High-energy gamma radiation from a black hole: slow accretion with friction

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We consider the production of high-energy gamma rays ($E_\gamma \sim 100$ MeV) near a black hole through the collision of protons in stable orbits in an accretion disk. Spherical accretion at a slow pace yields a self-consistent model. When the accretion rate is $\dot{M} \lesssim 10^{-3}$ of the Eddington limit, proton orbits in the gravitational field of the black hole contract, due to friction with the hot electron gas. We refer to this type of motion as *accretion with friction*. Gamma radiation results from the decay of neutral pions created by collisions between relativistic protons moving in the vicinity of stable orbits. We conclude by showing that this mechanism can give rise to a detectable gamma flux from both Galactic and extragalactic sources.

1. INTRODUCTION

Gamma rays generated through accretion onto a black hole have been thoroughtly discussed in the literature, beginning with the pioneering work of Shvartsman.[1] Specifically, high-energy gamma rays can be produced by the decay of neutral pions, which in turn are the products of nuclear interactions in the accretion flow.[2] The most convincing calculations have been based on a two-fluid model of the gas,[3–5] in which the temperature of the protons is much higher than that of the electrons, reaching perhaps $T_p \sim 10^{12}$ K. This situation arises if the characteristic time for proton energy losses exceeds the time it takes before they fall into the black hole, and the proton and electron components of the gas are thermally decoupled. Gamma-ray production in a high-proton-temperature accreting gas has been examined by a number of authors,[6–9] as has radiation produced in a strong magnetic field for the case of spherical accretion.[10–12]

The work of Meszaros and Ostriker,[8] in which they considered the formation of a standing shock wave around a black hole, is of particular interest as it relates to our approach. In their model, gas flows into the hole both radially and in a laminar stream right down to the accretion radius r_a. A spherical shock wave forms where the radial velocity exceeds the speed of sound. The location of the shock front can vary from a few gravitational radii out to hundreds, depending on the accretion rate and the gas temperature. In crossing the shock front, protons have their direction of motion randomized, and along the flow, their temperature rises to close to the gravitational temperature $T_g \sim 10^{12}$ K. It is assumed here that rapid dissipation keeps the small-scale magnetic field weak.

When the black hole is rotating, the efficiency with which high-energy gamma rays are produced is even higher. This is largely a result of the fact that the boundary of the region of stable motion in the Kerr metric is closer to the horizon than in the Schwarzschild metric. Motion near the limit of stability is therefore more emphatically relativistic, and the temperature of the proton component can be higher. Gamma radiation from the relativistic neighborhood of a rotating black hole is dealt with in Refs. 13–18. In a thick accretion disk around a rotating black hole, pions can be produced by collisions in the ergosphere between protons moving in nonequatorial orbits.[19]

Proton acceleration at a shock front near a black hole was proposed by Protheroe and Kazanas[20] to be the mechanism responsible for generating gamma rays in the nuclei of active galaxies through proton-proton collisions, and that idea was subsequently expanded upon.[21–24]

We also note here several mechanisms for generating gamma rays that are not directly connected with p-p interactions. These include synchrotron emission, inverse Comptonization of low-energy photons, and bremsstrahlung in relativistic jets emanating from compact courses.[25–31] This comprises far from a complete list of previous research on accretion-induced production of high-energy gamma rays.

In the present paper, we examine a model for low-rate spherical accretion onto a nonrotating black hole. That the black hole is nonrotating is not a fundamental limitation, but merely facilitates some of the lengthier calculations. All conclusions arrived at in this paper are valid for a rapidly rotating black hole as well. When the accretion rate $\dot{M}$ is low enough, proton energy losses are small, and their radial approach to the hole is a leisurely one along almost circular spiral trajectories. On the other hand, because of efficient Coulomb scattering from protons and other energy loss mechanisms, the electrons are thermalized and comprise a Maxwellian gas at temperature T_e, which is much lower than the effective kinetic temperature T_p of the protons in the relativistic region surrounding the black hole. The electrons move radially toward the black hole. Under these flow conditions, there is no chance for a shock wave to form. The proton distribution, which preserves the electrical neutrality of the plasma, stabilizes perturbations in the electron component of the gas associated with spatial density fluctuations. Neutral pions and gamma rays are then produced by interactions between protons in intersecting stable orbits.

2. A MODEL FOR SPHERICAL ACCRETION WITH FRICTION

In this section, we shall consider spherical accretion onto a nonrotating black hole; the accreting material is assumed to consist entirely of hydrogen. Distances will be measured in units of

$$r_g = GM/c^2 \approx 1.47\cdot 10^5 (M/M_\odot)\ \text{cm}, \tag{1}$$

where G is the gravitational constant, M is the mass of the black hole, and $M_\odot$ is the mass of the sun. The fundamental parameter driving this model is the dimensionless accretion rate

$$\dot{m} = \dot{M}/\dot{M}_{cr}, \tag{2}$$

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where the critical accretion rate

$$\dot{M}_{cr}\approx 1\cdot 10^{18}(M/M_\odot)\ \text{g/sec} \qquad (3)$$

corresponds to the Eddington luminosity

$$L_E=\alpha_0\dot{M}_{cr}c^2,$$

if mass is converted into energy with an efficiency $\alpha_0 = 0.1$. The regime that we treat here is characterized by a small value of $\dot{m}$; the density $n(r)$ of the accreting gas is low enough that the characteristic time for proton energy losses

$$\tau_p=(E^{-1}dE/dt)^{-1}$$

is much greater than the time for one orbital revolution $\tau_{\text{rev}} = 2\pi R/v$, where R is the orbital radius and v the orbital velocity.

The principal source of proton energy loss is the p-e Coulomb interaction, which results in small changes in proton energy, and behaves much like friction. Under these circumstances, the protons spiral into the black hole along practically circular trajectories, remaining largely in regular orbits.

In contrast, electrons are efficiently scattered by the Coulomb field of the protons, and experience sizable energy losses, forming a isotropic Maxwellian gas at temperature T_e as a result. The friction of protons against the high-temperature electron gas provides the mechanism that makes them spiral into the black hole.

To begin with, let us detail the energy losses suffered by particles in the present context.

1. For protons undergoing p-e Coulomb interactions in a hot gas at temperature T_e, and with initial energies

$$^3/_2T_e<E_p<(m_p/m_e)3/2T_e$$

the losses[32] are[1)]

$$(dE_p/dt)_e=1.6\cdot 10^{-16}(m_e/T_e)^{1/2}n_eE_p, \qquad (4)$$

where n_e is the electron density in cm^{-3}, and the numerical coefficient corresponds to c^{-1}.

2. Bremsstrahlung energy losses of electrons with kinetic energy $E_e = \frac{3}{2}kT_e$ are given by

$$(dE_e/dt)_{br}=8.2\cdot 10^{-11}(T_e/m_e)n_p,\quad T_e\ll m_e, \qquad (5)$$

$$(dE_e/dt)_{br}=1.2\cdot 10^{-10}(T_e/m_e)\Phi_r(T_e)n_p,\quad T_e\gg m_e. \qquad (6)$$

The numerical coefficients in Eqs. (5) and (6) are in units of eV/sec, and spatial densities are in cm^{-3}; $\Phi_r(T_e)$ is a tabulated function, some values of which are

T_e/m_e	$^2/_3$	$^4/_3$	$^{10}/_3$	$^{20}/_3$
Φ_r	0.73	1 0	1.6	1.9

3. Electron synchrotron losses (for electron energy equal to the mean) are

$$(dE_e/dt)_s=2\cdot 10^{-3}(T_e/m_e)H^2,\quad T_e/m_e\ll 1, \qquad (7)$$

$$(dE_e/dt)_s=1.5\cdot 10^{-3}(T_e/m_e)^2H^2,\quad T_e/m_e\gg 1, \qquad (8)$$

where H is the magnetic field strength in gauss, and the numerical coefficients are in eV/sec.

In the present model, in which the protons are not in thermal equilibrium with the electrons, the only losses that can be neglected are those due to p-p Coulomb scattering. In relativistic orbits close to the limit of stable motion, which lies at $r = 6r_g$, p-p collisions can lead to pion production. These energy losses quickly become catastrophic: Proton energies drop below the minimum needed to support stable motion, and the protons fall into the black hole. Only for a rapidly rotating black hole can one of a pair of protons colliding in the ergosphere manage to escape to infinity, by virtue of the Penrose process.

Hereafter, we limit our attention to the relativistic region of stable motion $6r_s < r < 30r_g$ and we assume that proton energy losses are small enough that the protons can move along regular spiral trajectories about the black hole. The self-consistency of this scenario depends on a criterion that we shall calculate below. First let us calculate the radial dependence of the gas density, bearing in mind throughout that electrical neutrality requires that $n_e(r) = n_p(r) \equiv n(r)$.

Making use of the expression for the velocity of a proton in a circular orbit relative to a locally stationary observer[33]

$$\beta(r)\equiv v(r)/c=[r_g/(r-2r_g)]^{1/2}, \qquad (9)$$

we find the kinetic energy of such a proton to be

$$E_p(r)=[(r-2r_g)^{1/2}/(r-3r_g)^{1/2}-1]\,m_p. \qquad (10)$$

Substituting (10) into (4), we obtain an equation relating $r(t)$, $\dot{r}(t)$, and $n(r)$. Taking into account the stationarity of the accretion flow

$$n(r)=\dot{M}/4\pi m_pr^2\dot{r}, \qquad (11)$$

we obtain the radial dependence of the effective radial velocity v_r and gas density n:

$$v_r=u_0\frac{r_g}{r}\left(\frac{r}{r_g}-3\right)\left(\frac{r-2r_g}{r-3r_g}\right)^{1/4}\left[\left(\frac{r-2r_g}{r-3r_g}\right)^{1/2}-1\right]^{1/2}, \qquad (12)$$

$$n(r)=n_0\frac{r_g}{r}\left(\frac{r}{r_g}-3\right)^{-1}\left(\frac{r-2r_g}{r-3r_g}\right)^{-1/4}\left[\left(\frac{r-2r_g}{r-3r_g}\right)^{1/2}-1\right]^{-1/2}, \qquad (13)$$

where

$$u_0=\tilde{u}_0\dot{m}^{1/2}(m_e/T_e)^{3/4}. \qquad (14)$$

$$n_0=\tilde{n}_0\dot{m}^{1/2}(M/M_\odot)^{-1}(T_e/m_e)^{3/4} \qquad (15)$$

with constants $\tilde{u}_0 = 1.0\cdot 10^{10}$ cm/sec and $\tilde{n}_0 = 2.2\cdot 10^{20}$ cm^{-3}. Once having solved (12) and (13), it is easy to check that the scattering of electrons from protons effectively isotropizes the electron gas.

The electron gas is heated by proton energy losses induced by p-e Coulomb interactions, and it is cooled by a combination of bremsstrahlung and synchrotron emission. The time scales for these processes are much shorter than the time it takes the electron gas to relocate in the radial direction, and the temperature T_e must therefore be found by equating the electron bremsstrahlung or synchrotron loss rates to the proton Coulomb losses. Neglecting all electron energy losses but those due to bremsstrahlung, we obtain the temperature T_e as a function of radius r from Eqs. (4)–(6):

r/r_g	6	8.3	3.7	13.1
T_e/m_e	6	4.8	4.5	3.9

This is the limiting case for very high temperature. All other cooling channels, including plasma wave generation, will lower the temperature. The opposite low-temperature limiting case can be relevant when there is a very strong magnetic field present. Assuming equipartition of energy to obtain a rough estimate of the strength of such a field, $H^2/8\pi \sim n(r)T_e$, we may use (4), (7), and (8) to obtain the following temperatures:

r/r_g	6	8.3	9.7	13.1
T_e/m_e	0.67	0.59	0.55	0.43

The closer one is to the black hole, the stronger the magnetic field, and at the last stable orbit $r = 6r_g$, it is

$$H\approx 2\cdot 10^7 \dot{m}^{1/4}(M/M_\odot)^{-1/2} \text{ G}.$$

This would appear to be an overestimate. In actuality, the magnetic field must be much weaker than implied by equipartition (see also the discussion in Refs. 6 and 8). An inhomogeneous magnetic field is induced by turbulence in the electron gas. But since the plasma is electrically neutral, the spatial distribution of protons in regular orbits about the black hole will stabilize the electron gas and suppress instability associated with large displacement. Turbulent motions of the electron gas are therefore bounded, and the associated magnetic field will be weak. The foregoing temperature estimates may be considered upper and lower bounds.

For the next step, we must show that the equilibrium density of positrons will be much lower than the gas density $n(r)$. This is basically a consequence of the relatively low surface density of gas around the black hole,

$$N=\int_{6r_g}^{\infty} n(r)\,dr\approx n_0 r_g \approx 3.2\cdot 10^{25}\dot{m}^{1/2}\left(\frac{T_e}{m_e}\right)^{3/4} \text{cm}^{-2}. \quad (16)$$

In the present instance, positrons are created via two processes:

1. $e+p\to p+e+e^-+e^+$,
2. $\gamma+\gamma\to e^+ + e^-$.

Positron annihilation can be neglected, inasmuch as their equilibrium density is maintained, on the one hand, by production processes, and on the other, by the advection of positrons in the accretion flow into the black hole. As a result, we obtain a simple equation for the equilibrium positron density:

$$\partial[r^2 v_r(r) n_{e^+}(r)]/\partial r = Q_{e^+}(r) r^2,$$

where $v_r(r)$ is given by Eq. (12), and $Q_{e^+}(r)$ is the number of positrons created per cubic centimeter per second. After some simple manipulations, we have for process 1

$$\frac{n_{e^+}(r)}{n(r)}\approx 2\frac{c}{\tilde{u}_0}\sigma_{br}(T)\tilde{n}_0{}^3 r_g\left(\frac{T_e}{m_e}\right)^{3/2}\ln\frac{R}{r}\approx 5\cdot 10^{-4}\left(\frac{T_e}{m_e}\right)^{3/2},$$

and for process 2,

$$\frac{n_{e^+}(r)}{n(r)}\approx\frac{2}{3}\dot{m}\frac{c}{\tilde{u}_0}\sigma_{\gamma\gamma}(T)\sigma_{br}{}^2(T)\tilde{n}_0{}^3 r_g{}^3\left(\frac{T_e}{m_e}\right)^3\frac{6r_g}{r}\ln^2\left(\frac{r}{3r_g}-1\right)$$
$$\approx 3\cdot 10^{-6}\left(\frac{\dot{m}}{10^{-3}}\right)\left(\frac{T_e}{m_e}\right)^3\frac{6r_g}{r}\ln^2\left(\frac{r}{3r_g}-1\right).$$

Here $\sigma_{br}(T)$ and $\sigma_{\gamma\gamma}(T)$ are the cross sections for processes 1 and 2 respectively, and R is the distance at which the temperature drops below the pair-creation threshold. For numerical estimates, we have taken $\sigma_{br}\approx\alpha r_e^2\approx 6\cdot 10^{-28}$ and $\sigma_{\gamma\gamma}=\frac{3}{16}\sigma_T\approx 1\cdot 10^{-25}$ cm^2 (σ_T is the Thomson cross section).

Finally, let us estimate the magnitude of the dimensionless accretion rate $\dot{m} = \dot{M}/\dot{M}_{cr}$, which when small ensures the self-consistency of our model. The latter requirement stems from the necessity of the orbital period of a proton

$$\tau_{rev}=2\pi r/v(r)=2\pi r_g c^{-1}\left(\frac{r}{r_g}-2\right)^{1/2}\frac{r}{r_g}$$

being small by comparison with the characteristic energy-loss time τ_p dictated by Eq. (4):

$$\frac{\tau_p(r)}{\tau_{rev}(r)}\approx 0.6\dot{m}^{-1/2}(T_e/m_e)^{3/4}f(r/r_g), \quad (17)$$

where

$$f(x)=1.59(x-2)^{1/4}\left(\frac{x-3}{x-2}\right)^{1/4}\left[\left(\frac{x-2}{x-3}\right)^{1/2}-1\right]^{1/2}$$

is a slowly varying function of its argument $x = r/r_g$, which is equal to 1 at $x = 6$ and 1.11 at $x = 100$. The requisite condition $\tau_p/\tau_{rev}\gg 1$ is satisfied when $\dot{m}\lesssim 1\cdot 10^{-3}$.

Will the accretion flow in our present problem be stable? The most natural factor inducing instability, a shock wave, cannot come into being in the region with which we are concerned, $r\lesssim 10r_g$. The electron gas in that region is relativistic, and the speed of sound there, $u_s\sim c/\sqrt{3}$, is much greater than the radial velocity of the bulk gas flow. In the zone $6 < r/r_g < 10$, the latter, according to Eq. (12) is

$$v_r\approx 3\cdot 10^7(\dot{m}/10^{-3})^{1/2}(m_e/T_e)^{3/4}\,\text{cm}\cdot\text{s}^{-1}.$$

These considerations will remain unaffected even if a shock wave does become established at large radii. In point of fact, the shock wave does distort the proton flow, but because of the collisionless nature of our model, the protons will be captured into regular stable orbits anew.

We expect that the gravitationally controlled motion of the protons will exert a stabilizing influence on the electron gas and prevent spatial instabilities from developing. The growth of electron-gas density fluctuations is restrained by the overall electrical neutrality of the plasma and by the fixed density of the accreting protons. Plasma waves, should they develop, will lead only to a slight cooling of the electron gas.

3. THE GENERATION OF GAMMA RAYS

One significant feature of the present model is the motion of the protons along regular, stable trajectories. As a result, all protons with velocity $\beta = v(r)/c$ will wind up in orbits at a radius $r = r_g(2+\beta^{-2})$, i.e., they will all follow trajectories that lie on a sphere of radius r centered on the black hole. At any point O on the sphere (see Fig. 1), proton trajectories will be uniformly distributed in azimuth over the range $(0, 2\pi)$. The probability per second of colliding with a proton is

$$W(r)=2\int_0^{\pi} d\theta\, n(r,\theta)\sigma(v)v(1-\beta^2\cos\theta), \quad (18)$$

where $n(r,\theta) = n(r)/2\pi$. Near the pion production threshold, we use Dermer's[34] approximation for the reaction cross section of $p+p\to p+p+\pi^0$, which is valid up to $p_{max} = 1.3$ GeV/c:

$$\sigma_{\pi^0}(p)=\sigma_0(p-p_{th})^{3.2}, \quad (19)$$

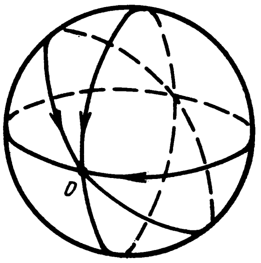

FIG. 1. All stable trajectories passing through the point O form a sphere centered on the black hole.

where p is the proton momentum in the rest frame of the second proton, $p_{th} = 0.8$ GeV/c, and $\sigma_0 = 32.6$ mb. The cross section (19) takes multiple pion production into account. The momentum p is related to the angle θ by

$$\cos\theta = \frac{1}{\beta^2}\left[1-(1-\beta^2)\left(\frac{p^2}{m_p{}^2}+1\right)^{-1/2}\right]. \qquad (20)$$

The total number of gamma rays produced per second by π^0 decays is

$$Q_\gamma = \eta \int_{r_{min}}^{r_{max}} 4\pi n(r) W(r) r^2\, dr, \qquad (21)$$

where $\eta \approx 0.9$ is a factor that makes allowance for the fact that some of the gamma rays are trapped by the black hole; $n(r)$ is given by (13), while r_{min} and r_{max} delimit the effective creation region for the π^0. When we integrate Eq. (21) over r, each proton is counted twice, but this is counterbalanced by the creation of two gamma rays in the decay of each π^0. Further simple calculations yield

$$Q_\gamma = 8\eta \dot{m} \sigma_0 c \tilde{n}_0{}^2 r_g{}^3 \left(\frac{M}{M_\odot}\right)\left(\frac{T_e}{m_e}\right)^{3/2} \int_{0,36}^{0,5} \frac{d\beta}{1-(1-\beta^2)^{1/2}}$$

$$\times \int_{0,78}^{p_{max}} \frac{p\,dp}{(1-\cos^2\theta)^{1/2}} \frac{(p-0,78)^{3,2}}{(p^2+1)^{1/2}} (1-\beta^2\cos\theta), \qquad (22)$$

where $p_{max} = 2\beta(1-\beta^2)^{-1}$, $\cos\theta$ is given by Eq. (20), and we have put $m_p = 1$ GeV. The numerical value of (22) is

$$Q_\gamma \approx 1.17\cdot 10^{40} (\eta/0.9)\, \dot{m}\, (M/M_\odot)\, (T_e/m_e)^{3/2}\ \text{sec}^{-1}. \qquad (23)$$

The gamma-ray energy spectrum corresponds to the decay of π^0 at rest; that is, $dQ_\gamma(E_\gamma)/dE_\gamma$ is constant up to $E_{max} \approx 70$ MeV.

We now wish to discuss the possibility that gamma rays are absorbed at the source. Equation (16) implies that there is little absorption in the gas. Self-absorption due to $\gamma + \gamma \to e^+ + e^-$ is also neglible, which can be shown by assuming that a monochromatic stream of gamma rays is emitted by a sphere of radius $R = 6r_g$ at a rate Q_γ. The absorption probability, i.e., the number of collisions between photons emitted radially from the surface of the sphere and participating in the process $\gamma + \gamma \to e^+ + e^-$, can be written in the form

$$\nu = \frac{Q_\gamma}{2\pi c R^2} \int_R^{r_{max}} dr \int_{(1-R^2/r^2)^{1/2}}^{\cos\theta_{min}} d\cos\theta\,(1-\cos\theta)\,\sigma(E_\gamma, \cos\theta), \qquad (24)$$

where $\sigma(E_\gamma, \cos\theta)$ is the cross section for two photons of energy E_γ to collide at an angle θ. The values of r_{max} and θ_{min} are derived from the condition $\varepsilon_{cm} = m_e$, where ε_{cm} is the photon energy in the center-of-mass frame. The usual calculations[35] then yield

$$\nu \approx \frac{3}{4\pi} \frac{Q_\gamma \sigma_T}{cR} \left(\frac{m_e}{E_\gamma}\right)^2 \left(\ln \frac{2^{1/2} E_\gamma}{m_e} - 1\right). \qquad (25)$$

With $E_\gamma \approx 59$ MeV, $R = 6r_g$, and the value of Q_γ specified by (23), we obtain

$$\nu \approx 1.5\cdot 10^{-4} \dot{m} (T_e/m_e)^{3/2}.$$

Similarly, it can be shown that absorption of bremsstrahlung and synchrotron gamma rays in the electron gas is also small.

It should be possible to detect sources of this type of gamma radiation at distances of up to 10 MPc. Detectable extragalactic sources might be massive black holes ($M \sim 10^6$–$10^7\ M_\odot$) embedded in gas clouds. These would be manifested by modest activity in the nuclei of normal galaxies like our own. Dead quasars, which might leave behind black holes with $M \sim 10^8$–$10^{10}\ M_\odot$ and a low accretion rate in the residual gas, comprise another possibly detectable family. According to (23), the flux expected from a source at a distance of 10 Mpc, with $M \sim 10^7\ M_\odot$ and an accretion rate $\dot{m} \sim 10^{-3}$, is $8\cdot 10^{-9}$ cm^{-2}sec^{-1}. A flux of this magnitude will be accessible to future detectors.

As for Galactic sources, the first place to look is the center of the Milky Way, which may conceivably harbor a black hole with $M \sim 10^6 M_\odot$. When $\dot{m} \sim 10^{-3}$, the expected flux could be as high as $8\cdot 10^{-4}$ cm^{-2} sec^{-1}. A one-solar-mass black hole with $\dot{m} \sim 10^{-3}$ residing in a gas cloud could be detected in gamma rays out to 1 kpc ($F_\gamma \approx 8\cdot 10^{-8}$ cm^{-2} sec^{-2}).

4. DISCUSSION

The pion-production process that we have considered in this paper takes place if protons approach a black hole while remaining in regular, stable orbits. This is the situation that will prevail in a low-density gas when the accretion rate $\dot{M}$ is also low. In turn, at low densities, the regular character of the motion of the protons becomes a stabilizing factor for the electron component of the gas, since the strict preservation of plasma electrical neutrality precludes instability. At high densities the gas motion can become turbulent, and accretion may take place quasiradially. In that event, our pion production mechanism is inapplicable. It is likewise inapplicable at accretion rates approaching the critical limit, where $\dot{m} \gtrsim 10^{-3}$. The characteristic time for proton orbits to evolve, as determined by p-e Coulomb interactions, then becomes shorter than the proton orbital period about the black hole, thereby taking accretion into the quasiradial regime. Consequently, our mechanism for generating high-energy gamma rays is inappropriate to such powerful sources as quasars and Seyfert galaxies.

We have limited our concern in the present paper to quasispherical accretion, but the calculations remain qualitatively correct for the case of disk accretion as well if the disk is thick and its density is low enough. In that case, protons near the inner edge of the disk can travel somewhat out of the equatorial plane, and p-p collisions can take place under geometrical conditions resembling those treated above.

We wish to stress once more that the appearance of a shock wave has little effect on the creation of pions in a collisionless (for protons) plasma. The proton flux crossing the shock front becomes isotropic, but because the motion is indeed collisionless, most of the protons are recaptured into stable orbits. One can easily see that in a thin disk, the necessary conditions for *p-p* collisions are not present.

We predict that it will be possible to detect both Galactic and extragalactic high-energy gamma-ray sources. The center of the Milky Way is presently the most likely candidate. If the observed luminosity $L \sim 10^{40}$ erg/sec is due to the accretion of gas onto a supermassive black hole with $M \sim 10^{6} M_{\odot}$, then the accretion rate must be of order $\dot{M} = L/\alpha_0 c^2 \approx 1 \cdot 10^{20}$ g/sec and $\dot{m} = \dot{M}/\dot{M}_{av} = 1 \cdot 10^{-4}$, where $\alpha_0 = 0.1$ is the mass-to-energy conversion efficiency. We see then from (23) that the expected flux is $F_\gamma \sim 10^{-4}$ cm^{-2} sec^{-1}, a level that should already have been observed if the putative black hole at the center of the Galaxy were embedded in a gas cloud. It is conceivable, however, that flareups in the activity at the center are the result of occasional (perhaps once every 10^4 years) events consisting of the tidal disruption of a passing star. Matter from the disrupted star would form an accretion disk that could not spawn pions. In that event, gamma radiation would be associated with accretion of the surroundings gas; the accretion rate would be lower than in the disk.

Within the Milky Way, there may also be a large number of stellar-mass black holes. If such a black hole were to wind up in a gas cloud, then for $M \sim 1 M_{\odot}$, $\dot{m} \sim 10^{-3}$, and a distance $r \sim 1$ kpc, the expected flux would be $F_\gamma \approx 8 \cdot 10^{-8}$ cm^{-2} sec^{-1}.

A flux of $F_\gamma \approx 8 \cdot 10^{-9}$ cm^{-2} sec^{-1} might arrive from massive black holes with $M \sim 10^7 M_{\odot}$, $\dot{m} \sim 10^{-3}$ at distances of 10 Mpc. A search for galaxies harboring such black holes might profitably be directed at the Virgo cluster.

[1)]Hereafter we write elementary-particle masses and temperatures in energy units.

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