

A SELF-CONSISTENT MODEL FOR THE FORMATION OF RELATIVISTIC OUTFLOWS IN ADVECTION-DOMINATED ACCRETION DISKS WITH SHOCKS

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ABSTRACT

In this Letter, we suggest that the relativistic protons powering the outflows emanating from radio-loud systems containing black holes are accelerated at standing, centrifugally supported shocks in hot advection-dominated accretion disks. Such disks are ideal sites for first-order Fermi acceleration at shocks because the gas is tenuous, and consequently the mean free path for particle-particle collisions generally exceeds the thickness of the disk. The accelerated particles are therefore able to avoid thermalization, and as a result a small fraction of them achieve very high energies and escape from the disk. In our approach, the hydrodynamics and the particle acceleration are coupled and the solutions are obtained self-consistently based on a rigorous mathematical treatment. The theoretical analysis of the particle transport parallels the early studies of cosmic-ray acceleration in supernova shock waves. We find that particle acceleration in the vicinity of the shock can extract enough energy to power a relativistic jet. Using physical parameters appropriate for M87 and Sgr A*, we confirm that the jet kinetic luminosities predicted by the theory agree with the observational estimates.

Subject headings: accretion, accretion disks — black hole physics — galaxies: jets — hydrodynamics

1. INTRODUCTION

Radio-loud active galactic nuclei (AGNs) are thought to contain supermassive central black holes surrounded by hot, two-temperature, advection-dominated accretion flows (ADAFs) with significantly sub-Eddington accretion rates. In these disks, the ion temperature $T_i \sim 10^{12}$ K greatly exceeds the electron temperature $T_e \sim 10^{10}$ K (e.g., Narayan et al. 1997; Becker & Le 2003). The observed correlation between high radio luminosities and the presence of outflows suggests that the hot ADAF disks are efficient sites for the acceleration of the relativistic particles powering the jets. This motivated Blandford & Begelman (1999) to consider the possibility of strong outflows of matter and energy from such systems, under the assumption of self-similarity and Newtonian gravity. The same scenario was extended to the general relativistic case by Becker et al. (2001). However, a single, global, self-consistent model for the structure of the disk/jet that includes a specific microphysical acceleration mechanism has not yet appeared in the literature. In this Letter, we explore the consequences of the presence of a shock in an ADAF disk for the acceleration of the observed nonthermal jet particles.

It is well known that for certain sets of outer boundary conditions describing the specific angular momentum and the specific energy of the gas supplied to the accretion flow, both shocked and shock-free solutions for the global disk structure are possible, although there is an upper limit on the amount of viscosity beyond which no shocks can form. However, it is clear that shocks can exist in both dissipative and inviscid flows (e.g., Chakrabarti & Das 2004). When both shocked and shock-free solutions are possible for the same set of upstream conditions, one must incorporate additional physical principles in order to determine which flow solution actually occurs. Since in general the shock solution possesses a higher entropy content than the shock-free solution, we argue based on the second law

of thermodynamics that shocks are preferred (see Becker & Kazanas 2001). Particle acceleration has not been studied in shocked disks in either the viscous or the inviscid cases, and therefore we concentrate here on the simpler situation of inviscid (adiabatic) flow. In this Letter we focus on disks containing isothermal shocks, because the jump in the energy flux at such a shock can be identified with the energy carried away from the disk by the relativistic particles in the outflow.

2. DYNAMICAL MODEL

The dynamical structure of the accretion flows considered here follows the model studied by Chakrabarti (1989) and Abramowicz & Chakrabarti (1990), which comprises an inviscid disk that either contains an isothermal shock or is shock-free. The model is based on the standard one-dimensional conservation equations for ADAFs, including utilization of the pseudo-Newtonian gravitational potential to simulate the effects of general relativity (Paczynski & Wiita 1980). In general, the flow is assumed to be adiabatic and in vertical hydrostatic equilibrium, although the entropy can jump at the shock. Under the inviscid flow assumption, the disk/shock model depends on two parameters, namely, the energy transport rate per unit mass, ϵ , and the specific angular momentum, ℓ . The value of ϵ is constant in ADAF-type flows since there are no radiative losses, although it will jump at the location of the isothermal shock if there is one in the disk. On the other hand, the value of ℓ remains constant since the disk is inviscid. The energy transport rate per unit mass can be written as

$$\epsilon = \frac{1}{2} \frac{\ell^2}{r^2} + \frac{1}{2} v^2 + \frac{1}{\gamma - 1} a^2 - \frac{GM}{r - r_s}, \quad (1)$$

where v denotes the radial velocity (defined to be positive for inflow), $r_s = 2GM/c^2$, and the adiabatic sound speed is given by $a = (\gamma P/\rho)^{1/2}$ for a gas with pressure P , mass density ρ , and specific heat ratio γ . Note that ϵ is defined to be positive for an inward flow of energy into the black hole. Furthermore, in this semiclassical approach, ϵ does not include the rest-mass contribution to the energy.

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Dissipation and radiative losses are unimportant in inviscid advection-dominated disks, and therefore the flow is in general adiabatic and isentropic, except at the shock location. In this situation, the “entropy parameter,”

$$K \equiv v a^{(\gamma+1)/(\gamma-1)} r^{3/2} (r - r_s), \quad (2)$$

is conserved throughout the flow except at the location of a shock (Becker & Le 2003). When a shock is present, we use the subscripts “−” and “+” to refer to quantities measured just upstream and just downstream from the shock, respectively. In the isothermal shock model, K and ϵ have smaller values in the postshock region (K_+ , ϵ_+) compared with the preshock region (K_- , ϵ_-). The entropy and energy lost from the disk are carried off by the outflowing high-energy particles, and the total entropy of the combined (disk/outflow) system is higher when a shock is present, as expected according to the second law of thermodynamics. By combining equations (1) and (2) and differentiating with respect to radius, we can obtain a “wind equation” of the form $(1/v)(dv/dr) = N/D$, where

$$N \equiv -\frac{\ell^2}{r^3} + \frac{GM}{(r - r_s)^2} + \frac{a^2(3r_s - 5r)}{(\gamma + 1)r(r - r_s)}, \quad D \equiv \frac{2a^2}{\gamma + 1} - v^2. \quad (3)$$

Critical points occur where the numerator N and the denominator D vanish simultaneously (see Fig. 1b for a specific example). The associated critical conditions can be combined with the energy equation (1) to express the critical velocity v_c , the critical sound speed a_c , and the critical radius r_c as functions of ϵ and ℓ . The value of K is then obtained by substituting r_c , v_c , and a_c into equation (2). In general, one obtains four solutions for the critical radius, denoted by r_{c1} , r_{c2} , r_{c3} , and r_{c4} in order of decreasing radius. Only r_{c1} and r_{c3} are physically acceptable sonic points, although the types of accretion flows that can pass through them are different. Specifically, r_{c3} is an X-type critical point, and therefore a smooth, global, shock-free solution always exists in which the flow is transonic at r_{c3} and then remains supersonic all the way to the event horizon. On the other hand, r_{c1} is an α -type critical point, and therefore any accretion flow originating at a large distance that passes through this point must display a shock transition below r_{c1} in order to cross the event horizon (Abramowicz & Chakrabarti 1990).

The radius of the isothermal shock, denoted by r_* , is determined self-consistently along with the structure of the disk by satisfying the velocity and energy jump conditions (Chakrabarti 1989),

$$R_*^{-1} \equiv \frac{v_+}{v_-} = \frac{1}{\gamma \mathcal{M}_-^2}, \quad \Delta\epsilon \equiv \epsilon_+ - \epsilon_- = \frac{v_+^2 - v_-^2}{2}, \quad (4)$$

where $\mathcal{M}_- \equiv v_-/a_-$ is the upstream Mach number at the shock location and R_* is the shock compression ratio. Because of the escape of energy at the isothermal shock, the energy transport rate ϵ drops from the upstream value ϵ_- to the downstream value ϵ_+ , and consequently $\Delta\epsilon < 0$. The flow must therefore pass through a new inner critical point located at $\hat{r}_{c3} < r_*$, which is computed using the downstream value ϵ_+ . This structure represents the transonic solution of interest here. The kinetic

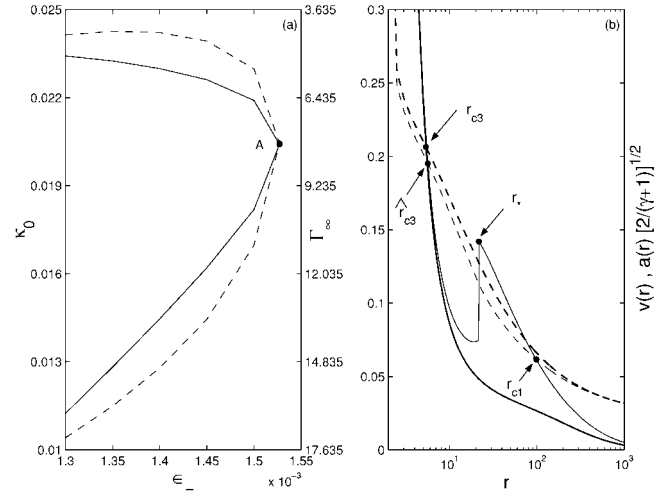


FIG. 1.—Results obtained using the M87 values for M , $\dot{M}$, and L_{jet} . (a) κ_0 (dashed line) and Γ_∞ (solid line) as functions of the upstream energy transport rate ϵ_- such that the energy conservation condition $\dot{N}_{\text{esc}} E_{\text{esc}} = L_{\text{jet}}$ is satisfied. (b) Flow velocity $v(r)$ (solid lines) and sound speed $a(r)[2/(\gamma + 1)]^{1/2}$ (dashed lines) are plotted as functions of the radius for the model A parameters, which are indicated by the letter “A” in panel a. In panel b, the thin lines represent the shocked solutions and the heavy lines denote the shock-free solutions. See the discussion in the text.

power lost from the disk at the isothermal shock location is related to observable parameters by writing

$$L_{\text{jet}} \equiv \dot{M} c^2 \Delta\epsilon, \quad (5)$$

where L_{jet} and $\dot{M}$ are the observed kinetic power and the accretion rate for a specific source, respectively. Hence we conclude that $\Delta\epsilon = L_{\text{jet}}/(\dot{M} c^2)$. For given observational values of M , $\dot{M}$, and L_{jet} , only certain combinations of the specific angular momentum ℓ and the upstream specific energy transport rate ϵ_- will yield the correct value for $\Delta\epsilon$. Hence we can view ϵ_- as the fundamental free parameter and determine ℓ using the energy constraint $\Delta\epsilon = L_{\text{jet}}/(\dot{M} c^2)$ once M , $\dot{M}$, and L_{jet} are specified for a particular source. The two sources of interest here are M87 and Sgr A*, which correspond to models A and B, respectively.

3. PARTICLE TRANSPORT EQUATION

The model for the diffusive particle transport employed here is similar to the scenario for the acceleration of cosmic rays first suggested by Blandford & Ostriker (1978), although in our situation the shock maintains a fixed location in the disk and the seed particles for the acceleration process are provided by the accretion flow itself. The energy of the injected seed particles is given by $E_0 = 0.002$ ergs, which corresponds to an injected Lorentz factor $\Gamma_0 \equiv E_0/(m_p c^2) \sim 1.3$, where m_p is the proton mass. Starting with values for ϵ_- and ℓ associated with a particular source, we can compute the injection rate for the seed particles, $\dot{N}_0$, by setting $\dot{N}_0 E_0 = \dot{M} c^2 \Delta\epsilon = L_{\text{jet}}$ (see eq. [5]), which ensures that the energy injection rate for the relativistic seed particles is equal to the energy loss rate for the background gas at the isothermal shock location. The particle injection rate is therefore given by $\dot{N}_0 = L_{\text{jet}}/E_0$, which represents one of the fundamental self-consistency requirements for the model.

The particle Green’s function $f_G(\epsilon, r)$ describes the energy and spatial distribution of the relativistic particles injected with en-

ergy E_0 at the shock location r_* . The associated number and energy densities of the relativistic particles are given by $n_r = \int_0^\infty 4\pi E^2 f_G dE$ and $U_r = \int_0^\infty 4\pi E^3 f_G dE$, respectively. Including the escape of the relativistic particles at the shock location, the energy moments $I_n \equiv \int_0^\infty 4\pi E^n f_G dE$ of the Green's function satisfy the one-dimensional steady state transport equation (T. Le & P. A. Becker 2004, in preparation)

$$\frac{d}{dr} (4\pi r H F_n) = 4\pi r H \left[\frac{(2-n)v}{3} \frac{dI_n}{dr} + \frac{\dot{N}_0 E_0^{n-2}}{4\pi r_* H_*} \times \delta(r-r_*) - c A_0 I_n \delta(r-r_*) \right], \quad (6)$$

where c is the speed of light, H_* is the disk half-thickness at the shock location, and the flux F_n is defined by

$$F_n \equiv -\frac{(n+1)v}{3} I_n - \kappa \frac{dI_n}{dr}, \quad (7)$$

which is obtained by integrating over the vertical structure of the disk. The terms on the right-hand side of equation (6) represent first-order Fermi acceleration, the particle source, and the escape of particles from the disk at the shock location, respectively. The terms on the right-hand side of equation (7) describe the diffusion and advection of particles, respectively, and κ denotes the spatial diffusion coefficient. The dimensionless parameter A_0 is a constant that represents the microphysical processes controlling the vertical escape of energetic particles from the disk in the vicinity of the shock, as discussed below.

To close the system, we must also specify the radial variation of the spatial diffusion coefficient κ in equation (7). Since the radial dependence of κ is not known precisely in this physical situation, we adopt the general form

$$\kappa(r) = \kappa_0 v(r) r_s \left(\frac{r}{r_s} - 1 \right)^\alpha, \quad (8)$$

where κ_0 is a dimensionless positive constant. Close to the event horizon, inward advection at the speed of light must dominate over outward diffusion. Conversely, in the outer region, we expect that diffusion will dominate over advection as $r \rightarrow \infty$. A careful analysis of the differential equation for the relativistic particle number density n_r obtained by setting $n = 2$ in equation (6) reveals that these conditions are both satisfied if $\alpha > 1$, and in our work we set $\alpha = 2$.

The spatial diffusion coefficient κ is related to the magnetic mean free path via the usual expression $\kappa = c\lambda_{\text{mag}}/3$. Analysis of the three-dimensional random walk of the escaping particles then yields

$$A_0 = \left(\frac{3\kappa_*}{cH_*} \right)^2, \quad (9)$$

where $\kappa_* \equiv (\kappa_- + \kappa_+)/2$. We need to specify the value of κ_0 in order to compute κ_- and κ_+ using equation (8). Since A_0 is independent of the particle energy, it follows that all of the relativistic particles have the same escape probability per unit time in the vicinity of the shock, and therefore the mean energy of the escaping particles is given by $E_{\text{esc}} = U_r/n_*$, where $n_* \equiv n_r(r_*)$ and $U_* \equiv U_r(r_*)$ denote the relativistic particle number and energy densities at the shock location. Note that E_{esc} is

proportional to E_0 and independent of $\dot{N}_0$. To obtain a self-consistent result, the product of the calculated escape energy E_{esc} and the particle escape rate N_{esc} at the shock location must be equal to the total energy lost from the background in the dynamical model, and therefore we set $N_{\text{esc}} E_{\text{esc}} = Mc^2 \Delta\epsilon = L_{\text{jet}}$. This condition is not automatically satisfied in general, and therefore it constrains the free parameters ϵ_- and κ_0 as explained below.

4. APPLICATIONS TO M87 AND SGR A*

In our numerical examples, the disk structures for M87 and Sgr A* are determined on the basis of observational estimates for the black hole mass M , the mass accretion rate $\dot{M}$, and the jet kinetic power L_{jet} . From observations of M87, we obtain $M \sim 3 \times 10^9 M_\odot$ (e.g., Ford et al. 1994), $\dot{M} \sim 1.3 \times 10^{-1} M_\odot \text{ yr}^{-1}$ (Reynolds et al. 1996), and $L_{\text{jet}} \sim 10^{43} - 10^{44} \text{ ergs s}^{-1}$ (Reynolds et al. 1996; Bicknell & Begelman 1996; Owen et al. 2000). We adopt the average value $L_{\text{jet}} = 5.5 \times 10^{43} \text{ ergs s}^{-1}$. For Sgr A*, observations indicate that $M \sim 2.6 \times 10^6 M_\odot$ (e.g., Schödel et al. 2002) and $\dot{M} \sim 8.8 \times 10^{-7} M_\odot \text{ yr}^{-1}$ (e.g., Yuan et al. 2002; Quataert 2003). The kinetic luminosity of the jet in Sgr A* is rather uncertain (see, e.g., Yuan 2000; Yuan et al. 2002). In our work, we adopt the value quoted by Falcke & Biermann (1999), $L_{\text{jet}} \sim 5 \times 10^{38} \text{ ergs s}^{-1}$.

Our goal here is to compute self-consistently the mean energy of the escaping particles, E_{esc} , while matching the observational values of M , $\dot{M}$, and L_{jet} for a specific source. We utilize natural gravitational units, with $GM = c = 1$ and $r_s = 2$, except as noted. Following Narayan et al. (1997), we assume approximate equipartition between the gas and magnetic pressures, and therefore we set $\gamma = 1.5$. Once the values of M , $\dot{M}$, and L_{jet} have been specified for the source, we select a value for the fundamental free parameter ϵ_- and then compute ℓ based on the energy conservation condition $\Delta\epsilon = L_{\text{jet}}/(Mc^2)$ (see § 2). The velocity profile $v(r)$ is determined by integrating numerically the wind equation $(1/v)(dv/dr) = N/D$, where N and D are given by equation (3). The associated solution for the isothermal sound speed $a(r)$ can then be computed using equation (2) based on the fact that K is constant in the adiabatic flow (except at the shock if there is one).

Once the velocity profile $v(r)$ has been obtained, the transport equation (6) can be solved numerically to determine the number and energy density distributions for the relativistic particles in the disk, $n_r(r) = I_2(r)$ and $U_r(r) = I_3(r)$. The energy conservation requirement $N_{\text{esc}} E_{\text{esc}} = L_{\text{jet}}$ then allows us to solve for the corresponding value of the diffusion coefficient constant κ_0 . On the basis of the observational values for M , $\dot{M}$, and L_{jet} associated with M87, in Figure 1a we plot the results obtained for κ_0 as a function of ϵ_- . In general, one obtains two roots for κ_0 if $\epsilon_- < \epsilon_{\text{max}}$, where $\epsilon_{\text{max}} = 0.001527$ for the M87 parameters and $\epsilon_{\text{max}} = 0.001229$ for the Sgr A* parameters. No self-consistent solutions exist if $\epsilon_- > \epsilon_{\text{max}}$, and a single root for κ_0 is obtained if $\epsilon_- = \epsilon_{\text{max}}$ (see Fig. 1a). For illustrative purposes, we develop two detailed models by setting $\epsilon_- = \epsilon_{\text{max}}$ for M87 and Sgr A*.

The parameter values for the two models are $\epsilon_- = 0.001527$, $\ell = 3.1340$, $\kappa_0 = 0.02044$, $\dot{N}_0 = 2.75 \times 10^{46} \text{ s}^{-1}$, $\kappa_* = 0.427877$, $A_0 = 0.0124$, $n_* = 1.19 \times 10^{44} \text{ cm}^{-3}$, $U_* = 1.42 \times 10^{42} \text{ ergs cm}^{-3}$, $N_{\text{esc}} = 4.61 \times 10^{45} \text{ s}^{-1}$, $E_{\text{esc}} = 0.0119 \text{ ergs}$, $r_* = 21.654$, $\epsilon_+ = -0.005746$, $r_{c1} = 98.524$, $r_{c3} = 5.379$, $\hat{r}_{c3} = 5.659$, $\mathcal{M}_- = 1.125$, $R_* = 1.897$, and $H_* = 11.544$ for model A (M87). For model B (Sgr A*), we have $\epsilon_- = 0.001229$, $\ell = 3.1524$, $\kappa_0 = 0.02819$, $\dot{N}_0 =$

$2.51 \times 10^{41} \text{ s}^{-1}$, $\kappa_* = 0.321414$, $A_0 = 0.0158$, $n_* = 1.93 \times 10^{39} \text{ cm}^{-3}$, $U_* = 2.10 \times 10^{37} \text{ ergs cm}^{-3}$, $N_{\text{esc}} = 4.56 \times 10^{40} \text{ s}^{-1}$, $E_{\text{esc}} = 0.0109 \text{ ergs}$, $r_* = 15.583$, $\epsilon_+ = -0.008749$, $r_{c1} = 131.874$, $r_{c3} = 5.329$, $\hat{r}_{c3} = 5.723$, $\mathcal{M}_- = 1.146$, $R_* = 1.970$, and $H_* = 7.672$. The corresponding dynamical solutions for $v(r)$ and $a(r)$ (with and without a shock) are plotted in Figure 1b for model A.

The terminal (asymptotic) Lorentz factor of the jet can be computed using $\Gamma_\infty = E_{\text{esc}}/(m_p c^2)$. In Figure 1a we plot Γ_∞ as a function of ϵ_- for the M87 parameters. Since in general we obtain two roots for κ_0 for each value of ϵ_- , we also obtain two Γ_∞ -values, with the larger one corresponding to the smaller κ_0 root. For models A (M87) and B (Sgr A*), we obtain $\Gamma_\infty = 7.92$ and 7.26 , respectively. The associated values for the mean energy boost ratio, E_{esc}/E_0 , are given by $E_{\text{esc}}/E_0 = 5.95$ and 5.45 , respectively. Conversely, we find that in the shock-free models with the same values for ϵ_- , ℓ , and κ_0 , the energy is boosted by only a factor of ~ 1.4 – 1.5 . This clearly establishes the essential role of the shock in efficiently accelerating particles up to very high energies, well above the energy required to escape from the disk. The self-consistency conditions $N_0 E_0 = L_{\text{jet}}$ and $N_{\text{esc}} E_{\text{esc}} = L_{\text{jet}}$ can be combined to conclude that $N_{\text{esc}} = N_0 E_0 / E_{\text{esc}}$, which implies that $\lesssim 20\%$ of the injected particles escape (vertically) through the surface of the disk. The remainder of the particles either are advected into the black hole or diffuse outward (horizontally) through the disk. Our predictions for the shock/jet location and the asymptotic Lorentz factor agree rather well with the observations of M87 reported by Biretta et al. (1999, 2002). The shock location predicted by the model in the case of Sgr A* is fairly close to the value suggested by Yuan (2000). However, no reliable observational estimate for the jet Lorentz factor in Sgr A* is currently available.

5. CONCLUSIONS

In this Letter, we have demonstrated for the first time the physical consequences of a standing shock in an inviscid ADAF

accretion disk for the acceleration of relativistic particles and the production of outflows. The global model presented here provides a single coherent explanation for the disk/outflow structure based on the well-understood concept of first-order Fermi acceleration in shock waves. The theory is based on an exact mathematical approach to the solution of the combined hydrodynamical and particle transport equations. Our work may help to explain the observational fact that the brightest X-ray AGNs do not possess strong outflows, whereas the sources with low X-ray luminosities but high levels of radio emission do. We suggest that the gas in the luminous X-ray sources is too dense to allow efficient Fermi acceleration of a relativistic particle population, and therefore in these systems, the gas simply heats as it crosses the shock. Conversely, in the tenuous ADAF accretion flows studied here, the relativistic particles are able to avoid thermalization because of the long collisional mean free path, resulting in the development of a significant nonthermal component in the particle distribution that powers the jets and produces the strong radio emission.

We have presented detailed results that confirm that the general properties of the jets observed in M87 and Sgr A* can be understood within the context of our disk/shock/outflow model. Although the specific numerical results discussed here are limited to the case of an inviscid disk, we believe that the inclusion of viscosity will not change the basic features of the inflow/outflow solutions obtained here because efficient first-order Fermi acceleration will also occur in viscous disks, provided they contain shocks. In addition to viscosity, future work must also address several other important considerations, such as the effects of general relativity and radiative transport on the disk/shock structure.

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