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TURBULENT ACCRETION ONTO MASSIVE BLACK HOLES

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ABSTRACT

Turbulent spherical accretion onto massive black holes ($M \sim 10^6$–$10^9\ M_\odot$) is considered. Temperatures of the order of 10^9–10^{10} K are found in the inner zone where the synchrotron frequencies are 10^{13}–10^{14} Hz. If the optical depth of the hole atmosphere is $\tau \sim 1$, Comptonization can produce a power-law spectrum of index ~ 1 from the synchrotron frequencies up to medium X-rays. If $\tau \gtrsim 1$, convection of radiation into the hole becomes important at small radii, reducing the efficiency of emission. The results are briefly compared with the observations.

Subject headings: black holes — galaxies: nuclei — galaxies: Seyfert — stars: accretion

I. INTRODUCTION

Accretion onto massive black holes has been proposed by various authors as a possible model of active galactic nuclei (for recent reviews, see Rees 1978; Ginzburg and Ozernoy 1977).

Observations of the X-ray emission of type 1 Seyfert galaxies (Elvis *et al.* 1978; Tananbaum *et al.* 1978) have shown that the nuclei of some of these objects have luminosities above 2 keV which are comparable or larger than in the optical band. In particular the energetics of NGC 4151, for which the spectrum is known up to 1 MeV, appears dominated by the emission at the highest frequencies (Di Cocco *et al.* 1977).

For these objects the standard picture of accretion disks is inadequate, since for the large central masses required in galactic nuclei (10^6–$10^{10}\ M_\odot$) the typical temperatures are in the UV band (Pringle and Rees 1972; Shakura and Sunyaev 1973). This leads one to introduce modifications, like the existence of a hot corona (Payne and Eardley 1978; Tsuruta 1977) or an ordered magnetic field (Blandford 1976, Lovelace 1976). Another possible approach is to consider spherical models of accretion (Fabian *et al.* 1976, Mészáros and Silk 1977).

In this paper we examine a spherical model involving turbulence and magnetic dissipation, which is an extension of the cases examined by Shvartzman 1971; Shapiro 1973*a*, *b*; Mészáros 1975; Maraschi and Treves 1977.

We have determined the temperature profile of the accreting gas for a rather wide range of hole masses (from 10^6 to $10^9\ M_\odot$) and find that the inner regions attain values within 10^9 and 10^{10} K. The energetically dominant emission mechanism from these regions is not bremsstrahlung, but rather cyclotron-synchrotron radiation from the thermal plasma. The typical frequencies of this radiation are between 10^{13} and 10^{14} Hz. However, when the optical depth to Thomson scattering becomes of order unity or larger, the Comptonization process is important, and the spectrum of the emergent radiation extends to $h\nu \approx 3kT$. We therefore attempt to explain the continuum emission of the Seyfert nuclei from the optical to the medium-energy X-rays with this process. For the very hard X-rays of NGC 4151 we suggest, as an important contributor, synchrotron emission from electrons acceleration to ultrarelativistic energies in the turbulent flow. Comptonization in a hot gas accreting onto a massive black hole has been recently discussed by Lightman, Giacconi, and Tananbaum (1978) in relation with X-ray flares from NGC 4151.

II. BASIC FORMULAE

The luminosity corresponding to an accretion rate $\dot{M}$ is

$$L = \epsilon \dot{M} GM/\bar{r} = 1.5 \times 10^{45}\ \dot{M}_{25}\ \text{ergs s}^{-1}\,, \qquad (1)$$

where $\bar{r}$ is an effective radius, which is taken to be

$$\bar{r} = 6GM/c^2 = 9 \times 10^{13}\ M_8\ \text{cm}\,, \qquad (2)$$

where M_8 is the mass of the hole in units of $10^8\ M_\odot$, and ϵ is the efficiency of conversion of the gravitational power into radiation.

Assuming free fall, the density is given by

$$n = 3.5 \times 10^9\ x^{-3/2} \dot{M}_{25} M_8^{-2}\ \text{cm}^{-3}\,, \qquad (3)$$

where x is the radial distance in units of $\bar{r}$.

The equipartition magnetic field is (Mészáros 1975)

$$B = 2.5 \times 10^3\ \dot{M}_{25}^{1/2} M_8^{-1}\ x^{-5/4}\ \text{gauss}\,, \qquad (4)$$

and the optical depth for Thomson scattering is

$$\tau_T = 0.42 \dot{M}_{25} M_8^{-1} x^{-1/2}\,.$$

The opacity for photon-photon pair production was recently considered by Cavallo and Rees (1978) and Lightman, Giacconi, and Tananbaum (1978), whence we take the approximate formula

$$\tau_{e^+e^-} \approx \frac{\sigma_T L(E_t)}{4m_e c^3}\,\bar{r}^{-1} = \frac{2.4 L_{44}(E_t)}{M_8}\,, \tag{5}$$

where $L(E_t)$ is the luminosity at the threshold energy $E_t = m_e c^2$.

For all reasonable choices of M and $\dot{M}$ appropriate for a galactic nucleus, the cooling time becomes less than the free-fall time and therefore the conversion of gravitational energy into radiation is efficient.

The dissipative heating is given by

$$\Gamma = 16.4 \dot{M}_{25} M_8^{-3} x^{-4} \text{ ergs cm}^{-3}\text{ s}^{-1}\,, \tag{6}$$

where it is supposed that 25% of the gravitational power is transformed in the heating of the gas, i.e., $\epsilon \approx 0.25$ (see Mészáros 1975).

Assuming a typical velocity between the turbulence cells of the order of the free-fall velocity, the induced electric field is (Maraschi and Treves 1977)

$$E \sim \frac{1}{c}(2GM/r)^{1/2} B \approx 4.3 \times 10^5 \dot{M}_{25}^{1/2} M_8^{-1} x^{-1/4}\ V/\text{cm}\,. \tag{7}$$

If the dominant energy loss for runaway electrons is synchrotron radiation, their emission is peaked at

$$h\nu = 0.1(3x)^{-1/2} \frac{2\pi mc^2}{\alpha} \approx 25 x^{-1/2} \text{ MeV}\,, \tag{8}$$

where α is the fine-structure constant.

III. THE TEMPERATURE PROFILES

The temperature T of the infalling gas as a function of x is obtained from the equation

$$\frac{dT}{dx} + \frac{T}{x} = \frac{4\pi m_p}{3\dot{M}}\,\bar{r}^3 x^2 (\Lambda - \Gamma)\,, \tag{9}$$

where Λ represents the energy loss rate per unit volume. Λ includes free-free and line emission, cyclo-synchrotron emission, and cooling due to electron scattering of the radiation. At large radii, line emission dominates over Γ, and a quasi-isothermal zone is formed ($T = 10^4$–10^5 K). From the condition

$$\Gamma = \Lambda = n^2 \tilde{\Lambda}(T) \tag{10}$$

the inner radius of this zone is estimated to be

$$x_{1s} \approx 1.4 \times 10^{-18} \dot{M}_{25}^{-1} M_8 \tilde{\Lambda}^{-1} \tag{11}$$

where $\tilde{\Lambda}$ is the cooling factor for line emission and depends both on the temperature and on the ionization and excitation processes. In the purely collisional case, we have adopted from Cox and Tucker (1969) $\tilde{\Lambda} = 10^{-22}$ erg cm^3 s^{-1} for $T = 10^4$ K, and $\tilde{\Lambda} = 10^{-21}$ for $T = 10^5$ K. We have integrated numerically equation (9) starting from an x_{1s} calculated assuming $\tilde{\Lambda}$ (10^5 K) $= 10^{-21}$ erg cm^3 s^{-1}.

Within x_{1s} the most important cooling mechanisms are cyclosynchrotron and Compton scattering. One has to take into account that the plasma is opaque to the synchrotron photons, and therefore the quantity of interest, as shown by Mészáros (1975), is the frequency ν^+ at which the plasma becomes transparent. We find that the extrapolation of the formulae for the cyclotron absorption coefficient given by Trubnikov (1958) and those given by Pacholczyk (1970) for the intermediate energy range $kT \approx m_e c^2$ lead to discrepancies in the estimate of ν^+. In the cases examined, cyclotron cooling is never important compared to the dissipative heating, while synchrotron cooling becomes important above 3×10^9 K, and we have therefore used the synchrotron formula given by Pacholczyck above that temperature.

The cooling rate due to the single scattering of thermal synchrotron photons by the electrons is given by

$$\Lambda_{es} = \frac{32}{3}\frac{\pi}{c^2}\frac{(kT)^2}{m_e c^2}\, n_e \sigma_T \nu^{+3}\,. \tag{12}$$

We have neglected in first approximation the effects of multiple Compton scattering, which can be important for $\tau_T \gtrsim 1$. This approximation will be briefly discussed in § IV*b*.

We have obtained the temperature profiles for various values of the hole mass ($10^6 < M < 10^9\ M_\odot$) and of the accretion rate. Some examples are shown in Figure 1. All the profiles are characterized by a rather steep rise starting at x_{1s} up to a maximum value between 10^9 and 10^{10} K, after which the temperature declines very slowly as x decreases.

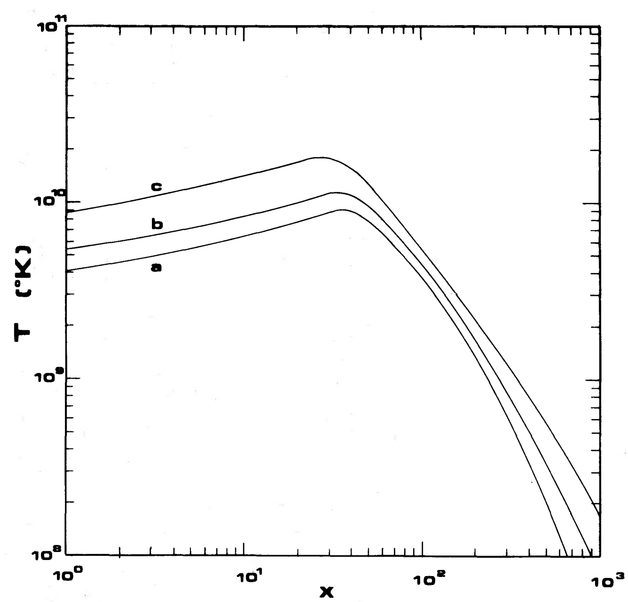

FIG. 1.—Temperature profiles for different masses and accretion rates. (*a*) $M = 10^6\ M_\odot$, $\dot{M} = 4.1 \times 10^{23}$ g s^{-1}. (*b*) $M = 10^7\ M_\odot$, $\dot{M} = 2.9 \times 10^{24}$ g s^{-1}. (*c*) $M = 10^{10}\ M_\odot$, $\dot{M} = 1.39 \times 10^{26}$ g s^{-1}.

Schematically, the temperature behavior could be approximated by two quasi-isothermal zones, one outside x_{is} at 10^4–10^5 K, the other well inside x_{is} at 10^9–10^{10} K, separated by a transition zone.

IV. SPECTRUM OF THE OUTGOING RADIATION

In this section we shall concentrate first on the spectrum of the radiation emitted by the thermal gas from the inner zone, $x < x_{is}$. The shape of the spectrum depends critically on the optical depth to Thomson scattering; therefore we shall consider separately the case $\tau \ll 1$ and $\tau \gtrsim 1$. Later we shall briefly discuss the nonthermal contribution.

a) $\tau \ll 1$

In this case the energetically dominant emission mechanism is synchrotron radiation from the 10^9–10^{10} K electrons, which yields a flat spectrum between 10^{13} and 10^{14} Hz. The contribution of Compton scattering is minor and does not exceed a frequency of 10^{14} Hz. The bremsstrahlung photons reach frequencies of the order of 10^{21} Hz, but the energy they carry is only a small fraction of the total emission power.

b) $\tau \gtrsim 1$

When the optical depth to Thomson scattering is of order 1 or greater, two effects come into play. The first is that convection of radiation into the hole becomes important and decreases the efficiency of energy extraction. The second is that Comptonization of the synchrotron photons can induce a drastic change in the form of the spectrum.

For consistency, let us first verify that a regime with $\tau > 1$ can be reached for luminosities less than the Eddington limit:

$$L_E = \frac{4\pi G M m_p c}{\sigma_T} = 1.3 \times 10^{46} M_8 \text{ ergs s}^{-1} .$$

From equation (1) the accretion rate must satisfy the condition

$$\dot{M} < \dot{M}_c \equiv \frac{4\pi m_p c \bar{r}}{\epsilon \sigma_T} = 8.7 \times 10^{25} M_8 \epsilon^{-1} \text{ g s}^{-1} . \tag{13}$$

Since $\tau(\bar{r}) = 2\sigma_T n(\bar{r})\bar{r}$, one has

$$\tau_T(\bar{r}) < \frac{2\bar{r}^{1/2}}{\epsilon} \frac{c}{(2GM)^{1/2}} ; \tag{14}$$

that is, for $\epsilon = 0.25$ and $\bar{r} = 3r_s = 6GM/c^2$, $\tau_T(\bar{r}) \leq 14$.

In order to estimate the importance of the convection process, we compare the free-fall time, $t_{ff} \sim r/v$, with the diffusion time:

$$t_d \approx \frac{n\sigma_T r^2}{c} = \frac{\sigma_T}{c} \frac{\dot{M}}{4\pi v m_p} , \tag{15}$$

where v is the infall velocity. The two time scales are equal at a radius r^+ given by

$$r^+ = \frac{\dot{M}}{\dot{M}_c} \frac{\bar{r}}{\epsilon} , \tag{16}$$

which is obviously greater than $\bar{r}$ when $\dot{M} > \epsilon \dot{M}_c$. Therefore, assuming that only the regions with $r > r^+$ contribute to the outgoing luminosity, one obtains

$$L_{out} = \epsilon \dot{M} G M / r^+ = \epsilon^2 \dot{M}_c (GM/\bar{r}) = \epsilon L_E . \tag{17}$$

Note that in this regime L_{out} is independent of the accretion rate. This result can be compared with that of Maraschi, Reina, and Treves (1976), who found that, in a disk model of accretion, L_{out} is independent of $\dot{M}$ for large values of $\dot{M}$. The computations by Kafka and Mészáros (1976) for spherical accretion, without heating, also show a similar effect, in the sense that in their model a radius appears, within which photons are transported by the gas into the hole.

The maximum optical depth for the outgoing radiation in the two regimes ($\bar{r} \lesseqgtr r^+$) depends on the accretion rate as follows:

$$\tau_T(\bar{r}) = 3.6 \frac{\dot{M}}{\epsilon \dot{M}_c} \quad \text{for } \dot{M} < \epsilon \dot{M}_c , \tag{18a}$$

$$\tau_T(r^+) = 3.6 \left(\frac{\dot{M}}{\epsilon \dot{M}_c}\right)^{1/2} \quad \text{for } \dot{M} > \epsilon \dot{M}_c . \tag{18b}$$

Let us now consider the effect of Comptonization on the spectral shape of the outgoing radiation. For a detailed discussion of this process, see Illarionov and Sunyaev (1972), Felten and Rees (1971), Shapiro, Lightman, and Eardley (1976), and Katz (1976). In the last paper, the Kompaneets equation, with the addition of a diffusive term, is studied for a monochromatic photon source at frequency $\nu = \nu_0$ surrounded by a static, homogeneous plasma sphere. It is shown that the resulting spectrum depends on the ratio of the radius R of the sphere to a critical radius

$$R_c \equiv \lambda \frac{\pi}{2} \left(\frac{m_e c^2}{kT}\right)^{1/2}$$

where $\lambda = (\sigma_T n)^{-1}$. For $R/R_c \gg 1$ the resulting spectrum approaches a Wien distribution, while for $R/R_c \ll 1$ the emission is still peaked at ν_0 but with a power-law tail of spectral index $\alpha > 1$. In particular, for $R/R_c = 0.8$, the spectrum has an index $\alpha = 1$ from $\sim \nu_0$ to $\nu = 3kT$, where T is the electron temperature. We adapt the results of Katz (1976) to our case by approximating the infalling gas with a homogeneous sphere of radius $\bar{r}$ (or r^+ if $r^+ > r^-$), density $n(\bar{r})$ or $n(r^+)$, and temperature $T = \xi\ 10^9$ K, where ξ, taken as $[T(\bar{r} \text{ or } r^+) + T_{Max}]/2 \times 10^9$, is of the order of 1–10, provided that r^+ is not much larger than the radius at which the temperature reaches $T = T_{\rm Max}$. In this approximation we have

$$R_c = 9.4 \times 10^{14} M_8{}^2 \dot{M}_{25}{}^{-1} \xi^{-1/2} \quad \text{for } \dot{M} < \epsilon \dot{M}_c\,, \quad (19a)$$

$$R_c = 9.4 \times 10^{14} M_8{}^2 \dot{M}_{25}{}^{-1} \xi^{-1/2} \left(\frac{\dot{M}}{\epsilon \dot{M}_c}\right)^{3/2} \quad \text{for } \dot{M} > \epsilon \dot{M}_c\,; \quad (19b)$$

and correspondingly

$$\bar{r}/R_c = 9.6 \times 10^{-2} M_8{}^{-1} \dot{M}_{25} \xi^{1/2} = 0.23\, \tau_T(\bar{r}) \xi^{1/2} = 0.83 \frac{\dot{M}}{\epsilon \dot{M}_c} \xi^{1/2}\,, \quad (20a)$$

$$r^+/R_c = 0.83 \left(\frac{\dot{M}}{\epsilon \dot{M}_c}\right)^{1/2} \xi^{1/2}\,. \quad (20b)$$

From this it appears that when $\dot{M}$ approaches $\dot{M}_c$, irrespective of the mass of the hole (except through ξ), Comptonization becomes the dominant effect in shaping the spectrum. Because of the approximations made, one cannot predict the exact form of the spectrum on the basis of the computations by Katz (1976) as a function of the parameters of the model, but it would seem that effective values of $R/R_c > 1$ can be attained even before $\dot{M} > \epsilon M_c$ (see Table 1). Correspondingly we would expect spectra which can be power laws from the optical to the hard X-rays, with spectral index which can be as small as 0.5 or even spectra which peak at $3kT$ with a rather minor energy flux (and also photon flux) in the optical-UV region.

However, when $\alpha < 1$, the power carried by the Comptonized photons is much larger than that emitted through the synchrotron process, and we expect that the temperature in the inner regions will be depressed compared to the values computed in the approximation of single Compton scattering. This effect still needs to be estimated in a self-consistent manner, but its consequence could be that the ratio R/R_c, which goes as $T^{1/2}$, hardly becomes greater than unity. Correspondingly a value of $\alpha \approx 1$ might represent a limit to the hardness of the outgoing spectrum.

c) *Contribution of Bremsstrahlung*

Bremsstrahlung contributes to the emission with a flat spectrum extending up to $\sim kT_e$. In order to estimate the luminosity, we approximate the atmosphere with an isothermal spherical crust of temperature $T = \xi\, 10^9$ K, of inner radius $\bar{r}$, and outer radius $x_{1s}\bar{r}$. The resulting bremsstrahlung luminosity is given by

$$L_{\rm brems} \sim 4.4 \times 10^{39} \dot{M}_{25}{}^2 M_8{}^{-1} \xi^{1/2} x_{1s} \text{ ergs s}^{-1}\,. \quad (21)$$

For the cases examined in Table 1, $L_{\rm brems}$ is between 10^{-3} and 10^{-2} of the total luminosity $L_{\rm tot}$. Therefore, if the spectral index of the Comptonized spectrum is $\alpha \sim 1$, the contribution of bremsstrahlung is negligible also at the highest frequencies.

d) *Nonthermal Radiation*

In the highly turbulent atmosphere of the accreting hole one can expect that the energy distribution function of electrons has a high-energy tail. This could be produced, for instance, by the mechanism studied by Maraschi and Treves (1977) and mentioned in § II. In this case synchrotron radiation yields photons of energy up to $\sim$25 MeV with a spectrum peaked at the highest frequency. The total luminosity due to this mechanism depends on the importance of the nonthermal tail relative to the Maxwellian distribution, and on the opacity of the atmosphere. In particular photon-photon pair production may be of importance. In fact, a photon of 25 MeV can produce a pair by scattering off another photon, if the energy of the latter is above $\sim$40 keV. Then equation (5) takes the form

$$\tau_{e^+e^-} \approx 3 \times 10^3 \frac{L(40\text{–}100 \text{ keV})}{L_{\rm E}}\,. \quad (22)$$

Therefore, if the luminosity at $\sim$100 keV is 3×10^{-4} or more of the Eddington luminosity, $\tau_{e^+e^-}$ becomes larger than 1 and the spectrum should be cut off at $\sim$1 MeV. This happens when $\dot{M} \gtrsim 5 \times 10^{-2}\, \dot{M}_c$, since in this case $\tau_{\rm T} \gtrsim 1$ and the 100 keV luminosity is comparable with the total luminosity. On the other hand, when $\dot{M} \ll 5 \times 10^{-2}\, \dot{M}_c$, the atmosphere is transparent for pair production and photons above 1 MeV are freely emitted.

TABLE 1

PARAMETERS CHARACTERIZING THE ATMOSPHERE OF AN ACCRETING MASSIVE BLACK HOLE

M_8	$\dot{M}_{25}$	$L_{\rm tot} = \epsilon GM\dot{M}/\bar{r}$ (ergs s^{-1})	$\xi \equiv [T(\bar{r}) + T_{\rm max}]/2 \times 10^9$ K	$\nu^+(r = \bar{r})$ (Hz)	$\bar{r}/R_c$	α
0.01	1.7×10^{-2}	2.4×10^{42}	8	1.5×10^{14}	0.44	2.7
	4.1×10^{-2}	6×10^{42}	6.7	1.4×10^{14}	1	0.8
0.1	1.4×10^{-1}	2×10^{43}	9.9	5.7×10^{13}	0.41	2.8
	2.9×10^{-1}	4.2×10^{43}	8.2	8.3×10^{13}	0.70	1.15
1	1.4	2×10^{44}	8.7	4.4×10^{13}	0.4	2.8
	8.2	1.2×10^{45}	6.1	5.2×10^{13}	1.9	0.12
10	13.9	2×10^{45}	13.2	1.16×10^{14}	0.47	2.4
	41.7	6×10^{45}	9.4	1.8×10^{14}	1.2	0.6

V. DISCUSSION

In this section we shall briefly compare the predictions of this model with the observations. We have shown that the temperatures attained by the infalling gas are of the order of 10^9–10^{10} K for a wide range of masses of the hole and of accretion rates. However, the dominant emission process is not bremsstrahlung, and the photons emitted in the synchrotron process are confined to the infrared portion of the spectrum. Only a scattering process can lead to a radically different frequency distribution in the emitted photons, and the efficiency of this process depends very critically on the accretion rate. When this rate increases, the scattering process yields power-law spectra extending up to $3kT$ and eventually, when it approaches the critical value corresponding to the Eddington luminosity, most of the emitted energy is concentrated in hard X-rays.

If the model applies to galactic nuclei, one would therefore expect a great variety in the frequency distribution of the central source. It is tempting to consider the "normal" galaxies with a nuclear IR excess and the Seyfert galaxies with a strong X-ray emission as two rather extreme manifestitations of the same process.

With regard to "normal" galaxies, Rieke and Lebofsky (1978) have measured the 10 μm flux from a complete sample and find that the IR excess leads to very high values of the ratio L_{10}/M, where L_{10} is the 10 μm luminosity of their nuclei, and M is the nuclear mass as deduced from the velocity dispersion in their star content. The authors propose as an explanation a contribution from a recent star burst, but admit that conditions quite different from those found in galactic H II regions must be assumed to explain the relative weakness of the radio free-free emission from these nuclei. In our model the excess IR luminosity could be explained by accretion on a central hole, when the accretion rate is small enough for the accreting gas to be essentially transparent to Thomson scattering. The IR source need not appear pointlike, if these nuclei contain enough dust to scatter the radiation and mask the central object.

At the other extreme, Seyfert nuclei are known to contain an "excess" of UV ionizing radiation, and in the last 2 years evidence has accumulated that they are also very powerful X-ray emitters. From Table 1 in Tananbaum *et al.* (1978) for a sample of Seyfert galaxies detected in X-rays, the ratio L_X/L_U (nuclear) is of the order of unity, corresponding to a spectral index between the two "frequencies" of about 1. In the present model, following the results of Katz (1976), it appears that when Comptonization is efficient enough to produce a "hard" optical spectrum up to the UV, then the spectrum extends up to $h\nu \approx 3kT$ with a constant spectral index $\alpha \approx 0.8$. Therefore, in this model the UV and the X-ray photons would have the same origin, and the correlation between UV and X-ray powers would find a straightforward explanation. On the other hand, it is important to realize that small changes in the ratio R/R_c imply large variations in the spectral shape: while with $R/R_c = 0.8$ the spectrum is a power law with $\alpha \approx 1$, with $R/R_c = 3$ most of the emergent flux is concentrated in the Wien peak at $3kT$, with a relatively minor contribution in the optical and UV. Unless the feedback on the temperature of the multiple scattering is such as to prevent R/R_c from becoming larger than unity (and this needs to be carefully checked), we would then expect from this model a rather wide range in the ratio L_X/L_U, and in particular values much greater than 1. The fact that such a spread in L_X/L_U has not been observed so far, may represent a problem, but it is also conceivable that the available data are affected by selection effects. As already pointed out by Katz (1976), one might find in objects with $L_X/L_U > 1$ the counterparts of some of the so far unidentified high-galactic-latitude X-ray sources, and these should be looked for especially among the sources which show time variability.

The X-ray power of Seyfert galaxies ranges from 10^{42} to 10^{44} ergs s^{-1}. Correspondingly the required $\dot{M}_{25}$ ranges from 10^{-2} and 1 and, in order for $\bar{r}/R_c$ to be about unity, the hole mass must range from 2×10^5 to 2×10^7 $M_\odot$. An accretion rate of 10^{25} g s^{-1} corresponds to 0.1 $M_\odot$ yr^{-1}. This is a rather extreme value if the accreting gas is fed by ordinary mass loss from stars in the nuclear bulge of a galaxy, but it is not inconceivable that the ordinary loss may be enhanced, at least not too far away from the central source, because of the atmospheric heating caused by the strong field of ionizing radiation.

Recent measurements of one of the Seyferts, NGC 4151, beyond the medium-energy X-rays, have led to the surprising result that its spectrum flattens above $E \approx 20$ keV and continues with $\alpha \approx 0$ up to 100 keV (Auriemma *et al.* 1978 and references therein). There is also an indication that this behavior may persist up to 1–10 MeV (Di Cocco *et al.* 1977), while recent observations by Shönfelder (private communication) indicate a cutoff at $\sim$3 MeV. It must be noted that if such a spectral distribution were a common feature of the Seyferts, their contribution to the background would exceed the measured value beyond $\sim$100 keV (Grindlay 1978; Bignami, Lichti, and Paul 1978); therefore, the case of NGC 4151 must be rather peculiar. In the present model one can attribute the high-energy X-rays to the nonthermal process proposed by Maraschi and Treves (1977) and relate the cutoff observed at 3 MeV to the opacity for pair production. The major problem, however, is in the energetics, because we should require that perhaps as much as 90% of the available energy goes into the relativistic particle channel. We do not know at the moment of a physical argument strongly prohibiting the above mentioned imbalance. The remarks made above on the "peculiarity" of the spectrum of NGC 4151 seem to imply that some special conditions are required, or a critical regime is required, for this situation to prevail.

The present model is based on the hypothesis of homogeneity of the flow onto the massive black hole, and this hypothesis has enabled one to distinguish two regimes of accretion, the former corresponding to a quasi-transparent atmosphere, the latter to the opaque case, which can be dominated by radiation convection. It is clear that our approximation cannot account for the variability of the emission observed in NGC 4151 (Tananbaum *et al.* 1978). This requires an inhomogeneity on a large scale, since the observed flares have powers comparable to the average power.

The problem of line emission has been ignored in the above discussion. It is well known that clumpiness of the gas is required in order to explain the observations. Moreover, our computation of the gas temperature is self-consistent only at small x since UV and X-ray radiation produced in the core will have an important effect in ionizing and heating the outer envelope.

The relaxation of the homogeneity hypothesis could also lead to a wider complexity of the spectrum of the Comptonized radiation, allowing direct explanation of the flattening of the spectrum of NGC 4151.

These last points indicate that inclusion of inhomogeneities in the flow is a necessary step for an improvement of the model.

REFERENCES

Auriemma, G., *et a* . 1978, *Ap. J.* (*Letters*), **221**, L7.
Bignami, G. F., Lichti, G. G., and Paul, J. A. 1978, *Astr. Ap.*, **68**, L15.
Blandford, R. D. 1976, *M.N.R.A.S.*, **176**, 465.
Cavallo, G., and Rees, M. J. 1978, *M.N.R.A.S.*, **183**, 359.
Cox, D. P., and Tucker, W. H. 1969, *Ap. J.*, **157**, 1157.
Di Cocco, G., *et al.* 1977, *Nature*, **270**, 31.
Elvis, H., Maccacaro, T., Wilson, A. S., Ward, M. J., Penston, M. V., Fosbury, R. A., and Perola, G. C. 1978, *M.N.R.A.S.*, **183**, 129.
Fabian, A. C., Maccagni, D., Rees, M. J., and Stoeger, W. R. 1976, *Nature*, **260**, 683.
Felten, J., and Rees, M. J. 1971, *Astr. Ap.*, **17**, 266.
Ginzburg, V. L., and Ozernoy, L. M. 1977, *Ap. Space Sci.*, **50**, 23.
Grindlay, J. 1978, *Nature*, **273**, 211.
Illarionov, A. F., and Sunyaev, R. A. 1972, *Soviet Astr.—AJ*, **16**, 45.
Kafka, P., and Mészáros, P. 1976, *Gen. Rel. Grav.*, **7**, 841.
Katz, J. I. 1976, *Ap. J.*, **206**, 910.
Lightman, A. P., Giacconi, R., and Tananbaum, H. 1978, *Ap. J.*, **224**, 375.
Lovelace, R. V. 1976, *Nature*, **262**, 649.
Maraschi, L., Reina, C., and Treves, A. 1976, *Ap. J.*, **206**, 295
Maraschi, L., and Treves, A. 1977, *Ap. J.* (*Letters*), **218**, L113
Mészáros, P. 1975, *Astr. Ap.*, **44**, 59.
Mészáros, P., and Silk, J. 1977, *Astr. Ap.*, **55**, 289.
Pacholczyk, A. G. 1970, *Radio Astrophysics* (San Francisco: Freeman).
Payne, D., and Eardley, D. M. 1977, *Ap. Letters*, **19**, 39.
Pringle, J., and Rees, M. J. 1972, *Astr. Ap.*, **21**, 1.
Rees, M. J. 1978, *Phys. Scripta*, **17**, 193.
Rieke, G. H., and Lebofsky, M. J. 1978, *Ap. J.* (*Letters*), **220**, L37.
Schönfelder, V. 1978, *Nature*, **274**, 344.
Shakura, N. J., and Sunyaev, R. A. 1973, *Astr. Ap.*, **24**, 337.
Shapiro, S. L. 1973*a*, *Ap. J.*, **180**, 531.
———. 1973*b*, *Ap. J.*, **185**, 69.
Shapiro, S. L., Lightman, A. P., and Eardley, D. M. 1976, *Ap. J.*, **204**, 187.
Shvartzman, V. F. 1971, *Soviet Astr.—AJ*, **15**, 377.
Tananbaum, H., Peters, G., Forman, W., Giacconi, R., Jones, C., and Avni, Y. 1978, *Ap. J.*, **223**, 74.
Trubnikov, B. A. 1958, *Soviet Phys. Doklady*, **3**, 136.
Tsuruta, S. 1977, *Astr. Ap.*, **61**, 647.

Note added in proof.—After this paper was completed, we found that L. Pozdnyakov, I. Sobol', and R. A. Sunyaev (*Soviet Astr.—AJ*, **21**, 708 [1977]) have treated Comptonization in a relativistic gas, $kT = nm_e c^2$, by a Monte Carlo method. They also give an approximate analytic expression for the index of the resulting power law spectrum $\alpha = -\ln \tau / \ln (12n^2)$. For the range of temperatures derived in our calculation their results are more appropriate than those of Katz. While the main conclusions remain unchanged, the range of parameters for which a power law spectral index from 2–1 is obtained is extended to lower values of τ.

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