PART 3.

Fundamental Physical Processes

Advective Disks

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Abstract. Recent work on advection-dominated accretion flows (ADAFs) is reviewed. The article concentrates on an optically thin branch of ADAFs which is present at mass accretion rates below a critical value $\sim (10^{-2} - 10^{-1})$ the Eddington rate. Models based on this branch have been quite successful at explaining a number of low-luminosity X-ray binaries and galactic nuclei, and some brighter systems. Some progress has also been made toward understanding the various spectral states of accreting black holes. It is argued that ADAFs may provide one of the best techniques for demonstrating the reality of event horizons in black holes.

1. Introduction

The physics of accretion is embodied in a set of basic equations which describe the conservation of mass, momentum, and energy, coupled with relations for viscosity, energy generation, and energy transfer. Solutions of these equations are much valued since they can be used to model systems we observe in astrophysics. Four solution branches have been discovered and studied so far (Chen et al. 1995 discuss the relationships among the branches):

1. Thin accretion disk solution (Shakura & Sunyaev 1973, Novikov & Thorne 1973): This classic solution is by far the best-known of the four branches and has been used for modeling a large number of systems (see Pringle 1981 and Frank, King & Raine 1992 for reviews).

2. Shapiro, Lightman & Eardley (1976, hereafter SLE) solution: This hot, optically thin branch makes use of a two-temperature plasma and produces power-law spectra at X-ray and γ-ray energies, consistent with observations of accreting black holes and neutron stars. However, the solution is violently unstable (Piran 1978).

3. Optically thick advection-dominated solution (Abramowicz et al. 1988): At sufficiently high mass accretion rates, a new branch of solution is present where most of the radiation emitted by the accreting gas is trapped in the flow and is advected into the central star. This regime of accretion was originally considered by Katz (1977) and Begelman (1978), but a full analysis was first presented in an important paper by Abramowicz et al. (1988).

4. Optically thin advection-dominated solution (Narayan & Yi 1994, 1995ab, Abramowicz et al. 1995): At sufficiently low mass accretion rates, another advective solution branch occurs. Here, the accreting gas has such a low density that it is unable to cool efficiently, and so the dissipated energy is advected into

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the central star as thermal energy. Advection of thermal energy is standard in models of spherical accretion (Shvartsman 1971, Ipser & Price 1977). The idea was mentioned in the context of rotating flows by Rees et al. (1982), but this mode of accretion was established as an independent solution branch only recently with the work of Narayan & Yi (1994, 1995ab) and Abramowicz et al. (1995). (See Narayan 1996b for a review.)

The present article reviews advection-dominated accretion flows (ADAFs), concentrating in particular on optically-thin ADAFs around black holes.

2.   Dynamics of Accretion Flows

2.1.   The Basic Equations

Consider a steady axisymmetric accretion flow. Mass conservation gives

$$\dot{M} = -4\pi R H \rho v = \text{constant}, \qquad H = C_H(c_s/\Omega_K), \qquad (2.1)$$

where $\dot{M}$ is the mass accretion rate, R is the radius, H is the vertical scale height, ρ is the mid-plane density, and v is the radial velocity. H is expressed in terms of the local isothermal sound speed $(c_s^2 = p/\rho)$ and the Keplerian frequency Ω_K, with a proportional constant C_H $(= (5/2)^{1/2}$ in the analysis of Narayan & Yi 1995b, hereafter NY95b).

In height-integrated form, the conservation laws of radial momentum, angular momentum and energy give the following ordinary differential equations:

$$v\frac{dv}{dR} = -\Omega_K^2 R + \Omega^2 R - \frac{1}{\rho}\frac{d}{dR}(\rho c_s^2), \qquad (2.2)$$

$$\frac{d\Omega}{dR} = \frac{v(\Omega R^2 - j)}{R^2 \nu}, \qquad (2.3)$$

$$\rho v T \frac{ds}{dR} = \rho \nu R^2 \left(\frac{d\Omega}{dR}\right)^2 - q^-, \qquad (2.4)$$

where ν is the kinematic coefficient of viscosity, $j\dot{M}$ is the inward flux of angular momentum (j is a constant independent of R), T is the temperature, s is the entropy, and q^- is the radiative cooling rate per unit volume. It is standard to write $\nu = \alpha c_s^2/\Omega_K$, with α assumed to be a constant.

The calculation of q^- requires a detailed analysis of radiation emission mechanisms and radiative transfer. In this paper, we are primarily interested in radiation processes in optically thin gases.

2.2.   Advection-Dominated Accretion

Let us write the energy equation (2.4) compactly as follows:

$$q^{\text{adv}} = q^+ - q^-, \qquad (2.5)$$

where q^+ refers to the viscous heating of the gas and q^{adv} is a term which describes advective transport of energy (usually a form of cooling).

Equation (2.5) shows that three regimes of accretion are possible, depending on the relative magnitudes of the three terms:

- $q^+ \approx q^- \gg q^{\mathrm{adv}}$. This corresponds to a cooling-dominated accretion flow, where the energy released via viscous dissipation is mostly radiated locally. The thin accretion disk and the SLE solution belong to this regime.

- $q^{\mathrm{adv}} \approx q^+ \gg q^-$. This corresponds to the two ADAF branches. Most of the viscous energy is advected with the gas.

- $-q^{\mathrm{adv}} \approx q^- \gg q^+$: Here we have a cooling flow (or Kelvin-Helmholtz contraction in the case of a star). Energy generation is negligible, but entropy is steadily converted to radiation as the gas accretes.

2.3. Dynamics of ADAFs — Self-Similar Solutions

Returning to the energy equation (2.4), let us write the right-hand side as $f(R)q^+$, where $f(R) = (q^+ - q^-)/q^+$. The actual form of f depends on the details of the cooling, but in the specific case of an ADAF we have $f(R) \to 1$ so that it is independent of R. Narayan & Yi (1994) showed that equations (2.1)–(2.4) then have an exact analytical self-similar solution. (A version of the solution was discovered earlier by Spruit et al. 1987 in a different context.) In the limit when $\alpha \ll 1$, the self-similar solution takes a particularly simple form (see the original papers for the full solution):

$$v \approx -\frac{3\alpha}{(5+2\epsilon)}\Omega_K R, \quad \Omega \approx \left(\frac{2\epsilon}{5+2\epsilon}\right)^{1/2}\Omega_K, \quad c_s^2 \approx \frac{2}{(5+2\epsilon)}\Omega_K^2 R^2, \quad (2.6)$$

where $\epsilon = (5/3 - \gamma)/(\gamma - 1)$, and γ is the ratio of specific heats. This solution reveals several interesting features of ADAFs.

Since γ is likely to be in the range 4/3 to 5/3, ϵ lies in the range 0 to 1. If $\alpha \gtrsim 0.1$ as some arguments seem to suggest (Narayan 1996a), then we see that the radial velocity is comparable to the free-fall velocity, $v \sim v_{ff} = \Omega_K R$. (The velocity is nevertheless subsonic.) The gas thus accretes much more rapidly than in a thin disk, where $v \sim \alpha(H/R)^2 v_{ff}$. The angular velocity is distinctly sub-Keplerian, and in fact tends to 0 as $\gamma \to 5/3$. What holds the gas up against gravity if it is sub-Keplerian? The answer is the pressure gradient, which can be quite large in ADAFs as seen from the fact that $c_s^2/v_{ff}^2 \sim 0.3 - 0.4$. The gas temperature is correspondingly quite high; in fact, optically thin ADAFs have almost virial temperatures. This is not unexpected as all the energy released during accretion remains in the gas as internal energy.

Since $c_s \sim v_{ff}$ the scale height is large: $H/R \sim 1$ (see eq. 2.1). Can the height-integrated equations (2.1)–(2.4) be trusted under such extreme conditions? Narayan & Yi (1995a) considered a general set of equations without height-integration and discovered another self-similar ADAF solution. Using this solution they showed that the height-integrated equations are a surprisingly good approximation to the exact problem even when $H \sim R$.

2.4. Global Solutions

To proceed beyond self-similarity, it is necessary to solve the full equations (2.1)–(2.4) with proper boundary conditions. This has been done for the case of accretion on to a black hole by several groups (Narayan, Kato & Honma 1997a, Chen, Abramowicz & Lasota 1997, Nakamura et al. 1996). These authors
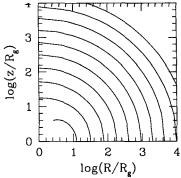
Figure 1.   Isodensity contours (two per decade of density) of a global ADAF with $\alpha = 0.3$ and $\gamma = 1.5$ (based on Narayan et al. 1997a, Appendix A). The contours are truncated at the sonic radius. The axes are expressed in units of the Schwarzschild radius, $R_g = 2GM/c^2$.

employ a pseudo-Newtonian potential (Paczyński & Wiita 1980) to simulate a Schwarzschild black hole, and obtain global transonic solutions for various choices of α and γ. They solve self-consistently for the sonic radius R_s (where the accreting gas goes supersonic) and the parameter j.

When $\alpha \gtrsim 0.01$, the global solutions reveal some unexpected features: (i) they remain sub-Keplerian at all radii, and (ii) they have no pressure maximum outside the sonic radius. Although these results were already apparent in the work of Matsumoto, Kato & Fukue (1985) and Narayan & Yi (1995a), they were nevertheless considered a surprise. Prior to the recent work, it had been widely assumed that when an ADAF approaches close to a black hole it would take the form of a thick torus with open funnels. The new results, however, suggest that ADAFs with moderately large α are not toroidal at all; in fact, ADAFs may be more similar to slowly rotating settling stars, as Fig. 1 shows.

Abramowicz et al. (1996) have extended the pseudo-Newtonian work to a full general relativistic analysis and have presented global solutions in Kerr geometry. Save for some technical issues related to causality in viscous interactions (see comments in Narayan et al. 1997a), it appears that the dynamics of ADAFs within the height-integrated approximation are now fully understood.

Much work remains to be done, however, on the two-dimensional structure in the Rz plane. Figure 1 is only an approximate representation, and real flows may well have complex circulation patterns and jets which cannot be captured by the approximations used so far.

Narayan & Yi (1994, 1995a) argued that ADAFs may be especially susceptible to producing bipolar outflows, but no consistent solutions have been found yet. ADAFs are also convectively unstable, especially along the axis (Narayan & Yi 1994, 1995a)), as confirmed by Igumenshchev et al. (1996). The physical consequences of the convection are not fully understood.

2.5.   Are There Shocks in ADAFs?

The global solutions described in §2.4 are free of shocks. However, Chakrabarti and his collaborators (cf. Chakrabarti & Titarchuk 1995) have claimed that

shocks are generic to ADAFs. Initially, Chakrabarti (1990) considered viscous flows under isothermal conditions, but more recently, with the increased interest in ADAFs, he has switched his attention to adiabatic flows. The clearest account of his results are found in Chakrabarti & Titarchuk (1995) where the authors claim: (i) low α flows ($\alpha \lesssim 0.01$) have shocks and high α ($\gtrsim 0.01$) flows do not, and (ii) low α flows have sub-Keplerian rotation at large radii while high α flows are Keplerian at all radii except very close to the black hole. Neither statement is confirmed in the work described in §2.4. In particular, no shocks are seen in any of the global solutions, which span a wide range of parameter values: α ranging from 10^{-3} to 0.3, and γ from 4/3 to 5/3.

What is the source of the discrepancey? Narayan et al. (1997a) suggested that it may lie in different philosophies regarding the angular momentum parameter j in equation (2.3). Chakrabarti simply assigns a value to j (in fact, different values in different papers) whereas this parameter ought to be treated as an eigenvalue and determined self-consistently through boundary conditions (see Popham & Narayan 1991 and Paczynski 1991).

In a viscous accretion flow the gas starts out with considerable angular momentum at the outer edge, but very little of that angular momentum is accreted by the central star since most of it is transported outward by viscosity. How much reaches the center is determined by the properties of the star. More specifically, there is a radius near the inner edge where the flow satisfies a "no-torque" condition, and $j = \Omega R^2$ at this radius. This well-known argument leads in the case of a thin disk to the result $j = \Omega_K(R_{in})R_{in}^2$, where R_{in} corresponds to the inner edge of the Keplerian flow. For more general global flows, it is not possible to write an explicit formula for j since neither the sonic radius nor the run of $\Omega(R)$ is known ahead of time. Instead, j is determined self-consistently by requiring the solution to satisfy boundary conditions on the inside appropriate to the properties of the star. Clearly, if one sets j equal to an arbitrary value the flow is likely to become self-inconsistent, perhaps leading to a shock. Indeed, self-inconsistency is likely because the value of j is known to be extremely insensitive to outer boundary conditions (cf. Narayan et al. 1997a), which means that a randomly chosen value will not correspond to any reasonable flow.

At a deeper level, if j is a global eigenvalue set by the star, how does the flow on the outside know the value of j? The answer is that information is transported via viscous stresses (cf. Popham & Narayan 1992). However, what happens when α is very small? Is it possible that when the viscous stress is weak information does not make it to the outside, thus causing a shock? No, because eq. (2.6) shows that $v \propto \alpha$. The smaller α gets, the slower the radial velocity becomes, and the longer the gas spends within the flow. Therefore, regardless of the value of α information always makes it back to the outside via viscosity, and the arguments of the previous paragraph are always valid. Some caveats are necessary, however. If $\alpha = 0$ exactly, rather than $\alpha \to 0$, then there is no angular momentum transfer at all and shocks are indeed generic. Also, if the flow varies on a time scale shorter than the viscous time (R/v), shocks may be possible. Finally, shocks do occur when inflowing gas first meets a rotating flow, e.g. when the accretion stream in a binary hits the disk. We are not concerned with such shocks.

3.   Optically Thin ADAFs

Following Abramowicz et al. (1995), consider a simplified model where electrons cool purely by free-free emission. In such a gas, the heating and cooling rates vary as $q^+ \propto n_i$ and $q^- \propto n_i n_e T_e^{1/2}$, where n_i and n_e are the ion and electron number densities, and T_e is the electron temperature. Since $n_i \sim n_e \propto \dot{m}$, it is clear that below a critical $\dot{m}_{\mathrm{crit}}$, q^- will be unable to match q^+ even if T_e is equal to the local virial temperature T_{vir}. The accretion must then proceed as an ADAF. (Here, $\dot{m}$ is the mass accretion rate in Eddington units, $\dot{m} = \dot{M}/1.39 \times 10^{18} m$ g s^{-1}, where m is the black hole mass in solar mass units.)

In actual fact, once electrons become relativistic, they cool via additional channels, notably synchrotron radiation and inverse Compton scattering, which can be very efficient. Can an ADAF survive in the presence of these cooling processes? Yes, because the viscous heating primarily affects ions while the cooling is almost entirely by electrons. At low densities, the transfer of energy from ions to electrons via Coulomb collisions is quite inefficient, creating a bottleneck for efficient cooling. In fact, under these conditions the plasma is expected to become two-temperature with ion temperature $T_i \gg T_e$ (SLE). From NY95b, the rate of viscous heating of ions scales as

$$\left(\frac{dE_i}{dt}\right)_{visc} = 9\alpha m^{-1} r^{-5/2} \text{ ergs s}^{-1}, \tag{3.1}$$

and the rate of energy transfer from ions to electrons varies as

$$\left(\frac{dE_e}{dt}\right)_{Coul} = 6\alpha m^{-1} \left(\frac{\dot{m}}{\alpha^2}\right) r^{-5/2} T_{10}^{-1} \text{ ergs s}^{-1}, \tag{3.2}$$

where r is the radius in Schwarzschild units and $T_{10} = T_e/10^{10}$K (typically $T_{10} \sim 0.1 - 1$). From these relations we see that the critical $\dot{m}$ below which a two-temperature ADAF is possible is

$$\dot{m}_{crit} \sim 1.5 T_{10} \alpha^2 \sim (0.3 - 1)\alpha^2. \tag{3.3}$$

This estimate is consistent with detailed calculations (NY95b, Narayan 1996a).

Esin et al. (1996) have shown that even when there is efficient coupling between ions and electrons (say by collective processes in the plasma, cf. Begelman & Chiueh 1988) ADAFs are still allowed, but $\dot{m}_{crit}$ is much lower.

3.1.   Properties of Two-Temperature ADAFs

The ion temperature in a two-temperature ADAF is very high, $T_i \sim 10^{12}$K/r, and the electron temperature is also quite high, $T_e \sim (10^9 - 10^{10})$K. The optical depth to electron scattering is usually low: $\tau_{es} \sim 10\alpha(\dot{m}/\dot{m}_{crit})$ (NY95b). These properties imply that such flows should produce hard Comptonized spectra extending up to a few hundred keV. ADAFs are therefore very attractive for modeling observations of accreting black holes (§§4, 5). Despite the high temperature, pair processes are not important because the radiation density is low (Bjornsson et al. 1996, Kusunose & Mineshige 1996).

From eqs. (3.1), (3.2), we see that q^-/q^+ and the luminosity vary as

$$\frac{q^-}{q^+} = \frac{(dE_e/dt)_{Coul}}{(dE_i/dt)_{visc}} = \frac{\dot{m}}{\dot{m}_{crit}}, \quad \frac{L}{L_{Edd}} \sim \frac{q^-}{q^+}\dot{m} = \frac{\dot{m}^2}{\dot{m}_{crit}}. \tag{3.4}$$

Low $\dot{m}$ ADAFs are thus intrinsically very inefficient, and the efficiency drops rapidly with decreasing $\dot{m}$. Not surprisingly, the most successful applications of ADAF models have been to extremely low-luminosity sources.

Where does all the energy in an ADAF go? Ultimately, the advected energy is accreted by the central star. What happens then depends on the nature of the star. If we have a normal star with a surface, the absorbed thermal energy will be re-radiated and the overall radiative efficiency (when this re-radiation is also included) will be similar to that of a standard thin accretion disk. However, if the central star is a black hole, then the advected thermal energy disappears through the horizon and there is no re-radiation. The overall luminosity of the source is then significantly below what one expects for the given $\dot{m}$. This is a very specific signature of black hole accretion. Indeed, the firm demonstration that an accreting source has an ADAF and that it is anomalously underluminous can be construed as "proof" that the particular accreting star is a black hole with an event horizon (Narayan, McClintock & Yi 1996). This is one of the best astrophysical techniques available today to demonstrate the reality of horizons.

3.2. Stability

Although the low $\dot{m}$ ADAF branch and the SLE branch are both hot and optically thin, they differ in one important respect; the ADAF solution is stable (Abramowicz et al. 1995, NY95b), whereas the SLE solution is violently unstable (Piran 1978).

Consider vertically integrated quantities as a function of the ion temperature T_i in an accretion flow with a fixed $\dot{m}$. The heating rate per unit area, $Q^+ = 2Hq^+$, is directly proportional to $\dot{m}$ and is independent of T_i.

At sufficiently low T_i, when the gas is optically thick (and geometrically thin), the cooling rate Q^- increases with increasing T_i because the temperature gradient goes up and the optical depth goes down. Beyond a certain critical temperature, however, the gas goes optically thin and, with increasing T_i, the density decreases and so does Q^- (since the cooling varies as ρ^2).

Figure 2a shows a schematic plot of Q^+ and Q^-. We see that two solutions are in general possible which satisfy the condition $Q^+ = Q^-$ required for a cooling-dominated flow: (i) a cool optically thick solution, and (ii) a hot optically thin solution. The former corresponds to the standard thin accretion disk. This is a stable solution (at least for fixed $\dot{m}$), in the sense that a slight perturbation of T_i causes the system to return to its equilibrium state (indicated by the arrows in Fig. 2a). The second solution corresponds to the unstable SLE branch.

The advective cooling term Q^{adv} is equal to the divergence of the advective energy flux and is proportional to the local thermal energy density of the gas. Roughly we expect $Q^{adv} \sim Q^+(T_i/T_{vir})$ (see NY95b for more details), where $T_{vir} \sim m_p c^2/kr$ is the virial temperature of the protons. Thus, when $T_i \sim T_{vir}$ it is possible to have an advection-dominated solution satisfying $Q^{adv} \sim Q^+$.

Figure 2b shows the situation when both cooling and advection are present. Now there are three solutions, of which two (the thin disk and the ADAF)
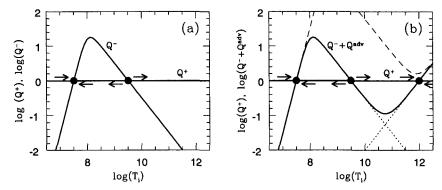
Figure 2.   (a) Shows Q^+ and Q^- (normalized to Q^+) as functions of T_i. The two dots represent equilibrium cooling-dominated solutions. The arrows indicate the direction of evolution of a non-equilibrium system; when $Q^+ > Q^-$, the system heats up, and when $Q^+ < Q^-$ it cools. The solution on the left (thin disk) is stable, while the one on the right (SLE) is unstable. (b) Includes advection. A third solution is present, which is hotter than SLE, and stable. The dashed line corresponds to a system with $\dot{m} > \dot{m}_{crit}$. Only one solution survives.

are stable, and one (SLE) is unstable. At sufficiently high $\dot{m}$ ($> \dot{m}_{crit}$) the SLE and ADAF solutions merge and disappear and we are left with only the thin disk branch. But for a wide range of $\dot{m} < \dot{m}_{crit}$, both the thin disk and ADAF are allowed. Which of these two stable solutions does a system select? NY95b discuss a number of possibilities; the current evidence suggests that nature prefers the ADAF branch whenever $\dot{m}$ is low (§4).

The above discussion refers to long-wavelength perturbations at constant $\dot{m}$. Kato, Abramowicz & Chen (1996) have carried out a more complete linear perturbation analysis and conclude that the low $\dot{m}$ ADAF branch may have weakly unstable modes. However, the modes do not grow sufficiently rapidly to threaten the global viability of the solution. Indeed, Manmoto et al. (1996) use the instabilities to explain X-ray shots in Cygnus X-1.

4.   Applications to Low Luminosity Systems

ADAF solutions have been applied to several low luminosity black holes, and model spectra have been fitted to observational data. Detailed codes have been developed for this purpose which include synchrotron and bremsstrahlung emission and Comptonization (see Narayan, Barret & McClintock 1997b for details). Some earlier models used the self-similar solution to estimate the run of density and pressure (§2.3), but the recent work employs global solutions (§2.4).

The models have been quite successful on two counts. First, the computed spectra are nearly always consistent with the available observations, though the data are usually meager as the sources are faint. Second, some of the systems have independent estimates of $\dot{m}$ and the sources appear to be extremely underluminous: $L/L_{Edd} \sim (10^{-5} - 10^{-3})\dot{m}$. An ADAF can clearly explain this.
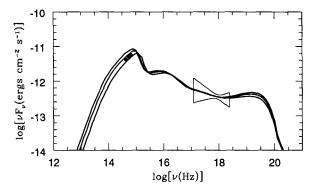
Figure 3.   The three dots show the optical flux of V404 Cyg in quiescence, and the "bow-tie" is the X-ray error box (2σ) from ASCA data. The solid lines are the spectra of three ADAF-based models with black hole masses (from left) of 16, 12, $8M_\odot$. (From Narayan et al. 1997b)

The following subsections describe various applications of ADAF models to low-luminosity systems. To summarize, in every low-luminosity system for which spectral data exist, the ADAF model provides a plausible explanation of the observations. Thus, for the moment at least, it appears that the ADAF paradigm is well supported by observations. The evidence for event horizons may thus be considered reasonably strong (§3.1).

4.1.   Quiescent X-ray Binaries

Narayan et al. (1996) showed that some puzzling observations of the soft X-ray transient (SXT) source, A0620–00, in its quiescent state could be explained with a model in which the accretion flow has two zones: (i) an outer zone ($r > r_{tr} \sim 10^3$–10^4) consisting of a thin disk, and (ii) an inner zone ($r < r_{tr}$) consisting of an ADAF. The model explains the available spectral data, and provides a plausible explanation for the recurrence time between outbursts (Lasota, Narayan & Yi 1996).

Narayan et al. (1996) also modeled the quiescent SXT, V404 Cyg, and predicted that it would have a hard X-ray spectrum with a photon index ~ 2. This has now been confirmed with ASCA data (Narayan et al. 1997b). Figure 3 shows a comparison between the model and observations. The fit is very good. The model is quite robust and fits the data for a wide range of black hole mass, binary inclination, transition radius r_{tr}, viscosity α, and magnetic field strength. The model also avoids several potential problems emphasized by Wheeler (1996).

V404 Cyg is the brightest known quiescent SXT and provides the best test of the ADAF paradigm. The evidence thus far is very encouraging.

4.2.   Quiescent Galactic Nuclei

The first galactic nucleus to which an ADAF model was applied was Sagittarius A* (Narayan, Yi & Mahadevan 1995, see also Rees 1982), the $\sim 10^6 M_\odot$ supermassive black hole at the center of our Galaxy. This source has long been known to be unusually dim for its estimated mass accretion rate ($L/L_{Edd} \sim 10^{-5}\dot{m}$).

The ADAF model provides a natural explanation of the low luminosity and also gives a reasonable fit to various flux measurements and upper limits of the source (ranging from radio to X-rays).

Following this work, Fabian & Rees (1995) applied the same model to other galaxies. Fabian & Canizares (1988) had earlier noted that the nuclei of bright ellipticals are much too dim considering the mass accretion rates estimated from cooling flows. If the accretion is via a thin accretion disk, the central black holes in these galaxies can be no larger than $\sim (10^7 - 10^8) M_\odot$, which conflicts with the much larger masses suggested by quasar evolution models. This problem can be reconciled if the black holes accrete through an ADAF (Fabian & Rees 1995, Mahadevan 1997).

5.   Applications to Luminous Systems

5.1.   X-ray Binaries

Narayan (1996b) constructed ADAF models with higher $\dot{m}$ than in quiescent systems and found that systems in which $\dot{m}$ approaches $\dot{m}_{crit}$ have very hard X-ray spectra with photon indices $\sim 1.5 - 2$. This led to the plausible suggestion that the so-called low state of black hole X-ray binaries may correspond to an ADAF flow with $\dot{m} \lesssim \dot{m}_{crit}$. Further, when $\dot{m}$ exceeds $\dot{m}_{crit}$, the outer thin disk moves in to smaller radii and the system switches to a pure thin accretion disk with a very soft spectrum. This corresponds to the high state. Thus the transition from low to high state corresponds to a switch from an ADAF to a thin disk. Chen & Taam (1996) have come up with a somewhat different proposal for the low-high transition.

Narayan & McClintock (1997) carried out a more detailed test of these ideas. Using a sequence of models with varying $\dot{m}$ they tried to explain the spectral evolution of Nova Muscae 1991 during outburst. The results are fairly encouraging (Fig. 4).

From these studies, the following five states of an accreting black hole can be identified, in order of increasing $\dot{m}$ (compare also with Nowak 1995): (i) At very low $\dot{m}$, the ADAF extends over a large range of r; if a thin disk is present at all, it is only at a large radius; this is the quiescent state. (ii) At higher $\dot{m}$, but still with $\dot{m} < \dot{m}_{crit}$, the model geometry is identical, but the luminosity is now much higher (recall eq. 3.4); this is the low state. (ii) For $\dot{m} \sim \dot{m}_{crit}$, the thin disk moves in toward smaller radii, and we have a situation where the ADAF extends over only a moderate range of r and the thin disk is energetically comparable to or even stronger than the ADAF; this is an intermediate state. (iv) At yet higher $\dot{m}$, the thin disk comes all the way down to the last stable orbit and there is only a weak corona; this is the high state. (v) Finally, as $\dot{m}$ approaches unity, the disk develops an energetically active corona (Haardt & Maraschi 1991), for reasons not yet understood; in this very high state, the corona may be viewed as an ADAF which happens to coexist with the thin disk. (Alternatively, an ADAF is just a corona without a disk.)

Ultra-relativistic jets appear to be associated with the very high state (e.g. GRS 1915+105 and GRO J1655–40 discussed at this meeting). The jets are most likely launched by the ADAF-like corona, which is consistent with the idea that ADAFs are prone to produce jets (§2.4). We would then expect the
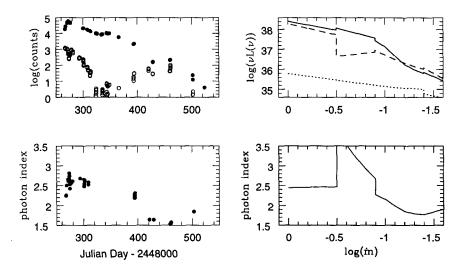
Figure 4. The top left panel shows the X-ray counts of the SXT, Nova Muscae 1991, during decline from outburst (Ebisawa et al. 1994); solid dots: 1.2–4.6 keV, open dots: 9.3–37 keV. The bottom left panel shows the observed photon spectral indices; the spectral index could not be measured between Day 320 and Day 390. The top right panel shows model fluxes (solid line: 1–6 keV, dashed: 6–40 keV, dotted: 2–10 eV), and the bottom right panel shows model spectral indices. The models reproduce the qualitative features seen in the data. (From Narayan & McClintock 1997)

intermediate state and low state also to have jets, though these outflows may be less energetic. Cyg X-1 and other low state black holes with nonthermal radio emission may be examples of this phenomenon. (Quiescent systems too should have outflows, but perhaps too weak to be observed.)

5.2. Galactic Nuclei

Lasota et al. (1996) made an ADAF model of NGC 4258 and showed that it explains the observed X-ray spectrum. Herrnstein et al. (this volume) show that the same model is also consistent with VLBI radio observations. Since the model gives a reasonably large $\dot{m} \sim 0.01$, NGC 4258 corresponds to a state that is in between the quiescent and low states of X-ray binaries. NGC 4258 is a LINER galaxy and Lasota et al. suggested that perhaps all LINERs have ADAFs. In terms of $\dot{m}$, LINERS probably range all the way from the quiescent state up to the intermediate state.

Carrying through the analogy with X-ray binaries, one expects that active galactic nuclei too should have multiple states. We have already mentioned the quiescent, low and intermediate states above. Seyfert galaxies probably span the range from the intermediate state to the low end of the high state, while bright quasars may correspond to the high and very high states. Again, in analogy with X-ray binaries, we would expect the most energetic jets to occur in quasars in

the very high state (these would correspond to FR II systems), while we would have lower power jets in systems corresponding to the low and intermediate states (FR I). In this picture, radio quiet quasars would correspond to the high state, where there is neither a separate ADAF nor an active corona; in terms of $\dot{m}$, radio quiet quasars should lie in between FR II and FR I systems.

Yi (1996) has argued that the break in the quasar luminosity function at redshift $z \sim 2$ might be the result of quasars switching from thin disk accretion (at $z > 2$) to ADAF accretion ($z < 2$). Such an evolution is expected just from the increase of black hole mass with time, aided perhaps by a decrease in the mass accretion rate. The model naturally explains the lack of bright quasars at the present epoch.

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Discussion

M. Nowak: Two questions: In the low state of Galactic black hole candidates, what is the dominant source of soft photons for Comptonization? Do you include the irradiation of the cool disk from the Compton cloud? Second, you predict that the low-to-high state transition is the optically thick disk working its way inward, whereas other theories may predict the opposite (an optically thick disk forming in the center and working its way outward). Can this be distinguished with the recent Cyg X-1 high state observations? Can you think of other tests?

R. Narayan: (1) In the low state as well as the quiescent state, the dominant source of soft photons is synchrotron radiation from within the ADAF. We do include irradiation of the outer disk by radiation from the inner ADAF, as well as irradiation of the ADAF by the cool disk. These interactions are important in the intermediate state, but not in the quiescent and low states. (2) We have not tried to model the recent Cyg X-1 high state data, but there are several tests of our ideas that are potentially possible with SXTs. We predict that during the rise phase of an SXT outburst the inner edge of the thin disk moves in, and that late in the decline, once the system switches from the high state to the intermediate state, the disk moves out. The variation of the inner edge should be seen in multi-color blackbody fits of the soft component of the spectrum. If an X-ray iron line is seen, this should show the variation of the Keplerian velocity at the inner edge. Also, we have clear predictions for the variation of the X-ray reflection component.

A. King: Two questions: (1) What causes the SXT outburst in your model? (2) In the case when the accreting object is a neutron star, am I right in thinking that the advective flow will crash into the surface and give back all the advected energy? So these systems cannot be "secret guzzlers"?

R. Narayan: (1) I believe the outburst is caused by a thermal limit cycle in the outer thin disk, exactly as in current models (e.g. Mineshige & Wheeler, ApJ, 343, 241, 1989). However, the fact that the disk is truncated at a large transition radius r_{tr} will doubtless cause important changes (cf. Lasota, Narayan & Yi 1996) and needs to be modeled in detail. (2) Yes, there ought to be a clear difference between black hole SXTs and neutron star SXTs, especially in quiescence, since the former are "secret guzzlers" (cf. Abramowicz & Lasota, Comments Ap, 18, 141, 1995, who introduced the term) according to our model while the latter cannot be secret guzzlers (see §3.1). There is some evidence that, in quiescence, Cen X-4 (a NS SXT) is brighter than A0620–00 (a BH SXT), which is consistent with this expectation. However, detailed ADAF models have not yet been developed for NS SXTs and so it is not clear if the observed spectra can be fitted.

M. Livio: Your models do not produce a secondary maximum in the optical-UV, while these are observed in many systems.

R. Narayan: Our models do not show any effect in the optical-UV band (dotted line in Fig. 4) when the soft and hard X-rays go through their "glitch." The X-rays are from the ADAF/corona, and the optical-UV are from the cool disk.

Perhaps we could make our model produce a secondary maximum in the optical light curve if we change the way we model the irradiation of the disk.

M. Begelman: In your model for the different spectral states of black hole candidates, a transition from low state to high state should be accompanied by a large increase in L_{bol}, since $\dot{M}$ is increasing and the advection-dominated zone is shrinking. But in data shown by Zhang this morning, $L(1-200 \text{ keV})$ was constant during the transition of Cyg X-1 from low state to high state. Can this be accommodated within your scheme?

R. Narayan: I believe the Cygnus X-1 "high state" is not a true high state but more like what I have called the intermediate state. We would predict an increase in L_{bol} from the low state to this state, exactly as you say. However, if the thin disk remains at a moderately large radius, then much of its radiation may lie outside the 1-200 keV band, and the observed L in the X-ray band may well remain constant even while L_{bol} is going up. In our model of SXT outbursts, for instance, during the rise phase (not discussed in the article) we actually see the X-ray flux *fall* briefly during a period when the bolometric luminosity is rising. So, Zhang's observations could be accommodated in principle, but we would need to model his data in detail before we can tell for sure.