THE FORMATION OF BROAD-LINE CLOUDS IN THE ACCRETION SHOCKS OF ACTIVE GALACTIC NUCLEI

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ABSTRACT

Recent work on the gas dynamics in the Galactic center has improved our understanding of the accretion processes in galactic nuclei, particularly with regard to properties such as the specific angular momentum distribution, density, and temperature of the inflowing plasma. With the appropriate extrapolation of the physical conditions, this information can be valuable in trying to determine the origin of the broad-line region (BLR) in active galactic nuclei (AGNs). In this paper, we explore various scenarios for cloud formation based on the underlying principle that the source of plasma is ultimately that portion of the gas trapped by the central black hole from the interstellar medium. Based on what we know about the Galactic center, it is likely that in highly dynamic environments such as this, the supply of matter is due mostly to stellar winds from the central cluster. Winds accreting onto a central black hole are subjected to several disturbances capable of producing shocks, including a Bondi-Hoyle flow, stellar wind-wind collisions, and turbulence. Shocked gas is initially compressed and heated out of thermal equilibrium with the ambient radiation field; a cooling instability sets in as the gas is cooled via inverse-Compton and bremsstrahlung processes. If the cooling time is less than the dynamical flow time through the shock region, the gas may clump to form the clouds responsible for broad-line emission seen in many AGN spectra. Clouds produced by this process display the correct range of densities and velocity fields seen in broad emission lines. Very importantly, the cloud distribution agrees with the results of reverberation studies, in which it is seen that the central line peak (due to infalling gas at large radii) responds more slowly to continuum changes than the line wings, which originate in the faster moving, circularized clouds at smaller radii. Finally, we provide an example of fitting an observed line profile using the parameters specified by our model.

Subject headings: accretion, accretion disks — galaxies: active — galaxies: nuclei — galaxies: Seyfert — line: profiles — shock waves

1. INTRODUCTION

The spectra of many active galactic nuclei (AGNs), including Seyfert galaxies and quasars, are distinguished by strong, broad emission lines, with a full-width at maximum intensity (FWHM) of $\sim 5000$ km s$^{-1}$, and a full-width at zero intensity (FWZI) of $\sim 20,000$ km s$^{-1}$ (e.g., Peterson 1997). From the observed strength of UV emission lines, we know that the temperature of the emitting plasma is on the order of a few times $10^4$ K (e.g., Osterbrock 1989), insufficient to produce the observed line widths via thermal (Doppler) broadening. Instead, bulk motions of the broad-line region (BLR) gases appear to be responsible for the line broadening.

Because reverberation studies show a direct response of emission-line strengths to continuum variability (e.g., Clavel et al. 1991), we know that the BLR gas must be photoionized by the continuum. The International AGN Watch consortium has carried out long-term optical and ultraviolet monitoring on a set of four Seyfert 1 galaxies: NGC 5548 (e.g., Korista et al. 1995; Peterson et al. 1999), NGC 3783 (Reichert et al. 1994; Stirpe et al. 1994), Fairall 9 (Rodríguez-Pascual et al. 1997; Santos-Lleó et al. 1997), and NGC 7469 (Wanders et al. 1997; Collier et al. 1998); and the broad-line radio galaxy 3C 390.3 (Dietrich et al. 1998; O'Brien et al. 1998). In all sources, it is observed that higher ionization lines respond faster than lower ionization lines. This would indicate that the former are found at smaller radii using a simple $r \propto \tau_{\text{delay}}$ argument. The response time for the same line varies by source, even when the luminosities are very similar, indicating that the simple $r_{\text{BLR}} \propto L^{1/2}$ rule alone does not determine the size of the BLR. Indeed, other factors, such as geometry, viewing angle, and spectral energy distribution (SED) may play equally important roles in determining its volume (e.g., Robinson 1995; Wandel 1997).

What the reverberation studies do tell us, however, is that the size of the BLR ranges from a few to several hundred light days, with a radial dependence on ionization state and possibly other physical properties. We also see that the response delay within a given source tends to increase with increasing luminosity (Peterson et al. 1999), consistent with the $r_{\text{BLR}} \propto L^{1/2}$ rule. The optical continuum displays little or no lag ($\tau_{\text{delay}} \lesssim 2$ days) in variability with respect to the ultraviolet continuum, and the amplitude of variations are typically weaker at longer wavelengths (Collier et al. 1998).

The ionization parameter, $U \equiv n_e/n$, which is the ratio of the number density of hydrogen ionizing photons, $n_e$, to the number density of hydrogen nuclei, $n$, determines the physical state of the BLR plasma. In modeling the BLR with photoionization simulations, we observe that a value of

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10^{-2} < U < 1 \) (e.g., Netzer 1990) is required to reproduce the correct line strength and ionization state of the gas. Because of the radial structure of the BLR, it is likely that the density of emitting gas varies with radius, within the range \( 10^8 \lesssim n \lesssim 10^{11} \text{ cm}^{-3} \) (e.g., Peterson 1997). However, Marziani et al. (1996) find that the BLR may extend to even higher densities \( (n \sim 10^{12.5} \text{ cm}^{-3}) \). In addition, the BLR may contain a mixture of optically thick and thin gases. Optically thick gas must be present to account for the variability of the low-ionization Mg II, Lyz, and Balmer lines (e.g., Ferland et al. 1992). Optically thin gas, on the other hand, may account for the Baldwin effect, a negative correlation between the ultraviolet emission-line equivalent width and continuum luminosity; and the Wamsteker-Colina effect, a negative correlation between the C IV \( \lambda 1549/\text{Lyz \text{ ratio}} \) and continuum luminosity (Shields, Ferland & Peterson 1995). Green (1996), however, suggests that these effects are due to changes in the SED with luminosity. The set of other important observational constraints include (1) that the absence of a deep Lyz absorption edge in AGN spectra indicates that the BLR gas must cover only a small fraction (5%-25%) of the continuum source (e.g., Bottorff et al. 1997), and (2) that the observed line strength to continuum ratio requires a small volume-filling factor \( (\sim 10^{-7} \); Netzer 1990).

The recent work in modeling (and observations of) the BLR suggests that the clouds are spread over a wide range of radii and may display a wide range in particle density, \( n \), and column depth, \( N_{HI} \), at each radius. An important result of the “locally optimally emitting cloud” (LOC) models is that the predicted integrated spectrum from all clouds in such a mixed population depends only weakly on the properties of individual clouds (Baldwin et al. 1995; Baldwin 1997; Korista, Baldwin, & Ferland 1998). The observed spectrum instead depends on the global properties of the BLR, such as the SED, elemental abundances, and the spatial distribution of clouds. A consequence of the weak dependence on many of the input parameters is that several different models may be able to account for at least some of the observed spectra.

1.1. A Sample of Current Models

There are several models in the literature that account for the origin and nature of the BLR. For example, Emmering, Blandford, \& Shlosman (1992) propose that the BLR is associated with magnetohydrodynamic winds originating in a dusty molecular accretion disk. Dense molecular clouds are loaded onto magnetic lines threading the disk and are centrifugally accelerated outward. Being exposed to the central continuum, these clouds are quickly photoionized and produce the observed emission lines. This model correctly “postdicts” both the observed shape and differential response time of the C IV \( \lambda 1549 \) line, with its mid-red wing portion responding fastest to continuum variations (Bottorff et al. 1997).

In a different model, Murray et al. (1995) propose that the broad absorption lines (BALs) seen in \( \sim 10\% \) of radio-quiet QSOs are produced in outwardly flowing radiation- and gas-pressure-driven winds rising from an accretion disk. These winds would also be partially responsible for the broad-line emission. Because this model requires shielding of the absorbing gas from soft X-rays, the absence of BALs in radio-loud quasars and Seyfert galaxies is explained by the fact that these objects are strong X-ray emitters (Murray \& Chiang 1995). Cassidy \& Raine (1996) present a similar model in which BLR clouds form as the result of the interaction of an outflowing wind with the surface of an accretion disk.

Alexander \& Netzer (1994) propose that the AGN broad-line emission originates in the winds or envelopes of bloated stars in the nuclear environment. They obtain good agreement with the line ratios and response features seen in AGNs, but they encounter some difficulty in reproducing the broad-line wings (Alexander \& Netzer 1997).

Finally, Perry \& Dyson (1985) propose that the BLR clouds are formed as the result of a cooling instability that occurs when an outflowing wind from the black hole encounters an “astrophysical obstacle” and is shocked. Rapid cooling in the shocked gas causes the plasma to clump into clouds, and the observed line widths are then due to cloud acceleration along the shocks.

1.2. The Analogy with the Galactic Center

In this paper, we take the approach that it may be worthwhile in formulating a model for the BLR to seek guidance from the galactic nucleus we know best—that of our own Galaxy. The evidence for the presence of a supermassive black hole, coincident with the radio source Sgr A* at the Galactic center, is now the most compelling of any such systems (for a recent review dealing mostly with the observational characteristics of this region, see Mezger, Duschl, \& Zylka 1996; for a summary of the theoretical status concerning Sgr A*, see Melia 1998 and Melia \& Falcke 2001). The motions of stars within 1 pc of Sgr A* seem to require a central dark mass of \( (2.61 \pm 0.35) \times 10^6 \text{ M}_\odot \) (Genzel et al. 1997; Ghez et al. 1998), in good agreement with earlier ionized gas kinematics and velocity dispersion measurements.

Our proximity to the Galactic center provides us with the rather unique opportunity of examining the gas dynamics surrounding such a massive pointlike object with unprecedented detail. Combined with multidimensional hydrodynamical simulations, this extensive body of multiwavelength data is opening our view into the complex patterns of plasma-plasma and plasma-stellar interactions. It is likely that many of Sgr A*‘s characteristics are associated with the liberation of gravitational energy as gas from the ambient medium falls into a central potential well (Melia 1994; Ruffert \& Melia 1994). There is ample observational evidence in this region for the existence of rather strong winds in and around Sgr A* itself (from which the latter is accreting), e.g., the cluster of mass-losing, blue, luminous stars comprising the IRS 16 assemblage located within several arcseconds from the nucleus. Measurements of high-outflow velocities associated with IR sources in Sgr A West (Krabbe et al. 1991) and in IRS 16 (Geballe et al. 1991), the H\( _2 \) emission in the circumnuclear disk (CND) from molecular gas being shocked by a nuclear mass outflow (Genzel et al. 1996; but see Jackson et al. 1993 for the potential importance of UV photodissociation in promoting this H\( _2 \) emission), broad Br\( \gamma \), Br\( \alpha \), and He I emission lines from the vicinity of IRS 16 (Hall, Kleinmann, \& Scoville 1982; Allen, Hyland, \& Hillier 1990), and radio-continuum observations of IRS 7 (Yusef-Zadeh \& Melia 1992), provide clear evidence of a hypersonic wind, with a velocity \( v_\infty \sim 500-1000 \text{ km s}^{-1} \), a number density of \( n_\infty \sim 10^{14-16} \text{ cm}^{-3} \), and a total mass-loss rate of \( \dot{M}_\infty \sim 3-4 \times 10^{-3} \text{ M}_\odot \text{ yr}^{-1} \), pervading the inner parsec of the Galaxy.
In recent years, several studies have addressed the question of what the physical state of this gas is likely to be as it descends into the deep gravitational potential well of the massive black hole. In the classical Bondi-Hoyle (B-H) scenario (Bondi & Hoyle 1944), the mass accretion rate for a uniform hypersonic flow is \( \dot{M} = \frac{\pi R_d^2 m_n n_{e} v_{w}}{GM} \), in terms of the accretion radius \( R_d \equiv \frac{2GM}{v_w^2} \). With the conditions at the Galactic center (see above), we would therefore expect an accretion rate \( \dot{M} \sim 10^{22} \text{ g s}^{-1} \) onto the black hole, with a capture radius \( R_d \sim 0.02 \) pc.

In reality, the flow past the supermassive black hole is not likely to be uniform, so this value of \( \dot{M} \) may be greatly underestimated. For example, one might expect many shocks to form as a result of wind-wind collisions within the cluster of wind-producing stars, even before the plasma reaches \( R_c \). With this consequent loss of bulk kinetic energy, it would not be surprising to see the black hole accrete at an even larger rate than in the uniform case. The implications for the gas dynamics in the region surrounding the black hole are significant. Coker & Melia (1997) have undertaken the task of simulating the B-H accretion from the spherical winds of a distribution of 10 individual point sources located at an average distance of a few \( R_d \) from the central object. The results of these simulations show that the accretion rate depends not only on the distance of the mass-losing star cluster from the accretor but also on the relative spatial distribution of the sources.

These calculations indicate that to fully appreciate the morphology of the gaseous environment surrounding the accretor, one must pay particular attention to the spatial distribution of specific angular momentum \( l \) in the accreting gas. Written as \( l = \sigma R_d \), where \( R_d \equiv 2GM/c^2 \) is the Schwarzschild radius, the (modeled) accreted \( \sigma \) can vary by 50% over \( \pm 200 \) yr with an average equilibrium value of 10–50 for the conditions in the Galactic center. This is interesting in view of the fact that earlier simulations based on a uniform flow—the “classic” Bondi-Hoyle accretion, producing a bow shock—resulted in \( \langle \sigma \rangle \sim 3–20 \). It appears that even with a large amount of angular momentum present in the wind, relatively little specific angular momentum is actually accreted. This is understandable, since clumps of gas with a high specific angular momentum do not penetrate to within \( R_d \). The variability in the sign of the components of \( \sigma \) suggests that if an accretion disk forms at all, it dissolves and reforms (perhaps) with a different sense of spin on a timescale of \( \sim 100 \) yr (Coker & Melia 1997).

The fact that AGNs are significantly more gas rich and display a more powerful array of phenomena than the Galactic center could mean that these ideas derived from the latter may not be valid in the case of the former. However, one area in which this type of gas morphology would certainly have a significant impact is in the structure and nature of the BLR. Our intention in this paper is therefore to frame our investigation of the BLR in AGNs with the conditions (i.e., clumping, distribution in specific angular momentum \( \sigma \), density, and temperature) we now believe to be prevalent in the Galactic nucleus, although scaled accordingly.

In § 2, we present a description of the model; in § 3 we discuss general results; in § 4 we test the plausibility of forming BLR clouds in specific shocks; and in § 5 we use the density and velocity distributions predicted by our model to fit an observed broad emission line.

2. OVERVIEW OF THE MODEL

2.1. An Accretion-Shock Scenario for the Production of BLR Clouds

We suggest that many of the observed properties of the BLR can be explained by a simple picture of cloud production within the accretion shocks surrounding the central black hole. For this, we adopt several of the ideas introduced in Perry & Dyson (1985, hereafter PD85) for the formation of clouds from cooling instabilities in these regions. In their model, a hypersonic, outflowing wind is incident upon a supernova remnant or other astrophysical obstacles, causing bow shocks to form around them. The shocked gas is compressed and heated, and is brought out of thermal equilibrium with the radiation field.

The equilibrium temperature, \( T_{eq} \), of a gas whose heating/cooling is dominated by radiative processes is determined by another ionization parameter (Krolik, McKee, & Tarter 1981),

\[
\Xi \equiv \frac{F_{ion}}{nkT_c},
\]

where \( F_{ion} \) is the ionizing flux between 1 and \( 10^3 \) ryd; then, \( T_{eq} = T_{eq}(\Xi) \). Note that \( \Xi = U_{\gamma}/E_{gas} \), where \( \gamma \) is the average energy of the ionizing photons and \( E_{gas} \) is the thermal energy of the gas. Because the shock temperature \( T_s \gg T_{eq} \), the shocked gas will rapidly cool via inverse-Compton and bremsstrahlung processes. If the cooling time is shorter than the dynamical time for the gas to flow along the shock, the cooled gas will clump and form clouds, which then stream along and behind the shock.

An important parameter in determining \( T_{eq} \) is the Compton temperature, \( T_c \), at which Compton heating and cooling processes balance (e.g., Krolik et al. 1981; Guillet 1986). The Compton temperature is highly sensitive to the shape of the continuum, particularly at high (X-ray and gamma-ray) energies; typical AGN spectra have \( T_c \approx 0.01–5 \) (Mathews & Ferland 1987), where \( T_c = T_c(10^7 \text{ K}) \). For these high values of \( T_c \), we can state a couple of generalities about \( T_{eq} \): if \( \Xi \gg 1 \), Compton processes dominate, and \( T_{eq} \approx T_c \); if \( \Xi \ll 1 \), collisional (bremsstrahlung) and recombination-line cooling dominate, and \( T_{eq} \approx 1–3 \times 10^4 \text{ K} \).

In our picture, we assume that gravitation dominates over the outward radiation pressure within the BLR (either because the outward radiation field is sub-Eddington, or because the radiative emission is anisotropic), allowing a hypersonic, accreting wind to feed the central black hole. In the AGN context, we reexamine the PD85 result for stellar wind bow shocks, and extend the idea of “astrophysical obstacles” to include Bondi-Hoyle accretion shocks and density perturbations due to wind-wind collisions and turbulence in the accretion flow. As discussed above, this is motivated by the recent simulations of the highly variable gas flows at the Galactic center.

2.2. Model Parameters

Our model requires the specification of several parameters, including the density, velocity, and temperature profiles of the accreting wind as functions of radius, as well as the intensity and Compton temperature of the continuum. In order to keep our arguments general, we choose to model the wind flow using simple dimensional requirements. We adopt the view that the accreting gas circularizes before it...
reaches the event horizon, thereby forming a disk at small radii. The existence of an accretion disk in AGNs is inferred from, e.g., the asymmetry observed in many sources (Brotherton 1996; Glenn, Schmidt, & Foltz 1994). Again writing the specific angular momentum as \( l \equiv mcR_s \), it is easy to show that the circularization radius is

\[
R_{\text{circ}} = \frac{2\sigma^2 R_s}{c}. \tag{2}
\]

If the central object dominates the gravitational potential, then the characteristic wind velocities as the free-fall velocity, \( v_{\text{ff}}(r) \approx v_{\text{ff}}(2GM/r)^{1/2} \). From mass conservation (which in the outer region gives \( M = 4\pi r^2 n_w v_w \)) and assuming that the bolometric luminosity is related to the mass accretion rate by

\[
L_{\text{bol}} = \epsilon M c^2, \tag{3}
\]

where \( \epsilon < 1 \) is the accretion efficiency, the wind density can be expressed as

\[
n_w(r) = \frac{L_{\text{bol}}}{4\pi \sqrt{2GM\epsilon} c^2 r^{3/2}}. \tag{4}
\]

The temperature \( T_w \) of the gas undergoing steady, spherical infall is described by the equation (Mathews & Ferland 1987)

\[
dT_w/dr = -\frac{T_w}{r} + \frac{F_c(T_w)}{v_w}, \tag{5}
\]

which we use as an approximation for our nonsteady flow. In equation (5), the first term represents compressional heating and the second term is due to radiative heating/cooling, for which \( F_c(T) \) is the cooling function of the gas (cf. eq. [A2] in Appendix A). Since we are mainly concerned with establishing a minimum temperature of the flow, we have neglected the effects of viscosity, shocks, and the dissipation of magnetic energy, all of which may raise the value of \( T_w \).

In a recent study, Wandel, Peterson, & Malkan (1999, hereafter WPM99) used reverberation data to infer central masses of \( M_8 \approx 0.02 - 4 \) [where \( M_8 \equiv M/(10^8 M_\odot) \)] for a sample of 17 Seyfert 1 galaxies and two quasars. They also determined a mass–monochromatic luminosity relation of \( \lambda L_\lambda(5100 \text{ Å}) \approx 10^{44} M_8^{1.25} \) ergs s\(^{-1}\) for these objects. Setting \( L_{\text{bol}} = \lambda L_{\text{bol}}(5100 \text{ Å}) \) for the monochromatic–bolometric luminosity relation, with \( f_{\text{bol}} = 10 \) consistent with the findings of Bechtold et al. (1987), we obtain

\[
L_{\text{bol}} = 10^{45} M_8^{1.25} \text{ ergs s}^{-1}. \tag{6}
\]

For the ionizing luminosity, we set \( L_{\text{ion}} = f_{\text{ion}} \lambda L_{\lambda}(5100 \text{ Å}) \), where \( f_{\text{ion}} = 5 \) has been chosen as a fiducial value (N.B.: WPM99 estimate \( f_{\text{ion}} \approx 10 \)). Under these assumptions, the only free parameters in the model are \( M, \epsilon, T_C, \) and \( \sigma \).

### 2.3. Optical Depth of the Flow

The observed absence of the Fe K-shell edge in most AGN spectra indicates that the intercloud medium must be optically thin to X-radiation (e.g., Mathews & Ferland 1987). This limit is written as \( \tau_K < 1 \), where

\[
\tau_K = \int_0^\infty \delta_{\text{circ}} \sigma_K n_w(r) dr \tag{7}
\]

is the Fe K-shell optical depth, \( \sigma_K = 2.3 \times 10^{-20} \text{ cm}^2 \) is the total K-shell cross section (e.g., Morrison & McCammon 1983), and \( \delta_{\text{Fe}} \) is the elemental abundance of iron. Assuming that \( \delta_{\text{Fe}} = 3.3 \times 10^{-3} \) (corresponding to the local ISM value; Dalgarno & Layzer 1987) and that all Fe ions in the flow retain at least two electrons (a conservative estimate, given the likely high temperature of the gas), then the condition for the flow to remain optically thin to X-radiation is

\[
\tau_K \geq \frac{0.03 M_8^{0.25}}{\epsilon}. \tag{8}
\]

This follows from the use of equations (3), (4), (6), and (7), with the appropriate definition of \( r_{\text{circ}} \) in equation (2). It is clear that this condition is met for reasonable values of \( \epsilon \) and \( \sigma \) (see § 1.2). Because \( \delta_{\text{Fe}} \approx \sigma \), the flow will then also be optically thin to Compton scattering.

#### 2.4. Cloud Formation

In order for clouds to form, the cooling time, \( t_{\text{cool}} \), of the shocked gas must be less than the dynamical time, \( t_{\text{dyn}} \), for the gas to be transported through the shock region; i.e., \( t_{\text{cool}} < t_{\text{dyn}} \). The cooling time can be calculated numerically (cf. eq. [A1] of Appendix A) if the initial temperature of the shocked gas is known. Assuming that the shock converts the ordered velocity of the flow into random (thermal) motions, the initial temperature should be \( T_i \approx n_{\text{hi}} \Delta v^2 / 3k \), where \( \Delta v^2 \) is the change in the square of the velocity across the shock.

We assume that the preshock conditions are those of the wind; i.e., \( n_w \) and \( v_w \) are used as the preshock density and velocity, respectively. At a strong shock, we have

\[
v_s = v_s^0 + v_s(t), \quad n_s = n_s^0 + n_s v_s^0, \tag{9}
\]

where the superscripts \( (n) \) and \( (t) \) refer to the normal and tangential velocity components, respectively, relative to the shock front, and \( v_s \) is the velocity of the shocked gas. In bow shocks, most of the kinetic energy of the incident flow is dissipated, so \( \Delta v^2 \approx v_s^2 \). For shocks between obliquely incident gas flows, it is the component of the wind velocity normal to the shock that is converted into thermal energy, so \( \Delta v^2 \approx v_s^2 \).

#### 2.5. Physical Properties of the Cooled Gas

As the shocked gas cools, it clumps to form clouds. Their physical characteristics, such as the number density \( n_c \), the ionization parameter \( U \), and the column depth \( N_H \), can be determined from the luminous ionizing flux, the SED, and \( n_c \). Finally, the column density of a cloud is given by \( N_H \approx n_c l_c \), where \( l_c \) is the cloud size.

In the PD85 model, the maximum cloud size is set by the coherence length, \( l_{\text{coherence}} \leq l_{\text{cool}} c_s \), where \( c_s \) is the sound speed in the shocked gas. If the cooling is isobaric and steady, \( l_c = (T_v/T_C)^{1/3} l_{\text{coherence}} \). It seems that this model is overly optimistic, however, since turbulent mixing is likely to be very important in any shock. Random motions of the turbulent fluid will disrupt coherence within the cooling gas; the maximum cloud size is then dictated by the smallest
scale at which turbulence persists. Unfortunately, this scale is not specified by our simple model, so \( N_{\text{H}} \) remains relatively undetermined.

2.6. Cloud Confinement

Krolik et al. (1981) first proposed the coexistence of cool, dense clouds (the source of the broad emission lines) confined by a hot, rarefied medium. They showed that it was possible, under the right spectral conditions, to have the two phases in pressure equilibrium (i.e., to have the same value of \( \Xi \)) and yet have vastly different temperatures. Unfortunately, two stable states can only coexist in very hard AGN spectra, with \( T_C \gtrsim 10^8 \) K. Most AGN spectra are much softer than this, effectively ruling out the two-phase pressure equilibrium condition (Fabian et al. 1986). Dense clouds emerging from the high-pressure shock environment will rapidly expand into the ambient flow at their sound speed (e.g., Reynolds & Fabian 1995). Therefore, unless some other confinement mechanism is introduced, such as a magnetic field (e.g., Emmering et al. 1992), the clouds produced within a shock are likely to survive only within the shock itself. As a result, the cloud motions are dictated by the shock motions, which are in turn dictated by the wind flow. In our picture (see below), the clouds are therefore not confined, but arise (and eventually disperse) throughout the accreting medium as shocks form and dissipate.

3. RESULTS

The condition for cooling, i.e., \( t_{\text{cool}} < t_{\text{dyn}} \), sets a minimum length scale for the shock region. Shocked gas in regions smaller than this size will simply flow out of the region before it has time to cool. From the discussion in § 2.4, we note that the dynamical time can be expressed as \( t_{\text{dyn}} \approx d_s/v_s^{(0)} \), where \( d_s \) is the size of the shock region. Our requirement for the minimum shock size is therefore

\[
d_s > v_s^{(0)} t_{\text{cool}}.
\]

In Figures 1 and 2, we plot the minimum value of \( d_s \) for a range of values in the parameters \( M \), \( \epsilon \), and \( T_C \). We have here set \( v_s^{(0)} = v_w/4 \), an upper limit that occurs when the colliding winds are incident normally. Our \( d_s^{(\text{min})} \) estimates are therefore rather conservative; for obliquely incident winds, smaller shock regions may suffice. It is reasonable to assume that shocks are possible sites for cloud formation only if \( d_s^{(\text{min})} \leq r \).

Figure 1 illustrates the effect of varying the central mass \( M \) (left) and accretion efficiency \( \epsilon \) (right). Increasing \( M \) has the effect of decreasing \( d_s^{(\text{min})} \) at any given radius, thereby extending the plausible cloud production region to larger radii. This is because both the luminosity (via eq. [6]) and wind density (via eq. [3]) increase with \( M \). With Compton cooling being proportional to the luminosity, and bremsstrahlung cooling (per particle) being proportional to the density (cf. eq. [A2]), the value of \( t_{\text{cool}} \) becomes smaller. Decreasing the value of \( \epsilon \) also decreases \( t_{\text{cool}} \); smaller values

![Figure 1](image1)

![Figure 2](image2)
shocks surrounding the stellar wind sources. In contrast to the outflow assumed by these authors. In this model, but now with an inflow (due to the accretion of ambient gas onto the central engine) to act as the agent of interaction with the winds from stars embedded within (although not comoving with) this plasma; this is in contrast to the outflow assumed by these authors. In this picture, broad-line clouds are produced within the bow shocks surrounding the stellar wind sources.

Note that cooling to equilibrium temperature, $T_{eq}$, thereby enhancing the bremsstrahlung cooling rate. The cooling times, with the cooling occurring slightly faster in the case of the softer continuum spectrum. However, it should be noted that is unlikely to remain the only relevant scale, since a Bondi-Hoyle shock is likely to break up into smaller scale shocks in a realistic (unsteady) flow. This would have the effect of reducing $t_{dyn}$. Production of BLR clouds by this mechanism is therefore dependent on the stability of the large-scale shock structure.

In this case, the size of the shock is determined by the standoff distance (Perry & Dyson 1985), which gives

$$d_{sw}^\text{(min)} \approx 3.1 \times 10^{29} \left( \frac{E_{36}}{n_w v_w v_o} \right)^{1/2} \text{[cm]},$$

where $E_{36}$ is the kinetic energy outflow rate in the stellar wind in units of $10^{36}$ ergs s$^{-1}$, and $v_o$ is the outflow velocity.

In Figure 4, we plot $d_{sw}^\text{(min)}$ using the fiducial values $v_o \approx 2000$ km s$^{-1}$ and $E_{36} \approx 100$, typical for Wolf-Rayet stars (although these are probably upper limits for a typical stellar population). It can be seen that the size of these shocks is probably too small for these to be viable sites for cloud production via radiative cooling, thus confirming the PD85 result.

4. MODELS OF SHOCK-FORMED CLOUDS

4.1. Stellar Wind Bow-Shock Model

We next study a sample of shock-producing mechanisms with the goal of determining plausible shock sites for cloud production. Let us begin by first considering the PD85 model, but now with an inflow (due to the accretion of ambient gas onto the central engine) to act as the agent of interaction with the winds from stars embedded within (although not comoving with) this plasma; this is in contrast to the outflow assumed by these authors. In this picture, broad-line clouds are produced within the bow shocks surrounding the stellar wind sources.

In the Bondi-Hoyle accretion process, a bow shock forms around the black hole when it accretes from a rather uniform, laminar flow. The length scale of the shock is roughly the accretion radius itself, i.e., $d_{sw}^\text{(BH)} \approx (0.1-1)r_4 \approx (0.1-1)r$ (see § 1.2). We have already seen in Figures 1 and 2 that the requirement $d_{sw}^\text{(BH)} < r$ can be met for a wide variety of parameters over the range of relevant radii. Therefore, Bondi-Hoyle shocks around the central mass concentration are plausible sites for cloud production.

In Figure 3, we plot the density of the cooled (cloud) gas for a range of parameter values. Note that the gas displays the range in densities over radii inferred for the BLR. We find that the density at any given radius increases with $M$, but is insensitive to the values of $\epsilon$ and $T_C$. Note also that the density increases with decreasing $r$; this is consistent with the results of modeling the BLR using photoionization codes (e.g., Kaspi & Netzer 1999).

4.2. Bondi-Hoyle Accretion Shock Model

In the Bondi-Hoyle accretion process, a bow shock forms around the black hole when it accretes from a rather uniform, laminar flow. The length scale of the shock is roughly the accretion radius itself, i.e., $d_{sw}^\text{(BH)} \approx (0.1-1)r_4 \approx (0.1-1)r$ (see § 1.2). We have already seen in Figures 1 and 2 that the requirement $d_{sw}^\text{(BH)} < r$ can be met for a wide variety of parameters over the range of relevant radii. Therefore, Bondi-Hoyle shocks around the central mass concentration are plausible sites for cloud production.

In Figure 4, we plot $d_{sw}^\text{(BH)}$ using the fiducial values $v_o \approx 2000$ km s$^{-1}$ and $E_{36} \approx 100$, typical for Wolf-Rayet stars (although these are probably upper limits for a typical stellar population). It can be seen that the size of these shocks is probably too small for these to be viable sites for cloud production via radiative cooling, thus confirming the PD85 result.

![Graph](image-url)
shock have roughly parallel velocities. In addition, the Bondi-Hoyle shock does not provide a sufficiently broad distribution of cloud properties inferred for an extended BLR. It is therefore unlikely that a single Bondi-Hoyle accretion shock could produce the broad-line profiles seen in AGN spectra.

4.3. Wind Collision and Turbulent Accretion Shock Model

The final source of BLR clouds we consider here is shocks produced by large-scale wind collisions and turbulence within the overall flow. The motivation for this is that realistic three-dimensional simulations of the accretion onto a massive nucleus from a distribution of wind sources (Coker & Melia 1997) indicate that a single Bondi-Hoyle bow shock is difficult to form or maintain. Instead, the stellar wind-wind collisions produce an array of shock segments and a consequent turbulent inflow toward the black hole. In this picture, clouds are produced continually throughout the extended BLR, so we avoid the problem of having to confine long-lived clouds; instead, clouds that evaporate upon leaving the shock region are continually replaced by newly formed clouds at other locations within the inflow.

The shock regions must be large to allow cooling to occur (cf. Figs. 1 and 2), but these are readily obtainable for a realistic flow. Because the shocks, and therefore the clouds themselves, are embedded within the overall accretion pattern, the velocity of the clouds is roughly equal to that of the captured wind; i.e., \( v_c \approx v_w \). Clouds that move at non-trivial velocities relative to the surrounding medium are subjected to disruption via Rayleigh-Taylor instabilities (Mathews & Ferland 1987). The winds, and therefore the clouds, display a \( v(r) \propto r^{-1/2} \) velocity field fully consistent with, e.g., the Peterson & Wandel (1999) conclusion that the BLR velocity fields in NGC 5548 mimic Keplerian motions about a single central mass. Assuming that a large number of shocks exist at different locations within the flow, it should be possible to reproduce the observed line profiles.

5. EMISSION LINE MODELING

The most important diagnostic for any dynamical model of the BLR is the emission line shape. As an illustrative example, we use the cloud density and velocity distributions predicted by our model to fit an observed line profile. In our procedure, we assume that the physical properties and structure of the BLR are time-independent; i.e., the line variability is due to changes in the continuum only. This is true if the dynamical timescale of the BLR is much longer than the observation time of the reverberation study, typically on the order of a few hundred days. The dynamical timescale, approximated by

\[
\tau_{\text{dyn}} \sim \frac{r_{\text{BLR}}}{v_{\text{FWHM}}} \approx 600 \left( \frac{v_{10}}{v_{5000}} \right) \text{ days},
\]

where \( v_{5000} \equiv v_{\text{FWHM}}/(5000 \text{ km s}^{-1}) \) and \( v_{10} \equiv v_{\text{delay}}/(10 \text{ days}) \), is typically only slightly higher than the observation time. This indicates that some of the line profile variations seen in AGN time studies may in fact be due to changes in the BLR structure.

To model an emission-line profile, we begin by assigning values to the parameters \( M, \epsilon, T_c, \) and \( \sigma \). A cloud population is then constructed from the spatial distribution function, \( f(r, \theta, \phi) r^2 \sin \theta \, dr \, d\theta \, d\phi \), via Monte Carlo sampling techniques. With our assumption of spherical symmetry outside of the circularization radius, the number of clouds located between radii \( r \) and \( r + dr \) is proportional to \( f(r) r^2 dr \). Since the prediction of \( f(r) \) is beyond the scope of our simple model, we assign a power-law form, \( f(r) r^2 dr \propto r^{5/2} dr \), with the exponent \( \xi \) as a free parameter.

The density of a given cloud is calculated from equation (10), and its velocity corresponds to the local accretion flow velocity. By analogy with the Galactic center, we assume that the flow carries nonzero angular momentum with \( \sigma \approx 10-50 \). The velocity components are then identified as

\[
v_r(r) = -\frac{\omega_c R_S}{r}, \quad v_t(r) = -\sqrt{v_w^2 - v_r^2},
\]

where \( v_r \) and \( v_t \) are the tangential (angular) and radial velocities, respectively. In keeping with the picture of a turbulent accretion process, we assign \( v_t \) a random orientation perpendicular to the radial vector.

The cloud size as a function of radius is not predicted by our model. As a simple approximation, we assume that the cloud size scales with the radius at which it forms, i.e., \( l_c \propto r \). This assumption determines how the cloud cross-sectional area and column density scale with distance from the central mass in this picture.

CLOUDY is used to calculate the \( C_\text{IV} \lambda 1549 \) line emission intensity from each cloud. This intensity must be adjusted for both anisotropic cloud emission and Doppler boosting effects. Each cloud's emission is assigned to a wavelength bin determined by a combination of its Doppler, gravitational, and cosmological redshifts, and a total line profile is constructed by summing over the contributions from each cloud. See Appendix B for a more complete description.

5.1. An Illustrative Example: NGC 5548

As an illustrative example, we model the well-studied Seyfert 1 galaxy NGC 5548. We attempt to fit both the mean line profile and line light-curve data from the 1989 IUE campaign carried out by the AGN Watch Consortium (Clavel et al. 1991). The observational data consist of 60 spectra taken over a period of \( \sim 240 \) days. We extract the \( C_\text{IV} \lambda 1549 \) line profile from each spectrum by subtracting the underlying continuum using a least-squares linear fit to the surrounding continuum. An integrated line flux is then computed from each line profile, and a mean line profile is constructed as the unweighted average of all observations.

Our modeled data consist of line profiles (calculated as described above) at times corresponding to the real observation times. Note that the time-varying profile is dependent on the continuum flux at earlier times. For purposes of modeling, we assume that the ionizing luminosity is linearly proportional to the observed \( \lambda 1337 \) continuum flux (determined from the continuum light curve) as a function of time, and is constant for all times prior to the first observation. As in the case of the observed data, an integrated line flux is computed from each line profile, and a mean line profile is constructed from the entire set of modeled profiles. These are then compared to the observed data using \( \chi^2 \) analysis.

We set \( M_k = 0.68 \), consistent with the findings of Peterson & Wandel (1999), who used reverberation data from a set of emission lines to estimate the central mass in NGC 5548, and assume \( \epsilon = 0.1 \). For the SED, we use the table agn spectrum of CLOUDY, similar to the "typical AGN" spec-
trum deduced by Mathews & Ferland (1987), but having a submillimeter break at 10 μm. This spectrum has a Compton temperature of $T_c \approx 1$. We allow $\sigma$ to vary in the range 10–50, and $\xi$ is a free parameter. A Gaussian profile ($v_{\text{FWHM}} = 1000$ km s$^{-1}$) centered at $\lambda 1580$ is added to the modeled broad line to represent emission from the narrow-line region. The relative fluxes of the broad- and narrow-line components are additional free parameters.

In Figure 5, we show $\chi^2$ contour plots for fits to the line profile (left) and line light curve (right). Both plots suggest better fits at the lower allowed values of $\sigma$. However, the line profile data point to a value of $\xi \approx -3$, whereas the $\chi^2$ from the light-curve data decreases monotonically with smaller values of $\xi$. Because of this divergence in $\xi$, a search for better fits at smaller values of $\sigma$ is unwarranted.

Figure 6 shows fits to the mean line profile (left) and line light curve (right) using $\sigma = 15$ and $\xi = -3$. For the line profile, observational data are indicated by filled circles with error bars. The modeled broad-line component is indicated by a thin solid curve and the narrow-line component by a dotted line; the heavy solid curve represents their sum. We obtain $\chi^2 = 0.86$ (60 degrees of freedom) for this fit. In the second plot, the observed continuum and line light curves are indicated by open and closed circles, respectively, and the modeled line light curve is indicated by the solid curve. The mean of each light curve has been normalized to unity, and the line light curves have been offset from the continuum light curve for clarity. We obtain $\chi^2 = 1.32$ (56 degrees of freedom) for this fit.

It is perhaps not surprising that a good fit to the observed data can be obtained from our model by choosing appropriate values for the free parameters, particularly in light of the LOC model results. A real test of our model will require full three-dimensional hydrodynamical simulations to determine the size and distribution of shock-formed clouds as a function of radius.

6. CONCLUSIONS

In this paper, we have considered a broad range of possible gas configurations in a wind accreting onto the central black hole, with physical conditions that may produce BLR clouds via cooling instabilities within shocks. We note that in order to reproduce the observed line shape in actual sources, the BLR clouds cannot all be produced within a single outer region such as a Bondi-Hoyle shock, since this does not account for the required range in cloud properties at smaller radii. Instead, we have found that the best scenario involves local cloud production throughout the overall accretion flow. We conclude that a viable model for the formation of the BLR is one in which ambient gas sur-

![Fig. 5.—$\chi^2$ contour plots for model fits to the line profile (left) and light curve (right)](image1)

![Fig. 6.—Model fits to the line profile (left) and light curve (right) using $\sigma = 15$ and $\xi = -3$. See text for details.](image2)
Cloud emission is dependent on relativistic and anisotropic emission effects, as well as the time-retarded continuum luminosity. Taking the line of sight to be along the $z$-axis, the effects of anisotropic line emission (Ferland et al. 1992; O'Brien, Goad, & Gondhalekar 1994) and Doppler boosting are given by (Corbin 1997)

$$
\epsilon(r, v, t') = \frac{\epsilon_0(r, t')}{2} \left\{ 1 - [2\alpha(r, t') - 1] \cos \theta \right\} \left( \frac{1 - v^2/c^2}{1 + v_z/c} \right)^3,
$$

Finally, we consider the argument that broad-line emission cannot be produced by discrete clouds (Arav et al. 1998). This reasoning is based on the assumption that each cloud has a fixed set of parameters, such as density, thickness, velocity, etc. In our model, with the clouds continually forming in regions of high turbulence, each cloud region can display a wide range of properties. Therefore, we suggest that the cross-correlations that appear with fixed cloud properties would vanish. Given the viability of this picture, it now remains to be seen whether the vast array of BLR phenomena observed in sources ranging from Seyfert galaxies to high-redshift quasars can be self-consistently accounted for with this single description. This is work in progress and the results will be reported elsewhere.

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APPENDIX A

CALCULATION OF THE COOLING TIMESCALE

Following Krolik et al. (1981) and Perry & Dyson (1985), we calculate the cooling timescale for a shocked gas to cool from an initial temperature, $T_i$, to the equilibrium temperature, $T_r$, using the expression

$$
t_{cool} = \frac{\int_{T_i}^{T_r} dT}{F_c(T)},
$$

where $F_c(T)$ is the net cooling rate. From Matthews & Ferland (1987), we have

$$
F_c(T) = \frac{4\sigma_T}{3m_e c^2} \frac{L_{bol}}{4\pi r^2} (T - T_c) + \frac{n\lambda_b T^{1/2}}{3k}.
$$

The first term is due to Compton heating/cooling processes, and the second term is due to bremsstrahlung (free-free) emission.

Equation (A2) is valid for $T \gtrsim 10^5$ K. Below this temperature, collisional and radiative transitions in the plasma cause very rapid cooling. Therefore, $T_r = 10^5$ K is taken as the lower limit of integration for the cooling processes described in this paper.

APPENDIX B

THE LINE PROFILE

The composite emission-line flux from a population of $N$ clouds as a function of wavelength and time is

$$
F_L(\lambda, t) \propto \sum_{i=1}^{N} A_i(r_i) \epsilon(r_i, v, t') T(r_i, t, t') \Lambda(r_i, v, \lambda),
$$

where $A_i(r_i)$ is the average cross sectional cloud area as a function of radius, $\epsilon(r_i, v, t')$ is the line emissivity including relativistic and anisotropic emission effects, $T(r_i, t, t')$ accounts for the geometrical delay between the cloud emission and observation time, and $\Lambda(r_i, v, \lambda)$ is a convolution term relating a cloud’s location and velocity to the observed line wavelength.

B1. LINE EMISSIVITY

Cloud emission is dependent on relativistic and anisotropic emission effects, as well as the time-retarded continuum luminosity. Taking the line of sight to be along the $z$-axis, the effects of anisotropic line emission (Ferland et al. 1992; O’Brien, Goad, & Gondhalekar 1994) and Doppler boosting are given by (Corbin 1997)

$$
\epsilon(r, v, t') = \frac{\epsilon_0(r, t')}{2} \left\{ 1 - [2\alpha(r, t') - 1] \cos \theta \right\} \left( \frac{1 - v^2/c^2}{1 + v_z/c} \right)^3,
$$
where $\epsilon_0(r, t')$ is the rest-frame emission, $a(r, t') \equiv \epsilon_{\text{ion}}/\epsilon_{\text{tot}}$ is the emission fraction that escapes the illuminated cloud face, and
\[ v' = v + c\sqrt{1 - R_\lambda/r - 1} \]  
(B3)
is the line-of-sight velocity including general-relativistic corrections.

The photoionization code CLOUDY (Ferland 1996; CLOUDY 90.04) was used to calculate both $\epsilon_0(r, t')$ and $a(r, t')$. CLOUDY calculates line intensities based the SED, $n_e$, $U$, and $N_H$. To calculate $U$ from a real, varying continuum source, we assume that the continuum is radiated isotropically and that the ionizing luminosity $L_{\text{ion}}(t')$ seen by the cloud is proportional to the continuum light curve $C(t')$ at retarded time $t'$,
\[ L_{\text{ion}}(t') = C(t')T_{\text{ion}}, \]  
(B4)
where $T_{\text{ion}}$ is the mean ionizing luminosity and $\overline{}$, the mean value of the continuum light curve, has been normalized to unity.

B2. GEOMETRICAL TIME DELAY AND OBSERVED WAVELENGTH

Radiation emitted at time $t'$ by a cloud located at $(r, \theta)$ will be observed at a delayed time $t = t' + (1 + \cos \theta)r/c$ (Peterson 1997). Therefore, $T(r, t, t')$ can be expressed as a delta function:
\[ T(r, t, t') = \delta \left( t' - \left[ t - \frac{r(1 + \cos \theta)}{c} \right] \right). \]  
(B5)
The observed wavelength $\lambda$ of line emission from a given cloud is related to rest-frame line wavelength $\lambda_0$ by
\[ \lambda = \lambda_0 \left(1 - \frac{R_\lambda}{r}\right)^{-1/2} \left(\frac{1 - v^2/c^2}{1 - v_h/c}\right)(1 + z), \]  
(B6)
where the correction terms on the right are the gravitational, Doppler, and cosmological redshifts, respectively. This relation, appearing as a delta function $\Delta(r, \psi, \lambda)$ in equation (B1), is used to map each cloud’s emission into predetermined wavelength bins.

REFERENCES

Alexander, T., & Netzer, H. 1994, MNRAS, 270, 781
Allen, D., Hyland, A., & Hillier, D. 1990, MNRAS, 244, 706
Arav, N., Barlow, T., Laor, A., Sargent, W. L. W., & Blandford, R. D. 1998, MNRAS, 297, 990
Baldwin, J. 1997, in IAU Colloq. 159, Emission Lines in Active Galaxies: New Methods and Techniques, ed. B. M. Peterson, et al. (ASP Conf. Ser. 113; San Francisco: ASP), 80
Baldwin, J., Ferland, G., Korista, K., & Verner, D. 1995, ApJ, 455, L119
Bechtold, J., Weymann, R. J., Lin, Z., & Malkan, M. A. 1986, ApJ, 301, 162
Bondi, H., & Hoyle, F. 1944, MNRAS, 104, 273
Brotherton, M. S. 1996, ApJS, 102, 1
Cassidy, J., & Raine, D. J. 1996, A&A, 310, 49
Clavel, J., et al. 1991, ApJ, 366, 64
Collier, S. J., et al. 1995, ApJ, 451, 498
Corbin, M. R. 1997, ApJ, 485, 517
Dalgarno, A., & Lazyer, D. 1987, Spectroscopy of Astrophysical Plasmas (Cambridge: Cambridge Univ. Press)
Dietrich, M., et al. 1998, ApJS, 115, 185
Emmering, R. T., Blandford, R. D., & Shlosman, I. 1992, ApJ, 385, 460
Fabian, A. C., Guilbert, P. W., Arnaud, K. A., Shafer, R. A., Tennant, A. F., & Ward, M. J. 1986, MNRAS, 218, 457
Ferland, G. J. 1996, HAZY, a Brief Introduction to CLOUDY (Univ. Kentucky, Dept. Phys. & Astron. Internal Rept.)
Ferland, G. J., Peterson, B. M., Horne, K., Welsh, W. F., & Nahar, S. N. 1992, ApJ, 387, 95
Geballe, T., Krisicinas, K., Bailey, J., & Wade, R. 1991, ApJ, 370, L73
Genzel, R., Eckart, A., Ott, T., & Eisenhauer, F. 1997, MNRAS, 291, 219
Genzel, R., Thatte, N., Krabbe, A., Kroker, H., & Tacconi-Garman, L. E. 1996, ApJ, 472, 153
Ghez, A. M., Klein, B. L., Morris, M., & Becklin, E. E. 1998, ApJ, 509, 678
Glenn, J., Schmidt, G. D., & Foltz, C. B. 1994, ApJ, 434, L47
Green, P. J. 1996, ApJ, 467, 61
Guilbert, P. W. 1986, MNRAS, 218, 171
Hall, D., Kleinmann, S., & Scoville, N. 1982, ApJ, 260, L53
Jackson, J. M., Geis, N., Genzel, R., Harris, A. I., Madden, S., Poglitsch, A., Stacey, G. J., & Townes, C. H. 1993, ApJ, 402, 173
Kaspi, S., & Netzer, H. 1999, ApJ, 524, 71
Korista, K., Baldwin, J., & Ferland, G. 1998, ApJ, 507, 24
Korista, K. T., et al. 1995, ApJS, 97, 285
Krabbe, A., Genzel, R., Drapatz, S., & Rotaciuc, V. 1991, ApJ, 382, L19
Krolik, J. H., McKee, C. F., & Tarter, C. B. 1981, ApJ, 249, 422
Marnichi, P., Sulentic, J. W., Dultzin-Hacyan, D., Calvani, M., & Moles, M. 1996, ApJS, 104, 37
Matthews, W. G., & Ferland, G. J. 1987, ApJ, 323, 456
Melia, F. 1994, ApJ, 426, 575
Melia, F. 1998, in Proc. 1998 Workshop on the Galactic Center, ed. H. Falcke et al. (Berlin: Springer), in press
Mezger, P. G., Duschl, W. J., & Zylka, R. 1996, A&A Rev., 7, 289
Morris, R., & McCammon, D. 1983, ApJ, 270, 119
Murray, N., & Chiang, J. 1995, ApJ, 454, L105
Murray, N., Chiang, J., Grossman, S. A., & Voit, G. M. 1995, ApJ, 451, 498
Netzer, H. 1990, in Active Galactic Nuclei, ed. T. J.-L. Courvoisier & M. Mayor (Berlin: Springer), 67
O’Brien, P. T., Goad, M. R., & Gondhalekar, P. M. 1994, MNRAS, 268, 845
O’Brien, P. T., et al. 1998, ApJ, 509, 163
Osterbrock, D. E. 1989, Astrophysics of Gaseous Nebulae and Active Galactic Nuclei (Mill Valley: University Science Books)
Peterson, J. J., & Dyson, E. J. 1985, MNRAS, 213, 665
Peterson, B. M. 1997, An Introduction to Active Galactic Nuclei (Cambridge: Cambridge Univ. Press)
Peterson, B. M., & Wandel, A. 1999, ApJ, 512, 195
Peterson, B. M., et al. 1999, ApJ, 510, 659
Reichert, G. M., et al. 1994, ApJ, 425, 582
Reynolds, C. S., & Fabian, A. C. 1995, MNRAS, 273, 1167
Robinson, A. 1995, MNRAS, 276, 934
Rodriguez-Pascual, P. M., et al. 1997, ApJS, 110, 9
Ruffert, M., & Melia, F. 1994, A&A, 288, L29
Santos-Lleo, M., et al. 1997, ApJS, 112, 271
Sheils, J. C., Ferland, G. J., & Peterson, B. M. 1995, ApJ, 441, 507
Stirpe, M. G., et al. 1994, ApJ, 425, 609
Wandel, A. 1997, ApJ, 490, L131
Wandel, A., Peterson, B. M., & Malkan, M. A. 1999, ApJ, 526, 579
Wanders, I., et al. 1997, ApJS, 113, 69
Yusef-Zadeh, F., & Melia, F. 1992, ApJ, 385, L41