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COUNTERROTATING ACCRETION DISKS

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ABSTRACT

We consider accretion disks consisting of counterrotating gaseous components with an intervening shear layer. Configurations of this type may arise from the accretion of newly supplied counterrotating gas onto an existing corotating gas disk. For simplicity we consider the case where the gas well above the disk midplane is rotating with angular rate $+\Omega$ and that well below has the same properties but is rotating with rate $-\Omega$. Using the Shakura-Sunyaev alpha turbulence model, we find self-similar solutions where a thin (relative to the full disk thickness) equatorial layer accretes very rapidly, essentially at free-fall speed. As a result, the accretion speed is much larger than it would be for an alpha disk rotating in one direction. Counterrotating accretion disks may be a transient stage in the formation of counterrotating galaxies and in the accretion of matter onto compact objects.

Subject headings: accretion disks — galaxies: evolution — galaxies: formation — galaxies: nuclei — galaxies: spiral

1. INTRODUCTION

The widely considered models of accretion disks have gas rotating in one direction with a turbulent viscosity acting to transport angular momentum outward (Shakura & Sunyaev 1973). However, recent observations point to more complicated disk structures in both galactic nuclei and on a galactic scale. Studies of normal galaxies have revealed counterrotating gas and/or stars in many galaxies of all morphological types—ellipticals, spirals, and irregulars (see reviews by Rubin 1994a, 1994b; and Galetta 1996). In elliptical galaxies, the counterrotating component is usually in the nuclear core and may result from merging of galaxies with opposite angular momentum. Newly supplied gas with misaligned angular momentum in the nuclear region of a galaxy may be important in the formation of radio-loud quasars if there is a rotating massive black hole at the galaxy's center (Scheuer 1992).

In a number of spirals and S0 galaxies, counterrotating disks of stars and/or gas have been found to co-exist with the primary disk out to large distances (10–20 kpc), with the first example, NGC 4550, discovered by Rubin, Graham, & Kenney (1992). Of interest here is the "Evil Eye" galaxy NGC 4826, in which the direction of rotation of the gas reverses going from the inner (180 km s^{-1}) to the outer disk (-200 km s^{-1}) with an inward radial accretion speed of more than 100 km s^{-1} in the transition zone, while the stars at all radii rotate in the same direction as the gas in the inner disk which has a radius ~1200 pc (Braun et al. 1994; Rubin 1994a, 1994b). The large-scale counterrotating disks probably do not result from mergers of flat galaxies with opposite spins because of the large vertical thickening observed in simulation studies of such mergers (Barnes 1992). Thakar & Ryden (1996) discuss two different possibilities: (1) that the counterrotating matter comes from the merger of an oppositely rotating gas-rich dwarf galaxy with an existing spiral and (2) that the accretion of gas occurs over the lifetime of the galaxy with the more recently accreted gas counterrotating. Subsequent star formation in the counterrotating gas may then lead to counterrotating stars. The two-stream instability between counterrotating gas and corotating stars may enhance the rate of gas accretion (Lovelace, Jore, & Haynes 1996).

An important open problem is how counterrotating gas disks form and what their structures are on galactic scales and on the scale of disks in active galactic nuclei. Here, we investigate accretion disks consisting of counterrotating gaseous components with gas at large z rotating with angular rate $+\Omega(r)$ and gas at large negative z rotating at rate $-\Omega(r)$. The interface between the components at $z \sim 0$ constitutes a supersonic shear layer and is sketched in Figure 1. Related advection dominated accretion flows have been studied by Narayan & Yi (1995). A configuration similar to that shown in Figure 1 may arise from the accretion of newly supplied counterrotating gas onto an existing corotating disk. It might at first be supposed that powerful Kelvin-Helmholtz instabilities heat the gas to escape speed and rapidly destroy the assumed configuration. However, supersonic shear layers exist and exhibit gross stability in stellar and extra-galactic jets (Hardee, Cooper, & Clarke 1994). In the counterrotating disk, matter approaching the equatorial plane from above and below has angular momenta of opposite signs with the result that there is angular momentum annihilation at $z = 0$, the matter loses its centrifugal support and accretes at essentially free-fall speed. On the other hand, accretion disks rotating in one direction are modeled assuming a turbulent viscosity which is crucial for the outward transport of angular momentum (Shakura & Sunyaev 1973). The counterrotating disks can also be expected to be turbulent owing in part to the Kelvin-Helmholtz instability, and turbulent viscosity can transport angular momentum outward in the large $|z|$ regions of the disk.

2. THEORY AND SELF-SIMILAR SOLUTIONS

Here, we consider a model of a stationary viscous disk consisting of co- and counterrotating gas. For axisymmetric disklike flows in cylindrical coordinates $(r,\ \phi,\ z)$ with $\boldsymbol{v} = [v_r(r,z), v_\phi(r,z), v_z(r,z)]$, the momentum and continuity equations are

$$\rho(\boldsymbol{v}\cdot\boldsymbol{\nabla})v_r = \frac{\rho v_\phi^2}{r} - \frac{\partial\tilde{p}}{\partial r} + \rho g_r + \frac{2}{r}\frac{\partial}{\partial r}\left(r\eta\frac{\partial v_r}{\partial r}\right) + \frac{\partial}{\partial z}\left[\eta\left(\frac{\partial v_r}{\partial z} + \frac{\partial v_z}{\partial r}\right)\right] - \frac{2\eta v_r}{r^2}, \qquad (1)$$

$$\rho(\boldsymbol{v}\cdot\boldsymbol{\nabla})v_\phi = \frac{-\rho v_r v_\phi}{r} + \frac{1}{r^2}\frac{\partial}{\partial r}\left[r^3\eta\frac{\partial(v_\phi/r)}{\partial r}\right] + \frac{\partial}{\partial z}\left(\eta\frac{\partial v_\phi}{\partial z}\right), \qquad (2)$$

$$\rho(\boldsymbol{v}\cdot\boldsymbol{\nabla})v_z = -\frac{\partial\tilde{p}}{\partial z} + \rho g_z + \frac{1}{r}\frac{\partial}{\partial r}\left[r\eta\left(\frac{\partial v_r}{\partial z} + \frac{\partial v_z}{\partial r}\right)\right] + 2\frac{\partial}{\partial z}\left(\eta\frac{\partial v_z}{\partial z}\right). \qquad (3)$$

$$\frac{1}{r}\frac{\partial}{\partial r}(r\rho v_r) + \frac{\partial}{\partial z}(\rho v_z) = 0. \qquad (4)$$

Here, $\rho(r,z)$ is the gas density, $\tilde{p}(r,z) \equiv p + \frac{2}{3}\eta(\boldsymbol{\nabla}\cdot\boldsymbol{v})$ with p the pressure, $\boldsymbol{g} = -\boldsymbol{\nabla}\Phi$ the gravitational acceleration with $\Phi(r,\ z)$ the potential, $\eta = \rho\nu$ the dynamic viscosity, and $\nu = \eta/\rho$ the kinematic viscosity. Because the microscopic viscosity is negligible for the conditions considered, we assume ν arises from small-scale turbulence and can be approximated by the "alpha" prescription of Shakura (1973) and Shakura & Sunyaev (1973) as $\nu = \alpha c_s H$, where H is the full half-thickness of the disk, which is assumed to be thin ($H \ll r$); c_s the sound speed; and α a dimensionless constant much less than unity. We also assume the second viscosity is zero so that the stress tensor $T_{ij} = \rho v_i v_j + p\delta_{ij} + T_{ij}^\nu$ with the viscous contribution $T_{ij}^\nu = -\rho\nu[\partial v_i/\partial x_j + \partial v_j/\partial x_i - \frac{2}{3}(\boldsymbol{\nabla}\cdot\boldsymbol{v})\delta_{ij}]$.

We seek self-similar solutions to equations (1)–(4) of the form

$$v_r(r,z) = -u_r(\zeta)V_c(r), \quad v_\phi(r,z) = u_\phi(\zeta)V_c(r),$$
$$v_z(r,z) = -\left[\frac{h(r)}{r}\right]u_z(\zeta)V_c(r),$$
$$\rho(r,z) = \rho_0(r)\,Z\left[\frac{z}{H(r)}\right], \qquad (5)$$

where $\zeta \equiv z/h(r)$ is the dimensionless vertical distance in the disk with $h(r)$ a length scale identified subsequently, and $V_c(r) = [r(\partial\Phi/\partial r)|_{z=0}]^{1/2}$ is the circular velocity of the gas. The potential is due in general to disk and halo matter and a central object. Here and subsequently we neglect the pressure force in the radial equation of motion. We consider cases where $V_c(r) = (GM/r_0)^{1/2}(r_0/r)^n$, with r_0 a constant length scale. The value $n = 1/2$ corresponds to a Keplerian disk around an object of mass M, and $n = 0$ to a flat rotation curve applicable to flat galaxies. We assume that the disk is not strongly self-gravitating and that it is vertically isothermal so that its overall half-thickness is $H \sim (c_s/V_c)r$, which is the same as for a disk rotating in one direction. Thus, for the "alpha" viscosity, $\eta = \alpha p/\Omega_c$, where $\Omega_c \equiv V_c/r$.

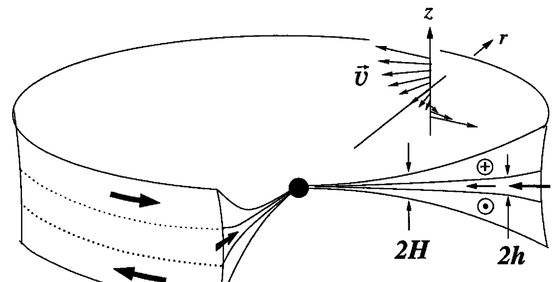

FIG. 1.—Structure of two apposed, counterrotating accretion disks and the midplane boundary layer. The inset shows the three dimensional view of the velocity field for the $n = 1/2$ case shown in Fig. 2. The velocity variation is analogous to that in the Ekman layer of a rotating fluid (such as the ocean) where the Coriolis force balances the viscous force (see, for example, Batchelor 1967, p. 197).

Substitution of equations (5) into equations (1) and (2) gives

$$u_r'' = u_\phi^2 + nu_r^2 - 1 + k_h\zeta u_r u_r' - u_z u_r' + 2\delta^2\zeta u_r' + O(\epsilon^2), \qquad (6)$$

$$u_\phi'' = (n-1)u_r u_\phi - u_z u_\phi' + k_h\zeta u_r u_\phi' + 2\delta^2\zeta u_\phi' + O(\epsilon^2). \qquad (7)$$

Here, a prime denotes a derivative with respect to ζ, $k_h \equiv (r/h)(dh/dr)$, $\epsilon \equiv h/r \ll 1$, and $\delta \equiv h/H$. We have made the natural identification $h^2 = \nu r/V_c$, which implies $\delta = \sqrt{\alpha}$. We consider the ordering

$$\epsilon^2 \ll \delta^2 \ll 1.$$

For $\delta^2 \ll 1$, equation (3) simplifies to $\partial_z p = -\rho\partial_z\Phi$, which determines the isothermal density profile $Z(\zeta) = \exp(-\delta^2\zeta^2)$. From the assumed scaling $\rho_0(r) \propto r^{-\beta-1/2}$, equation (4),

$$u_z' = \left[\beta + n - \frac{1}{2} - 2k_H\delta^2\zeta^2\right]u_r + k_h\zeta u_r' + 2\delta^2\zeta u_z, \qquad (8)$$

determines v_z, where $k_H \equiv (r/H)(dH/dr)$. For consistency of the self-similar solutions, δ and k_h must be independent of r, which implies $k_H = \text{constant}$.

Equations (6), (7), and (8) constitute a closed system for the dimensionless functions $u_r(\zeta)$, $u_\phi(\zeta)$, and $u_z(\zeta)$. Applying equations (6), (7), and (8) to the flow suggested in Figure 1, we infer that $u_r(\zeta)$ is an even function of ζ, while $u_\phi(\zeta)$ and $u_z(\zeta)$ are odd functions. For $|\zeta| \gg 1$, we impose $u_r \to 0$ and $u_\phi \to \pm 1$ so as to have $u_r'' \to 0$ and $u_\phi'' \to 0$. In a more complete treatment with the $O(\epsilon^2)$ terms retained in equations (6) and (7), the solution at large $|\zeta|$ approaches that of Shakura & Sunyaev (1973) with $u_r = O(\epsilon^2)$.

A steady counterrotating disk may result from gas continuously supplied at large r. Similarly, steady state solutions where gas is entrained and delivered vertically onto the faces of the disk may exist. We assume that $\rho v_z \to 0$ as $|z| \to \infty$.

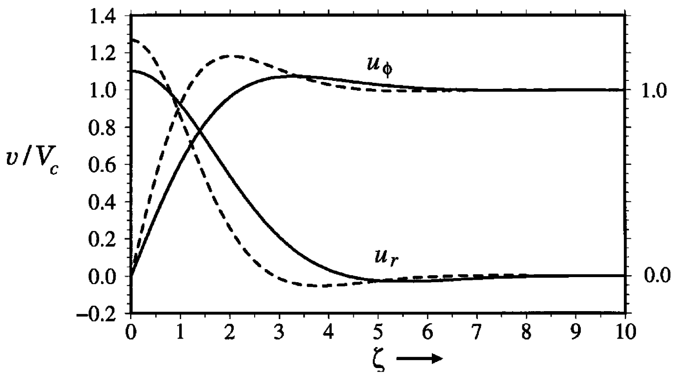

FIG. 2.—Solutions to eqs. (6), (7), and (8) for $\delta = 0.1$ as a function of $\zeta = z/h(r)$. The solid lines correspond to a Keplerian disk ($n = 1/2$, $\beta = 1$) with $u_r(0) = 1.10011267$ and $u'_\phi(0) = 0.66265042$. The dashed curves correspond to $n = 0$ (flat rotation curve) and $\beta = 3/2$ and have $u_r(0) = 1.26728062$ and $u'_\phi(0) = 1.12573559$. These values are close to those for the $\delta = 0$ "zero-order" solutions, $u_r(0)/2 - u_r(0)^3/12 = [u'_\phi(0)]^2 = 1.10011267$ for ($n = 1/2$), and $u_r(0) = [u'_\phi(0)]^2 = 1.26728062$ for ($n = 0$). All solutions reach their asymptotic values for $\zeta \to \infty$ and the vertical velocity $v_z/V_c \approx 0$ on this scale. Note that the midplane infall velocity for the Keplerian case is shear-reduced by a factor ~ 0.77 from the free-fall speed $(2GM/r)^{1/2}$.

Considering only this subset of solutions, we demand the accretion rate

$$\dot{M} = 2\pi \int_{-\infty}^{\infty} dz\, r\rho(r, z)|v_r(r, z)|$$

$$\simeq 2\pi r \rho_0(r) V_c(r) h(r) \int_{-H/h}^{H/h} d\zeta\, u_r(\zeta) e^{-\delta^2 \zeta^2} , \qquad (9)$$

be a constant. A sufficient condition for this is $\beta = k_h - n + 1/2$ and $\delta = h/H$ is r-independent. For an alpha disk rotating in one direction, $\dot{M}_{SS} = 2\pi r \Sigma |v_r|_{SS}$, with the accretion speed $|v_r|_{SS} \sim \alpha c_s H/r \ll c_s$, where $\Sigma = \int dz\, \rho$ is the disk's surface mass density (Shakura & Sunyaev 1973). In contrast, for a counterrotating disk equation (9) gives $\dot{M}_{CR} = 2\pi r \Sigma |v_r|_{CR}$, with accretion speed (averaged over z) $|v_r|_{CR} \sim (h/H) V_c$. For the same α, we have $|v_r|_{CR} \gg |v_r|_{SS}$: the accretion speed of the counterrotating disk is much larger than that of the disk rotating in one direction.

Consider now the energy dissipation and radiation of the disk. The viscous dissipation per unit area of the disk is

$$D(r) = \int_{-H}^{H} dz\, \rho\nu \left[2\left(\frac{\partial v_r}{\partial r}\right)^2 + 2\left(\frac{v_r}{r}\right)^2 + 2\left(\frac{\partial v_z}{\partial z}\right)^2 + \left(r \frac{\partial \Omega}{\partial r}\right)^2 + \left(\frac{\partial v_z}{\partial r} + \frac{\partial v_r}{\partial z}\right)^2 + \left(\frac{\partial v_\phi}{\partial z}\right)^2 - \frac{2}{3} (\nabla \cdot \boldsymbol{v})^2 \right], \qquad (10)$$

where $\Omega = v_\phi / r$. For a disk rotating in one direction, the dominant contribution to $D(r)$ is from the $\partial\Omega/\partial r$ term and this gives $D_{SS}(r) \approx \Sigma\nu(n + 1)^2 (V_c/r)^2 = \dot{M}_{SS}(V_c/r)^2 (n + 1)^2/(3\pi)$. In contrast, for a counterrotating disk, the dominant contributions are from the terms $\partial v_r/\partial z$ and $\partial v_\phi/\partial z$ with the result that $D_{CR}(r) \sim \Sigma\nu V_c^2/(hH) \sim \dot{M}_{CR}(V_c/r)^2/(2\pi)$. The dissipated energy is radiated from the faces of the disk if it is optically thick in the z-direction and thus $D = 4acT^4/(3\kappa\Sigma)$, where a is the radiation constant, c the speed of light, κ the opacity, and T is the internal disk temperature. For the general scaling $\kappa \propto \rho^a T^b$, the counterrotating disk thus has $T \propto r^{-\xi}$ with $\xi = (3 + 2a + n)/(4 - b + a/2)$ and $h \propto T^{1/2} r^{1+n} \propto r^{k_h}$. For example, for Kramer's opacity $a = 1$, $b = -3.5$, we have $\xi = (5 + n)/8$ and $k_h = (11 + 15n)/16$, whereas for electrons scattering opacity $a = 0$, $b = 0$, and $\xi = (3 + n)/4$ and $k_h = (11 + 9n)/8$. The apparent surface temperature of the disk $T_{\rm eff}(r)$ is given by $D(r) = 2\sigma T_{\rm eff}^4$, where $\sigma = ac/4$ is the Stefan-Boltzmann constant. The dependence for the counterrotating disk, $T_{\rm eff} \propto r^{-(n+1)/2}$, is the same as for an alpha disk rotating in one direction. Similarly, the spectrum, F_ω, obtained by integrating the Planck function $B_\omega[T_{\rm eff}(r)]$ over the surface area of the disk, $F_\omega \propto \omega^{(3n-1)/(n+1)}$, is the same as for a disk rotating in one direction. For an optically thin disk, $D = 2H\Lambda$, where $\Lambda(\rho, T)$ is the emissivity of the disk matter.

In the limit $\delta = 0$ so that $\beta = k_h - n + 1/2$, equation (8) gives $u_z = k_h \zeta u_r$, which reduces equations (6) and (7) to $u''_r = u_\phi^2 + n u_r^2 - 1$ and $u''_\phi = (n - 1) u_r u_\phi$, respectively. These have the integral

$$(u'_\phi)^2 - \frac{1 - n}{2} (u'_r)^2 + (1 - n) u_r u_\phi^2 + \frac{n(1 - n)}{3} u_r^3 - (1 - n) u_r = \text{const}. \qquad (11)$$

Evaluating equation (11) at $\zeta = 0$ (where $u'_r = 0$ and $u_\phi = 0$) and $\zeta \to \infty$ (where $u_r = 0$ and $u_\phi = 1$) gives $[u'_\phi(0)]^2 = (1 - n) u_r(0) - n(1 - n) u_r^3(0)/3$. We utilize this result by first finding the single independent initial condition $u_r(0)$ for $\delta = 0$, then allow $\delta \neq 0$, and perturbatively find the proper values of $u_r(0)$ and $u'_\phi(0)$.

Figure 2 shows solutions to equations (6), (7), and (8) for $\delta = h/H = 0.1$. Both Keplerian ($n = 1/2$, $\beta = 1$) and galactic ($n = 0$, $\beta = 3/2$) solutions are shown. Note that for $3 \lesssim |\zeta| \lesssim 7$, the solutions have a region with $u_r < 0$ indicating spiraling outflows on both sides of the midplane. For large ζ, the density fall off gives $\rho v_z \to 0$. The solutions remain qualitatively similar for small δ *and* small, non-zero ϵ.

3. DISCUSSION

Counterrotating gas supplied to the outer part of an existing corotating gas disk accretes with a much higher speed than for gas rotating in one direction. Comparing accretion rates of a standard α-disk ($\dot{M}_{SS}$) with that of a counterrotating disk ($\dot{M}_{CR}$) of the same Σ, we find $\dot{M}_{CR}/\dot{M}_{SS} \sim \delta/\epsilon^2 = (r/H)^2/\sqrt{\alpha}$, which is much larger than unity.

Thermal and/or dynamical instabilities may destroy the counterrotating accretion flows described above. The relative importance of thermal and dynamical instabilities can be estimated by comparing the thermal dissipative timescale, $\tau_Q \equiv \Sigma c_s^2/D_{CR}(r) \sim (\epsilon^2/\delta^3)(H/c_s)$ with dynamical and viscous timescales, $\tau_{\rm dyn} \equiv H/c_s \sim r/V_c$ and $\tau_\nu \equiv r^2/\nu \sim \epsilon^{-2}(H/c_s)$, respectively. Thus for $\delta^3 > \epsilon^2 \ll 1$, we have $\tau_Q < \tau_{\rm dyn} \ll \tau_\nu$, and thermal equilibrium is maintained on timescales of possible flow instabilities. The Kelvin-Helmholtz instabilities are likely to be the most important (Ray & Ershkovich 1983; & Choudhury & Lovelace 1984). If the counterrotating disk is treated as a vortex sheet without radial inflow, then a local stability analysis [$\boldsymbol{k}^2 H^2 \geq 1$ with $\boldsymbol{k}$ the (r, ϕ) wavevector] indicates unstable warping for wavenumbers $|k_\phi/k_r| < \sqrt{2}(c_s/\Omega r) \ll 1$ and a maximum growth rate of $|k_r| c_s/2$. Instability of short wavelengths, $\lambda_r \ll 2\pi H$, will lead to a mixing of the $\pm\Omega$ components in a layer $\Delta z \sim 1/|k_r| \ll H$ and rapid accretion of this layer. Thus, these perturbations will be rapidly advected inward. Instability of the longer wavelengths, $\lambda_r \sim 2\pi H$, have a growth rate $\omega_i \sim 1/(2\tau_{\rm dyn})$. These perturbations may destroy the assumed configuration, or they may also contribute to turbulence that enhances the boundary layer accretion.

Accretion of counterrotating gas by an existing corotating gas disk may be a transient stage in the formation of counterrotating galaxies and in the accretion of matter onto rotating black holes in active galactic nuclei. We find that newly supplied counterrotating gas drags inward the old corotating gas with an equal mass of old and new gas accreting rapidly. Thus the old corotating gas may be entirely "used up" (dragged to the center of a galaxy or into a black hole) if the mass of newly supplied gas exceeds that of the old gas disk. Accretion onto the faces of an existing disk may not have the symmetry shown in Figures 1 and 2. There may instead be two layers of rapid radial inflow bounding an older existing gas disk near the midplane as sketched in Figure 3. Solutions for this configuration can be composed from those of Figure 2 if $h \ll H$.

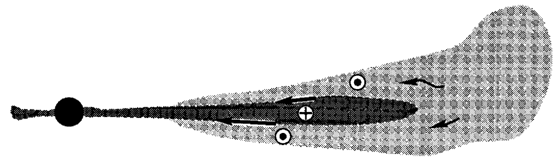

FIG. 3.—Schematic drawing of accretion of newly supplied counterrotating gas ($\odot$) induced by viscous interaction with an existing disk of corotating gas ($\oplus$).

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