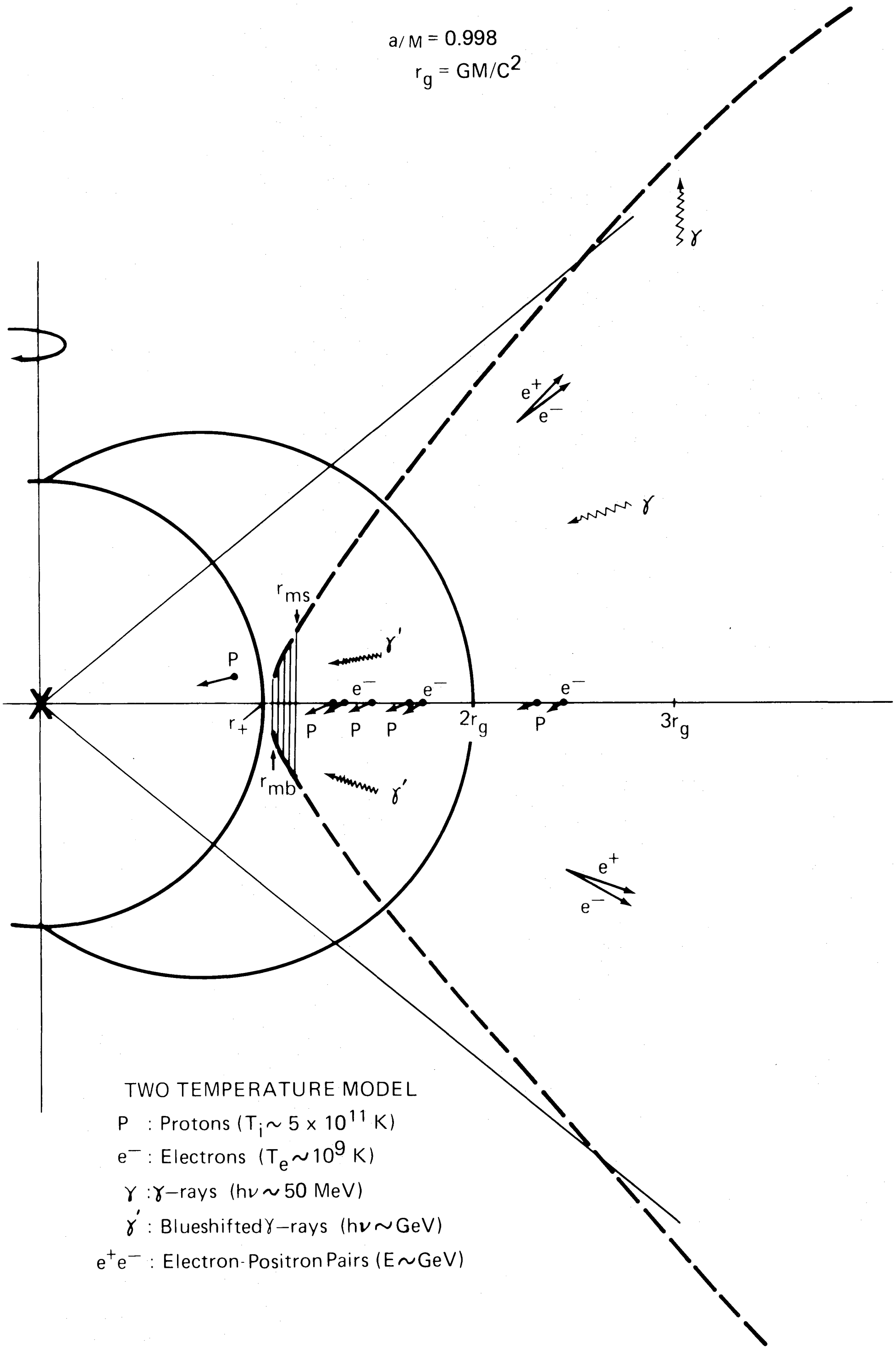

FIG. 1.—A possible PPP scenario. The ergosphere of a canonical black hole is shown to scale, with the event horizon (r_+), the r_{mb} to r_{ms} target region, the outer boundary of the ergosphere at $2r_g$, where r_g is the gravitational radius. Gamma rays (γ) produced from the very hot protons ($T_i \sim 5 \times 10^{11}$ K) expected in the two-temperature model enter the ergosphere. Radially moving blueshifted gamma rays (γ') scatter off the tangentially moving protons in the target region. When the protons get injected through the event horizon, the resultant e^+, e^- pairs are ejected predominantly within 40° of the equator. Direct orbits are assumed for the infalling plasma. The dashed line indicates the boundary of the inner, very hot, accretion disk. Note that PPP would operate in other hot accretion disk models.

While more detailed calculations are required to test the viability of the PPP process as a dominant power source of compact extragalactic sources, the present work shows that such Penrose processes cannot be ignored in realistic calculations for spatially thick accretion disks. A theory for thick accretion disks has not yet been worked out (Bardeen 1973); but when an attempt is made to solve this very challenging problem, PPP should certainly be included.

One test of our theory is the observation of resultant pair-annihilation γ-rays (0.5 MeV), which would be generated by the pairs when they are annihilated.

This γ-ray intensity should be higher the more active the source; the more pairs that are produced, the easier it is for them to escape the inner hot region and eventually be annihilated; whereas, if few pairs are produced, they would suffer severe synchrotron and Compton losses in the inner disk and their annihilation radiation would be masked by the γ-rays emitted by the disk itself. The detection of the 0.5 MeV line from the direction of the galactic center (Leventhal, MacCallum, and Stang 1978) is consistent with our theory. Forthcoming γ-ray detection experiments with the *HEAO* would be extremely useful to see if the high-energy bremsstrahlung γ-rays associated with the GeV pairs radiating in the disk could be observed as well.

It has been shown that PPP may provide a common mechanism to explain the intense activity in quasars and active galaxies. Assuming that the entire luminosity arises from the accreting matter (which includes the PPP process but is not limited to it), the Eddington limit (eq. [3]) relates the accretion rate to the mass of the black hole (see eqs. [3] and [4]). We find that PPP is an important process that may help to explain the power requirements, variability, and energies of relativistic electrons that are seen in the variable, compact extragalactic objects. This process operates as long as a copious supply of high-energy γ-rays is emitted from the inner disk of a canonical Kerr black hole. This requirement is consistent with the requirement (see relation [12]) that this inner region is spatially thick, since in certain accretion models (e.g., the two-temperature model) spatially thick, very hot inner disks are predicted. In our picture the differences between the various active galaxies and quasars are explained by differences in masses of the black holes, accretion rates, and general environment in the vicinity of the black hole. More calculations are needed to examine this process in realistic, self-consistent accretion disk scenarios. We believe that a search of 0.5 MeV γ-rays should be undertaken to try to detect this radiation from the most active extragalactic compact objects.

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