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A WIND AND SHOCK MODEL FOR ACTIVE GALACTIC NUCLEI

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ABSTRACT

The properties of active galactic nuclei (AGNs) are discussed in terms of a new supercritical wind model. A supercritical wind from a central source is decelerated in a shock at a distance of about 1 light-year from the center. This wind-shock model explains satisfactorily the overall spectrum of AGNs in the frequency range 10^8 Hz $\leq \nu \leq 10^{25}$ Hz, the specific order of the fluctuation time scales as well as the formation of the broad-line–emitting clouds.

Subject headings: galaxies: nuclei — plasmas — quasars — radiation mechanisms — X-rays: sources

I. INTRODUCTION

X-ray observations of active galactic nuclei (AGNs) clearly show that the compact regions of radio galaxies, Seyfert galaxies, quasars, and BL Lac objects have soft X-ray luminosities that are a substantial fraction of their total luminosities (Kriss, Canizares, and Ricker 1980; Ku, Helfand, and Lucy 1980; Zamorani *et al.* 1981). Seyfert galaxies have a mean energy spectral index of 0.7 (Mushotzky *et al.* 1980). The soft X-ray flux of quasars is compatible with a power law of index 0.4 (Zamorani *et al.*), a value also observed for those sources where hard X-ray observations are available (e.g., 3C 273 and Cen A). The extrapolation of this soft X-ray flux then suggests that active galactic nuclei may have their maximum power output in the hard X-ray domain of a few 100 keV. These spectral features typically occur when a soft photon source is Comptonized in a hot thermal gas. The temperature of this hot gas is in the range of 10^9 K as demonstrated, e.g., by the overall X-ray spectrum of Cen A (Baity *et al.* 1981; Pietsch *et al.* 1981).

The observed properties of AGNs can be explained in terms of a wind and shock model described below. This picture accounts for the generation of the global spectrum of AGNs for frequencies $10^8 \leq \nu \leq 10^{25}$ Hz and for the emission of the broad optical lines. The model predicts specific relations between the minimal fluctuation times observed in the various spectral domains. The nature of the central energy source does not have to be specified; it can be imagined as any object dominated by radiation pressure (a supermassive star exceeding the Eddington limit in its nuclear luminosity or a large accretion disk around a supermassive compact object). We only assume that the energy released by this central source is carried away by a supercritical wind, in the sense that the kinetic energy flux in the wind, L_k, is expected to be comparable to the Eddington luminosity,

$$L_k \approx L_{\rm ED} = \left(1.2 \times 10^{45}\ \text{ergs s}^{-1}\right) M_7. \tag{1}$$

The properties of the wind are then given by general energetic requirements. We find that these winds are optically thick in the central region and therefore that the central source is hidden behind a photosphere and is not directly accessible to observations. The energy stored in the wind is released as UV radiation from its photosphere and, as any other type of radiation, from the interaction zone between the wind and the ambient medium.

II. ENERGY BUDGETS IN THE WIND-SHOCK MODEL

The total luminosity, $L_{\rm AGN}$, is given by the sum of the photosphere luminosity, $L_{\rm UV}$ and the kinetic energy, L_k, carried by the wind (the thermal energy in the wind is completely negligible). In the following, we only consider spherically symmetric winds generated some 100 Schwarzschild radii from the center. The wind velocity is therefore $v_w = \beta c \approx v_{\rm escape} \lesssim 0.1c$. This allows us to parametrize the mass loss rate as

$$\dot{M} = 200\, \dot{M}_{\rm ED}\left(L_k/L_{\rm ED}\right)\beta_{-1}^{-2} = \left(3\, M_\odot\ \text{yr}^{-1}\right) L_{45}\beta_{-1}^{-2}, \tag{2}$$

where the Eddington mass loss rate is given in the usual

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way: $\dot{M}_{\rm ED} = 4\pi GM/c\,\kappa_{\rm es} = (1.4 \times 10^{24}\ {\rm g\ s^{-1}})M_7$. Supercritical winds will therefore typically flow at a rate of a few $M_\odot$ per year if the central object has a mass of $\sim 10^7\ M_\odot$. This parametrization can be used to calculate the electron-scattering depth in the wind. In this way, we obtain the following expression for the radius, $R_{\rm sc}$, of the "last scattering surface" ($\tau_{\rm es} = 1$):

$$R_{\rm sc} = \kappa_{\rm es} L_k/(2\pi c^3\beta^3) = (2.3 \times 10^{15}\ {\rm cm})\,L_{45}\beta_{-1}^{-3}. \tag{3}$$

The photosphere is in general inside the last scattering surface, i.e., $R_{\rm sc} = \delta R_\star$ with $\delta \gtrsim 1$. This allows us to express the photosphere temperature, $T_\star$, in terms of the independent parameters of the wind model:

$$T_\star = (2.2 \times 10^4\ {\rm K})(L_{\rm UV}/L_k)^{1/4}\delta^{1/2}\beta_{-1}^{3/2}L_{45}^{-1/4}, \tag{4}$$

For $L_k \approx L_{\rm UV}$, the photosphere will have a temperature in the range of 20,000–30,000 K and is therefore a strong source of UV radiation (the emergent flux will have in general a modified blackbody spectrum, Felten and Rees 1972). It is then also clear that the thermal energy carried by the wind is small compared with its kinetic energy.

Malkan and Sargent (1982) have recently determined blackbody temperature and the UV-emitting surface for four Seyfert galaxies and the QSO 3C 273. They find that the additional component, necessary to fit the infrared-optical-UV spectra of these sources, is well described by a blackbody at 20,000–30,000 K and emitting areas of 6×10^{30}–3×10^{33} cm^2. Therefore, we identify this UV component with emission from the photosphere of the wind. If we assume that $L_{\rm UV} \approx L_{\rm ED}$, then the observed emitting areas correspond directly to the masses of the central objects with typical values of $5 \times 10^6\ M_\odot$ for the Seyfert galaxies and $6 \times 10^8\ M_\odot$ for 3C 273.

Radio observations of compact sources (Kellermann and Pauliny-Toth 1981) and synchrotron lifetime arguments indicate that the synchrotron emitting regions contain magnetic fields up to the order of a few Gauss. Since the kinetic energy density in the wind decreases as r^{-2}, the formation of a bow shock is unavoidable at a distance R_M from the central source determined by pressure balance:

$$R_M = \{\dot{M}\beta c/[4\pi(B^2/8\pi)]\}^{1/2}$$
$$= (7.7 \times 10^{17}\ {\rm cm})\,\dot{M}_{26}^{1/2}\beta_{-1}^{1/2}B^{-1}, \tag{5}$$

where B is the magnetic field in the compressed region (taken in units of gauss). The magnetic field lines are largely toroidal, since during compression only the component perpendicular to the wind motion is amplified. As a result of this initial compression period, two shock fronts, S and S', form; S is an inward facing shock in the wind and S' an outward traveling shock. These two shocked regions are separated by a contact surface (magnetopause) located at R_M. The pressure confining the AGN is thus of a magnetic rather than a thermal nature. A thermal confinement would require a column density largely incompatible with the observed values.

In the compression region $r \leq R_M$, the wind slows down and its kinetic energy heats the plasma. We show in the following that the UV photons from the photosphere of the wind will be Comptonized to X-ray energies in this compression region, whereas the millimeter-infrared-optical and γ-fluxes are generated in the high B-field region $r \geq R_M$.

III. THE PLASMA IN THE SHOCKED REGION

a) *The Compression Region*: $R_S \leq r \leq R_M$

The density of the wind plasma is known at the inner edge of the compression region,

$$\rho_w = \dot{M}/4\pi R_S^2 v_w = (8 \times 10^{-19}\ {\rm g\ cm^{-3}})\,\dot{M}_{26}R_{S17}^{-2}v_{w9}^{-1}, \tag{6}$$

and the conservation laws for a strong shock and equation (5) provide us then with the density in the compression region

$$n_S = n_w(\gamma_{\rm G}+1)/(\gamma_{\rm G}-1)$$
$$= (1.0 \times 10^5\ {\rm cm^{-3}})\,B^2(30\beta)^{-2}(R_M/R_S)^2, \tag{7}$$

where $\gamma_{\rm G} = c_p/c_v$. The corresponding pressure for a nonrelativistic gas is then given by

$$p_S = 3j_w^2/4\rho_w = (3/4)(B^2/8\pi)$$
$$= 3 \times 10^{-2}B^2\ {\rm dynes\ cm^{-2}}, \tag{8}$$

with $j_w = \rho_w v_w$ as the mass flux in the wind. Under the assumption that the shocked gas behaves as an ideal (nonrelativistic) gas, the temperature, T_S, in the shocked layer now only depends on the wind velocity:

$$T_S = p_S/2kn_S = (3/32k)\,m_{\rm H}c^2\beta^2 = (1.1 \times 10^9\ {\rm K})\,v_{w9}^2. \tag{9}$$

This relation shows that the temperature of the shocked layer depends critically only on the wind velocity, while the pressure is merely given by the magnetic field in the outer shocked region. Stellar wind velocities, $v_w \approx 1000$ km s^{-1}, would produce shocked plasmas with a temperature of ~ 10^7 K. A cluster of OB and WR stars as a central engine can therefore not satisfy the energy constraints observed for AGNs. The above expressions for the physical quantities n_S, p_S, and T_S of the compression layer are mean values; the spatial distribution of these quantities over the whole region must be obtained from a solution of the time-dependent MHD equations.

b) The Synchrotron Domain: $r \geq R_M$

Magnetic fields of a few gauss are necessary in this region to explain the optical synchrotron emission, since

$$\nu_S = 3eB\gamma^2/(4\pi m_e c) = (4.2 \times 10^{14}\,\mathrm{Hz})\,B\gamma_4^2. \tag{10}$$

The magnetic field and the γ-factors of the synchrotron electrons degrade with the distance from the magnetopause and fall to the values observed in the self-absorbed region, which possibly coincides with the outer shock boundary. The power emitted by each relativistic electron due to synchrotron and inverse Compton losses is given by (Rybicki and Lightman 1979)

$$P = (4/3)\sigma_T c\gamma^2\beta^2 u_{\mathrm{tot}} = (2.7 \times 10^{-6}\,\mathrm{ergs\,s^{-1}})\,\gamma_4^2 u_{\mathrm{tot}}, \tag{11}$$

where u_{tot} is the total energy density in the magnetic field and the photon field

$$u_{\mathrm{ph}} = (0.3\,\mathrm{ergs\,cm^{-3}})\,L_{\mathrm{UV}45}R_{17}^{-2}. \tag{12}$$

Since, for magnetic fields of a few gauss, both energy densities are comparable, we expect that inverse Compton emission is important for the region immediately outside the magnetopause. The continuous renewal of these optical synchrotron electrons (and positrons) must be seen in direct relation to the inner shocked region.

This outer shocked region is also the natural place for the generation of the small-scale jets (see, e.g., Rees, Begelman, and Blandford 1981). This domain coincides with the compact radio-emitting regions seen in many AGNs (Readhead and Pearson 1982). In strong sources with large mass fluxes, the plasma of the inner shocked region may pour out where the magnetic field is weak, i.e., at the pole sections of this "magnetic bubble." VLBI jets are observed to be generated in a region of a few light-years. This dimension coincides quite well with the location of the magnetopause as estimated in equation (5). The magnetic field plays then a dominant role in the confinement of the jet matter, and instabilities are the main driving mechanism for the production of the jets. In addition, (preferentially) one-sided small-scale jets are formed in this process. (A more detailed discussion will be given in Camenzind and Courvoisier 1983).

IV. THE OVERALL SPECTRUM OF AGNS

The high temperature and the high density in the compression layer, $r \leq R_M$, make Comptonization an efficient cooling mechanism for this plasma since the electron-scattering depth is about 1/10

$$\begin{aligned}\tau_N &= n_S\sigma_T R_M(\Delta_S/R_M)\\ &= 5.0 \times 10^{-2}\dot{M}_{26}^{1/2}(30\beta)^{-3/2}B(R_M/R_S)(\Delta_S/R_S),\end{aligned} \tag{13}$$

where we assumed that the thickness of the shocked layer, Δ_S, is an appreciable fraction of R_S itself. The parameter τ_N does not correspond, however, to the true electron-scattering depth of the layer, since electron-positron pairs are already present at a temperature of 10^9 K. The true scattering depth is

$$\tau_{\mathrm{es}} = (n_+ + n_-)\sigma_T\Delta_S \geq \tau_N. \tag{14}$$

The exact value of τ_{es} follows from a solution for the pair production rate. It is, however, important to note that the temperature of a plasma with $0.1 \lesssim (kT_e/m_e c^2) \lesssim 1$ and $\tau_N \lesssim 0.1$ remains extremely stable since an increase in the heating rate of the plasma leads to the creation of pairs rather than to an increase in the plasma temperature (Lightman 1982; Svensson 1982). Thermal Comptonization as a possible source of X-rays in quasars was proposed for the first time by Katz (1976).

In this semirelativistic regime, the energy spectral index of the Comptonized photon spectrum is given by an approximate expression (Pozdnyakov, Sobol, and Sunyaev 1979; Takahara 1980), for 0.1 keV $\lesssim E \lesssim kT_e$,

$$\begin{aligned}\alpha_x &= (-\log\tau_{\mathrm{es}} + 2/(\boldsymbol{T}_e + 3))/\log(12\boldsymbol{T}_e^2 + 25\boldsymbol{T}_e),\\ &\approx 0.5 \text{ for } \boldsymbol{T}_e \approx 0.5 \text{ and } \tau_{\mathrm{es}} \approx 1.\end{aligned} \tag{15}$$

For small electron-scattering depths, $\tau_{\mathrm{es}} \lesssim 0.1$, and temperatures $\boldsymbol{T}_e \lesssim 0.5$ ($\boldsymbol{T}_e = kT_e/m_e c^2$), the energy spectral index exceeds unity. Compact and strong sources are therefore expected to have flat spectra (like Cen A, 3C 273, and strong QSOs as well as Seyfert galaxies), while in more extended sources (or extremely weak sources) $\alpha_x \gtrsim 1$. The exact behavior of the spectrum in the region $E \gtrsim kT_e$ cannot be discussed analytically for semirelativistic temperatures.

The γ-ray tail of the Comptonized spectrum can be strongly influenced by annihilation contributions in hot sources with $\boldsymbol{T}_e \approx 1$ (Lightman and Band 1981). In addition, electron-positron pairs produced in the compression region are accelerated to relativistic energies by a strong turbulence in the shock (see, e.g., Axford 1981); power laws for the electron distributions with indices between 2 and 3 are naturally produced by this mechanism. They explain the observed spectral indices $\alpha_S \approx 1$ for the infrared-optical synchrotron emission.

The kinetic energy of the wind is thus used partly to heat the Comptonization region and partly to accelerate the synchrotron electrons (and positrons). In general, the Comptonized luminosity should exceed the synchrotron luminosity (unless the turbulence in the shocked layer is in equipartition). Both luminosities can be comparable in very strong sources, where (1) sufficient energy is available to create numerous pairs (if, e.g., a high temperature is reached) and (2) the pairs are accelerated by a high turbulence level.

Since the energy density in the UV photons from the photosphere is comparable to the energy density in the magnetic field of the synchrotron region, a large, inversely Comptonized flux occurs at the characteristic frequency

$$\nu_{\rm IC} \approx \gamma^2 \nu_{\rm UV} = 10^{23}\ {\rm Hz}\ \gamma_4^2 \nu_{\rm UV\,15}. \tag{16}$$

A broad spectrum is generated in the MeV range by this flux and the inversely Comptonized X-rays with a cutoff in the GeV region (the stochastic acceleration of electrons and positrons in the shock is limited by inverse Compton losses). This total γ-ray flux must be comparable to the total synchrotron loss in compact and strong sources. It is interesting in this respect that the maximum given by equation (16) (which is still in the Thomson limit) corresponds to the maximum observed in the COS B γ-ray flux from 3C 273 (for a detailed spectrum of 3C 273 see, e.g., Ulrich 1981). The γ-ray luminosity is indeed of the same order as the synchrotron luminosity. Since 3C 273 is one of the strongest UV sources (Malkan and Sargent 1982), somewhat weaker sources will have inversely Comptonized γ-ray fluxes below the COS B limit of $\sim 1.5 \times 10^{-10}$ ergs cm^{-2} s^{-1}, and possibly at lower energies, since the stochastic acceleration is also less efficient.

V. THE BROAD-LINE REGION

A widely accepted model for the broad-line–emitting region (BLR) is that of small, cold, dense clouds in a hot intercloud gas (for different models explaining the clouds as formed by thermal instabilities in a wind, see, e.g., Beltrametti 1981). The existence of a hot compression layer in our wind model for AGNs provides a natural source for cloud formation. Pressure balance for a thermal intercloud medium requires that $n_S T_S = n_c T_c$, where c stands for cloud. With the above values for the compression region, we obtain the density in the cloud condensations:

$$n_c^0 \approx (3/8)(B^2/8\pi)(kT_c)^{-1}$$
$$= (1.1 \times 10^{10}\ {\rm cm}^{-3}) B^2 T_{c4}^{-1}. \tag{17}$$

Here, n_c^0 is the density of the cold condensations in the cloud formation front; these clouds feel the strong central UV flux $L_{\rm UV}$ and farther away the X-rays produced in the compression region. The clouds, born with zero initial velocity, are then accelerated by the UV radiation pressure. The acceleration is only inhibited by gravitation if the central mass exceeds a critical mass, M_c, given by

$$M_c = (3f_\Gamma/4)(c^2/8\pi G)\alpha_B(h\nu_{\rm H}/kT_c)(\dot{M}/m_{\rm H})(\beta/c^2)$$
$$= (6.7 \times 10^9\ M_\odot)\dot{M}_{26} T_{c4}^{-1}\beta_{-1}, \tag{18}$$

where $f_\Gamma \approx 1$ for spectra close to the blackbody. As a consequence, the theory of radiatively accelerated clouds as developed by Blumenthal and Mathews (1979) directly applies to our situation. Clouds with an initial density of $\sim 10^{10}$ cm^{-3} are radially accelerated to velocities of $\sim$ 10,000 km s^{-1} and densities in the observed regime of 10^8–10^{10} cm^{-3} (Mathews 1982). These clouds produce emission-line profiles which are essentially logarithmic for pure UV acceleration. Shock heating and the more or less isotropic Comptonized photon field make Lyα profiles symmetric, since a gross anisotropy in the Lyα profiles would result from a one-sided irradiation of the clouds.

The level of ionization that best accounts for the equivalent widths of typical QSO broad lines corresponds to a ratio of incident ionizing photon density to electron density in the clouds, $\Gamma = n_{\rm ph}/n_c$, of around 10^{-2} (Davidson and Netzer 1979). We can now use the above shock front conditions to express Γ in terms of the characteristic wind model parameters:

$$\Gamma = (4/3)(L_{\rm UV}/L_k)(kT_c/h\nu_{\rm H})\beta$$
$$= 0.84 \times 10^{-2}(L_{\rm UV}/L_k)T_{c4}\beta_{-1}. \tag{19}$$

The wind model accounts thus very naturally for the observed properties of the QSO broad emission lines. The absence of strong emission lines in BL Lac objects could be due to higher pressures in the shocked region (given by higher magnetic fields). The cold condensations now become too dense, $n_c^0 \gtrsim 10^{11}$ cm^{-3}, and the UV acceleration mechanism breaks down (Mathews 1982). In addition, cloud evaporation stabilizes the shocked plasma.

VI. CONCLUSION

We discussed in this *Letter* a model which describes most of the observed properties of AGNs. The wind model includes a strong UV flux from the wind photosphere, which is then partly Comptonized in the interaction zone between the strong supercritical wind and the ambient medium of the galactic nuclei. Electrons and positrons accelerated to highly relativistic energies in the shock produce the IR and optical synchrotron spectrum and at the same time a high γ-ray flux by inverse Comptonization of the UV and X-ray photons.

Physical arguments strongly constrain the wind parameters. The only free parameters of the model are the UV luminosity, L_{UV}, the kinetic energy flux, L_k, and the magnetic field, B, in the nuclear region. Future work will be needed to decide what type of central objects can emit winds as described in § II. We stress, however, that the existence of a UV photosphere constrains the smallest observable dimension to some 100 Schwarzschild radii. As a consequence of the stratified structure of the wind-shock model, we find the following relations between the minimum fluctuation times, t, in the different spectral domains:

$$t(\mathrm{UV}) < t(X) < t(\mathrm{opt}), \qquad t(\mathrm{IR}), \quad \text{and} \quad t(\gamma). \tag{20}$$

We also expect fluctuations in the UV flux to be strongly correlated with the X-ray fluctuations, which must lag behind the former ones by a typical Comptonization time.

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