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IMPLOSIVE ACCRETION AND OUTBURSTS OF ACTIVE GALACTIC NUCLEI

R. V. E. LOVELACE[1]
Institute of Astronomy, Madingley Road, Cambridge CB3 0HA, England, UK
M. M. ROMANOVA
Space Research Institute, Russian Academy of Sciences, Moscow, Russia
AND
W. I. NEWMAN
Departments of Earth and Space Sciences, Astronomy, and Mathematics, University of California, Los Angeles, CA 90024
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ABSTRACT

A model and simulation code have been developed for time-dependent axisymmetric disk accretion onto a compact object including for the first time the influence of an ordered magnetic field. The accretion rate and radiative luminosity of the disk are naturally coupled to the rate of outflow of energy and angular momentum in magnetically driven ($\pm z$) winds. The magnetic field of the wind is treated in a phenomenological way suggested by self-consistent wind solutions. The radial accretion speed $u(r, t)$ of the disk matter is shown to be the sum of the usual viscous contribution and a magnetic contribution $\propto r^{3/2}B_p^2/\sigma$, where $B_p(r, t)$ is the poloidal field threading the disk and $\sigma(r, t)$ is the disk's surface mass density. An enhancement or variation in B_p at a large radial distance leads to the formation of a soliton-like structure in the disk density, temperature, and B-field which propagates implosively inward. The implosion gives a burst in the power output in winds or jets and a simultaneous burst in the disk radiation. The model is pertinent to the formation of discrete fast-moving components in jets observed by very long baseline interferometry. These components appear to originate at times of optical outbursts of the active galactic nucleus.

Subject headings: accretion, accretion disks — galaxies: active — galaxies: jets — galaxies: nuclei

1. INTRODUCTION

High-resolution very long baseline interferometer (VLBI) radio observations of the brightest quasars has shown that the jets (0.1–10 pc) typically consist of one or more bright components moving with apparent "superluminal" speed with respect to the nuclear component. The extrapolated zero-point time of their formation often coincides with the brightness amplification of the continuum radiation in the infrared, optical, and X-ray wavebands. This correlation was noticed for several well-studied quasars, such as 3C 273, 3C 345, and 3C 120 (Babadzhanyants & Belokon 1986; Belokon 1988, 1991; Bregman et al. 1986; Zensus et al. 1990). VLBI monitoring of the quasar 3C 273 after an optical/infrared flux density outburst in 1988 revealed the appearance of a new component in the jet at the predicted time (Krichbaum et al. 1990). The data appear to confirm the earlier hypothesis of Kinman (1977) that both observed features, jet formation and flux-density outburst occur simultaneously and reflect the same physical process.

Instabilities of nonmagnetized alpha-disks (Shakura & Sunyaev 1973) have been studied intensively as models for dwarf-nova eruptions (see, e.g., Meyer & Meyer-Hofmeister 1982; Faulkner, Lin, & Papaloizou 1983; Cannizzo & Wheeler 1984). The nonmonotonic temperature dependence of the disk opacity can lead to a thermal runaway which rapidly increases the accretion speed and gives an optical outburst. An analogous mechanism in the disks of T Tauri stars has been proposed to explain the outbursts of FU Orionis objects (Hartmann & Kenyon 1985; Clark, Lin, & Pringle 1990; and Bell & Lin 1993). Flux-density outbursts of active galactic nuclei have been modeled as thermal instabilities of disks by Mineshige & Shields (1990), but jet formation was not included.

In the present work, we develop a theory for outbursts of disks threaded by an ordered magnetic field. The twisting of this field acts to drive winds or jets from the disk surfaces which carry away angular momentum and energy and thereby increases the accretion speed. The increased accretion speed can in turn amplify the B field and lead to runaway or implosive accretion and explosive wind or jet formation. We develop a model and present sample simulation results for time-dependent axisymmetric accretion disks including the influence of an ordered magnetic field. This work draws on the study by Chen & Taam (1992) of time-dependent nonmagnetic α-disks, and on the work of Lovelace, Wang, & Sulkanen (1987, hereafter LWS) on magnetized disks that of Lovelace, Berk, & Contopoulos (1991, hereafter LBC), and Lovelace, Romanova, & Contopoulos (1993, hereafter LRC) on magnetically driven winds and jets. However, the present work does not include a self-consistent solution for the wind magnetic field. Instead, we parameterize this field so as to obtain approximate accord with wind models.

2. THEORY

The basic equations are

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \boldsymbol{v}) = 0\,, \tag{1a}$$

$$\nabla \times \boldsymbol{B} = \frac{4\pi}{c}\boldsymbol{J}\,, \tag{1b}$$

$$\boldsymbol{J} = \sigma_e(\boldsymbol{E} + \boldsymbol{v} \times \boldsymbol{B}/c)\,, \tag{1c}$$

[1] On leave from Department of Applied Physics, Cornell University, Ithaca, NY 14853.

$$\nabla \times \boldsymbol{E} = -\frac{1}{c}\frac{\partial \boldsymbol{B}}{\partial t}, \tag{1d}$$

$$\rho \frac{d\boldsymbol{v}}{dt} = -\nabla p + \rho \boldsymbol{g} + \frac{1}{c}\boldsymbol{J} \times \boldsymbol{B} + \boldsymbol{F}^{\rm vis}. \tag{1e}$$

Here $\boldsymbol{v}$ is the flow velocity in the disk, ρ is the density, σ_e is the effective electrical conductivity, $\boldsymbol{F}^{\rm vis}$ is the viscous force density, $p = \rho k_{\rm B} T/\mu + aT^4/3$ is the total, gas plus radiation, pressure (with $k_{\rm B}$ the Boltzmann constant, μ the mean particle mass, and a the usual radiation constant), and $\boldsymbol{g}$ is the gravitatonal acceleration. Outside of the disk, dissipative effects are considered to be negligible ($\sigma_e \to \infty$, $\boldsymbol{F}^{\rm vis} = 0$, etc.). We assume a low-mass disk and neglect general relativistic effects so that $\boldsymbol{g} = -\nabla\Phi_g$ with $\Phi_g = -GM/(r^2 + z^2)^{1/2}$, where M is the mass of the central object. Equations (1) are supplemented later by an equation for the conservation of energy in the disk. Even field symmetry (LWS) is assumed so that $B_r(r, z, t) = -B_r(r, -z, t)$, $B_\phi(r, z, t) = -B_\phi(r, -z, t)$, and $B_z(r, z, t) = +B_z(r, -z, t)$. Previous treatments of disks with ordered B fields have assumed stationary solutions (see, however, Newman, Newman, & Lovelace 1992).

The plasma flow in a thin disk, can be described approximately by

$$\boldsymbol{v} = -u(r, t)\hat{\boldsymbol{r}} + v_\phi(r, t)\hat{\boldsymbol{\phi}} + \left(\frac{z}{h}\right)\left(\frac{dh}{dt}\right)\hat{\boldsymbol{z}},$$

where u is the accretion speed, $2h(r, t) \ll r$ is the disk thickness, and $d/dt = \partial/\partial t - u\,\partial/\partial r$ is the convective time derivative following a ring of disk matter. Integration of the continuity equation (1a) over the vertical (z) thickness of the disk gives

$$\frac{\partial}{\partial t}(r\sigma) = \frac{\partial}{\partial r}[(r\sigma)u], \tag{2}$$

where $\sigma(r, t) = \int_{-\infty}^{\infty} dz\rho(r, z, t)$ is the disk's surface mass density. Equation (2) neglects the small fractional mass loss due to outflows from the surfaces of the disk ($z = \pm h$) (see LBC).

Integration of the r-component of the Navier-Stokes equation (1e) over z gives

$$\sigma\frac{du}{dt} = \frac{\partial}{\partial r}P + \frac{\sigma(v_{\rm K}^2 - v_\phi^2)}{r} - \int dz\, F_r^{\rm mag} - \int dz\, F_r^{\rm vis},$$

or

$$\frac{\partial}{\partial t}(r\sigma u) = \frac{\partial}{\partial r}\left[r\sigma\left(u^2 + \frac{P}{\sigma}\right)\right] + \sigma\left[v_K^2 - v_\phi^2 - \left(\frac{P}{\sigma}\right)\right] - r\int dz\, F_r^{\rm mag} - r\int dz\, F_r^{\rm vis}, \tag{3}$$

where $P = \sigma k_{\rm B} T/\mu + 2haT^4/3$ is the height integrated total pressure, $T(r, t)$ is the midplane temperature of the disk, and $v_{\rm K} \equiv (GM/r)^{1/2}$ is the Keplerian velocity at r. Assuming even field symmetry, the radial magnetic force can be written as

$$\int dz\, F_r^{\rm mag} = \frac{1}{2\pi}(B_r B_z)_h - \frac{1}{r^2}\frac{\partial}{\partial r}\left[\frac{2hr^2\langle B_\phi^2 - B_r^2\rangle}{8\pi}\right] - \frac{\partial}{\partial r}\left[\frac{2h\langle B_z^2\rangle}{8\pi}\right] + \frac{1}{4\pi}\left(\frac{\partial h}{\partial r}\right)(B_\phi^2 + B_z^2 - B_r^2)_h, \tag{4}$$

where $\langle\cdots\rangle \equiv \int_{-h}^{h} dz(\cdots)/(2h)$, and the h subscript indicates that the quantity is evaluated at $z = h$. The dominant term, $(B_r B_z)_h/2\pi$, represents the radial force on the disk due to the external magnetic field.

The viscous force in equations (1e) and (3) is assumed to arise from an isotropic turbulent viscosity ν_t with Stokes's hypothesis regarding the second viscosity coefficient (Schlichting 1968); that is, the momentum flux density tensor due to the turbulence has zero trace. Thus,

$$\int dz\, F_r^{\rm vis} = -\frac{4}{3}\frac{\partial}{\partial r}\left[\frac{\sigma\nu_t}{r}\frac{\partial}{\partial r}(ru)\right] + \frac{2u}{r}\frac{\partial}{\partial r}(\sigma\nu_t). \tag{5}$$

The magnitude of the turbulent viscosity is assumed to be given by an α-model (Shakura 1973 and Shakura & Sunyaev 1973) based on the gas pressure of the disk; namely, $\nu_t = (2/3)\alpha(p_{\rm gas}/\rho)(r/v_{\rm K})$, where α is a dimensionless constant less than or of order unity. Sakimoto & Coroniti (1981) give physical arguments for this form of the viscosity law for AGN disks and show that for this law the inner part of the disk is thermally stable under typical conditions where radiation pressure is dominant. The possible contribution to the turbulent momentum flux due to small-scale magnetic field fluctuations is assumed included in α (Eardley & Lightman 1975; Coroniti 1981; Balbus & Hawley 1992; Kaisig, Tajima, & Lovelace 1992).

Integration over z of the ϕ component of the Navier-Stokes equation (1e) gives

$$\frac{\partial}{\partial t}(r\sigma l) = \frac{\partial}{\partial r}\left[(r\sigma l)u + r^3\sigma\nu_t\frac{d\omega}{dr}\right] + \frac{1}{2\pi}(r^2 B_\phi B_z)_h + \frac{\partial}{\partial r}\left[\frac{hr^2\langle B_r B_\phi\rangle}{2\pi}\right] - \frac{1}{2\pi}\left(\frac{\partial h}{\partial r}\right)(r^2 B_r B_\phi)_h, \tag{6}$$

where $l \equiv rv_\phi$ is the specific angular momentum, and $\omega \equiv v_\phi/r$. The dominant magnetic term, $(1/2\pi)(r^2 B_\phi B_z)_h$, represents the torque on the disk due to the external magnetic field associated with the outflows from the disk's surfaces.

The z-component of the Navier-Stokes equation (1e) gives the condition for vertical hydrostatic balance. This can be written as

$$\left(\frac{h}{r}\right)^2 + b\left(\frac{h}{r}\right) - \frac{2p}{\rho v_{\rm K}^2} = 0, \tag{7}$$

where p is the total, gas plus radiation pressure, and $b \equiv r[(B_\phi)_h^2 + (B_r)_h^2]/(4\pi\sigma v_{\rm K}^2)$ (Wang, Lovelace, & Sulkanen 1990). For $b^2 \ll 8p/(\rho v_{\rm K}^2)$, this equation gives the well-known relation $h/r = (2p/\rho v_{\rm K}^2)^{1/2}$ (Shakura & Sunyaev 1973), while for $b^2 \gg 8p/(\rho v_{\rm K}^2)$ it gives $h/r = 2p/(\rho v_{\rm K}^2 b)$ which is smaller than $(2p/\rho v_{\rm K}^2)^{1/2}$ owing to the compressive effect of the magnetic field external to the disk (Wang et al. 1990).

The solution for the magnetic field $\boldsymbol{B}$ outside of the disk ($|z| > h$) is matched to the field solution inside the disk at $z = \pm h$ as discussed by LWS. We first treat the internal field solution which is described by equations (1b), (1c), and (1d). These can be combined to give the "induction equation,"

$$\frac{\partial \boldsymbol{B}}{\partial t} = \nabla \times (\boldsymbol{v} \times \boldsymbol{B}) - \nabla \times (\eta\nabla \times \boldsymbol{B}), \tag{8}$$

where the magnetic diffusivity $\eta \equiv c^2/(4\pi\sigma_e)$ has the same units as kinematic viscosity. We assume that the magnitude of η is comparable to that of the turbulent viscosity ν_t (Bisnovatyi-Kogan & Ruzmaikin 1976; Parker 1979); that is, we let $\eta = D\nu_t$ with $D = O(1)$.

For the even field symmetry assumed, B_z is an even function of z. Because $\nabla \cdot \boldsymbol{B} = 0$, this implies $\Delta B_r/r \approx \Delta B_z/h$. Then, the variation of B_z from $z = 0$ to $z = h$ is $\Delta B_z \approx (h/r)(B_r)_h$. Assuming $(B_r)_h \lesssim B_z$ as discussed below, it follows that $\Delta B_z \ll B_z$; that is, the variation of B_z with z within the disk is negligible. The z-component of equation (8) is

$$\frac{\partial B_z}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r}(urB_z) + \frac{1}{r}\frac{\partial}{\partial r}\left[\eta r\left(\frac{\partial B_z}{\partial r} - \frac{\partial B_r}{\partial z}\right)\right].$$

For even field symmetry, B_r is an odd function of z and consequently $\partial B_r/\partial z \approx (B_r)_h/h$. The condition for validity of this equality is $h^2 |\partial^3 B_r/\partial z^3| \ll |\partial B_r/\partial z|$. Thus,

$$\frac{\partial}{\partial t}(rB_z) = \frac{\partial}{\partial r}\left[(rB_z)u - \frac{\eta r(B_r)_h}{h} + \eta r\frac{\partial B_z}{\partial r}\right]. \quad (9)$$

Inside the brackets, the term proportional to u describes the inward advection of the poloidal field with the accretion flow, the term proportional to $\eta(B_r)_h$ describes the diffusive outward drift of the poloidal field, and the term proportional to $\eta(\partial B_z/\partial r)$ represents the radial diffusion of the poloidal field. Equation (9) is of central importance to the transport of a large-scale B-field in an accretion disk. An alternative derivation of equation (9), obtained by integrating the equation preceding it from $z = 0$ to $z = h$, gives the same result apart from small terms $[O(\partial h/\partial r)]$.

In contrast with equation (10) fo $B_z(r, t)$, the values of $(B_r)_h$ and $(B_\phi)_h$ are determined by the field solution external to the disk. External MHD outflow solutions are discussed by Blandford & Payne (1982) and LBC. The timescale for setting up or changing the outflow is the Alfvén propagation time over distances of order r in the low-density external plasma. Therefore, this timescale is much shorter than any timescales for changes in the disk. The external field solutions for magnetically driven outflows obey the relation

$$(B_r)_h = \beta_r B_z(r, t), \quad (10)$$

where β_r is a constant of order unity. Blandford & Payne give the condition $\beta_r \gtrsim 1/3^{1/2} \approx 0.577$ for MHD outflows. On the other hand, the MHD solutions of LBC and LRC give $\beta_r \approx 1$ for outflows. The fact that there is definite vaue of β_r for MHD outflows can be understood qualitatively by considering the net force on a fluid particle in the direction of its poloidal motion away from the disk surface. For the conditions of interest here, this force must include the contribution of radiation pressure with the result that the net force has the form $F \propto (h/r) - (1 - 3\beta_r^2)(z/r)$ for $z/r \ll 1$ and $1 - \omega^2/\omega_K^2 \ll h/r$. The first term is due to the radiation pressure, while the terms $\propto z$ result from the centrifugal force and the gravitational force of the central object. The slow magnetosonic point of the outflow occurs where $F = 0$ (LRC) at $z_s \sim h/(1 - 3\beta_r^2)$. In order for there to be only a small fractional mass outflow from the disk as assumed earlier, z_s must be larger than h. For example, for $z_s = 3h$, $\beta_r \approx 0.47$. Of course, the actual value(s) of β_r in MHD outflows requires the self-consistent solution of the MHD equations outside the disk. However, the model and numerical results presented here do not depend sensitively on β_r. Therefore, we adopt $\beta_r = 1$. Consistent with the detailed models (Blandford & Payne 1982 and LBC), we assume $[B_\phi(r, t)]_h = -\beta_\phi B_z(r, t)$ with β_ϕ is a constant of order unity. We adopt $\beta_\phi = 1$.

Conservation of energy in the disk can be written as an equation for the midplane temperature of the disk $T(r, t)$,

$$c_v \sigma \frac{dT}{dt} = c_v(\gamma - 1)Th\frac{d}{dt}\left(\frac{\sigma}{h}\right) + D_{v\phi} + D_{vr} + D_O - 2F_R - \frac{1}{r}\frac{\partial(rh\mathscr{F}_r)}{\partial r}. \quad (11a)$$

Here $c_v \equiv k_B/[\mu(\gamma - 1)]$ is the specific heat; γ is the specific heat ratio;

$$D_{v\phi} = \sigma v_t\left[r\left(\frac{\partial\omega}{\partial r}\right)\right]^2 \quad (11b)$$

is the viscous dissipation (per unit area of the disk) due to the shear in the azimuthal motion; D_{vr} is the viscous dissipation due to the radial motion and is given by Chen & Taam (1992); and

$$D_O = \frac{4\pi}{c^2}\int_{-h}^{h} dz\eta |J|^2 = \frac{\eta}{4\pi h}\left[2(B_r)_h^2 + \frac{3}{5}(B_\phi)_h^2\right] \quad (11c)$$

is the ohmic dissipation in the disk;

$$F_R = \frac{2acT^4}{3\kappa\sigma} = \sigma_B T_{\rm eff}^4 \quad (11d)$$

is the radiative energy flux per unit area from the top or bottom surface of the disk which is assumed optically thick; a and σ_B are the usual radiation constants; κ is the opacity; $T_{\rm eff}$ is the effective temperature of the surface of the disk; and $\mathscr{F}_r = -[(8acT^3)/(3\kappa\rho)](\partial T/\partial r)$ is the radial flux of radiation in the disk.

The central object is considered to be a Schwarzschild black hole so that the inner radius of the disk is $r_i = 6GM/c^2 \approx 0.9 \times 10^{14}$ cm $(M/10^8\ M_\odot)$. The "i" subscript denotes values at r_i. The dimensionless radial distance is $R \equiv r/r_i$. The reference speed is the Kepler velocity at r_i, $v_{Ki} = (GM/r_i)^{1/2}$ (neglecting general relativity). All velocities are measured in units of v_{Ki}. The reference time is $1/(2\pi)$ of the Kepler period at r_i, $t_i = r_i/v_{Ki} \approx 1100$ s $(M/10^8\ M_\odot)$. The dimensionless time is $T = t/t_i$. The context distinguishes T from the temperature also denoted T.

We are mainly interested in the inner part of the disk ($R \lesssim 10^3$) and therefore consider that the opacity is due to electron scattering. In the absence of an ordered B field, the stationary disk solution, with viscosity proportional to gas pressure and vanishing viscous stress at r_i, has: $\sigma = \sigma_i J^{3/5} R^{-3/5}$, $T = T_i J^{2/5} R^{-9/10}$, and $u = u_i J^{-3/5} R^{-2/5}$, where $J(R) \equiv 1 - R^{-1/2}$, $\sigma_i \propto (\dot{M}^3/\alpha^4 M^2)^{1/5}$, $T_i \propto (\dot{M}^2/\alpha M^3)^{1/5}$, and $u_i \propto (\alpha^4 \dot{M}^2/M^3)^{1/5}$, with $\dot{M}$ the mass accretion rate. The singular dependences of the disk quantities as $R \to 1$ are treated in the numerical simulations by taking boundary conditions on u, σ, and T at $R \to 1$ corresponding to the above dependences with J replaced by $J = 1 - R^{-1/2} + \delta$ with $\delta \ll 1$.

We measure the surface density in units of σ_i, the temperature in units of $\mu v_{Ki}^2/k_B$, and the magnetic field in units of B_0, where B_0 is a characteristic value of the B_z field. In the dimensionless equations, the magnetic terms all have a dimensonless coefficient $\epsilon_i \equiv r_i(B_0)^2/(2\pi\sigma_i v_{Ki}^2)$, which is the square of the ratio of an Alfvén speed to v_{Ki}. The radiative luminosity of the nonmagnetized equilibrium disk is $L_0 = GM\dot{M}/(2r_i) \approx 0.48 \times 10^{45}$ ergs s^{-1} $(\dot{M}/0.1\ M_\odot\ {\rm yr}^{-1})$. For reference, the Eddington luminosity is $L_{\rm Edd} \approx 1.3 \times 10^{46}$ ergs s^{-1} $(M/10^8\ M_\odot)$. We write $\epsilon_i = r_i^2 \alpha T_i B_0^2/(\dot{M} v_{Ki})$, where T_i is the dimensionless temperature constant for $B = 0$ equilibrium disk. Thus

$$\epsilon_i \approx 10^{-6}\left(\frac{\alpha}{0.1}\right)^{4/5}\left(\frac{M}{10^8\ M_\odot}\right)^{7/5}\left(\frac{0.1\ M_\odot\ \mathrm{yr}^{-1}}{\dot{M}}\right)^{3/5}\left(\frac{B_0}{5.7\ \mathrm{kG}}\right)^2 .$$

The B fields considered here are weak in the sense that they have a negligible influence on the radial force balance. Also, the B field influence on the disk thickness is small compared with that of the radiation and gas pressure. A useful dimensionless measure of the radiation pressure is $\epsilon_R = 2(p_{\mathrm{rad}})_i/(\sigma_i v_{\mathrm{K}i}^2) \approx 0.05(10^8\ M_\odot/M)(\dot{M}/0.1\ M_\odot\ \mathrm{yr}^{-1})$, where $(p_{\mathrm{rad}})_i \equiv aT_i^4/3$. Thus, $\epsilon_R \approx 1.33(L_0/L_{\mathrm{Edd}})$. For the reference values $M = 10^8\ M_\odot$, $\dot{M} = 0.1\ M_\odot\ \mathrm{yr}^{-1}$ and $\alpha = 0.1$, one find $T_i \approx 2.8 \times 10^{-6}$ (or $T_i \approx 3.2 \times 10^6$ K), $\sigma_i \approx 5.5 \times 10^6$ g cm^{-2}, and $\epsilon_R \approx 0.05$. This time for a fluid particle to spiral in from a radius $R_0 \gg 1$ due to viscosity ($B = 0$) is very long, of the order of 90 yr $R_0^{7/5}$, for the mentioned reference values.

Numerical results have been obtained using two codes. The first uses a second-order Lax-Wendroff method (Richtmeyer & Morton 1967) to evolve: $R\sigma$ (eq. [2]); $R\sigma u$ (eq. [3]); $R\sigma l$ (eq. [6]; RB_z (eq. [9]); and $R\sigma T$ (eq. [11]); with h is gotten from equation (7). The second, much faster quasi-static code, is similar to the first except that (1) we neglect the inertial term in the radial equation of motion (3), which then takes the form $l^2 = GMr + (r^3/\sigma)(\partial P/\partial r)$ + magnetic + viscous terms, and; (2) we neglect $\partial l/\partial t$ in the angular momentum equation (6) which then takes the form $ur\sigma(\partial l/\partial r)$ = viscous + magnetic terms and directly gives $u(r, t)$. The Lax-Wendroff method is also used in this code to evolve $R\sigma$, RB_z, and $R\sigma T$. We find essential agreement between the two codes for dimensionless times up to 10^5 (limited by our computer resources). The results discussed here are based on the second faster code. The inner radius of the disk is treated as a free boundary in the respect that σ, T, and B_z needed at the first grid point each full time step are obtained by linear extrapolation. At the outer boundary, T is obtained by linear extrapolation; σ and B_z are taken to be a constant equal to their initial values. In the absence of a magnetic field, both codes give convergence to the Shakura-Sunyaev solution when started from even highly nonequilibrium conditions.

In the above mentioned quasi-static approximation, the radial accretion speed is $u = u_v + u_B$, where u_v is the contribution due to viscosity, and $u_B = -r(B_\phi)_h B_z/(\pi v_{\mathrm{K}} \sigma) \approx \beta_\phi B_z^2 r^{3/2}/[\pi\sigma(GM)^{1/2}]$ is the accretion speed which results from the removal of angular momentum by the outflows. Thus, the dynamics of a magnetized disk is controlled by the evolution of $B_z(r, t)$ and $\sigma(r, t)$. In the B_z equation (9), the inward advection ($\propto u$) is counteracted by the outward diffusive drift ($\propto \eta$), while the diffusion term is generally smaller. The nonlinear dependence of u_B on B_z suggests that an enhancement or bump in B_z at, say, r_0 can make the inward advection of B_z dominant over the outward drift. Consequently, advection will cause both σ and B_z to increase with time. This can give rise to what we term implosive accretion.

A disk initialized with a distributed B_z field, for example, $B_z = B_0 R^{-5/4}$, suggested by Blandford & Payne (1982), remains approximately steady for a long time ($T \sim 10^6$) with an approximately constant ratio $\mathscr{R}$ of the power in the $\pm z$ winds or jets to the radiated power of the disk. The power flow to the winds per unit radial distance is $-2\pi(r_i \sigma_i v_{\mathrm{K}i}^3)\epsilon_i \tilde{\omega} R^2 \tilde{B}_\phi \tilde{B}_z$, where the over tildes indicate dimensionless quantities. We find $\mathscr{R} \approx 4.5(\epsilon_i/10^{-6})$ for $10^{-7} < \epsilon_i < 10^{-5}$ (the range studied). The disk's σ and T remain approximately equal to the Shakura-Sunyaev values while u is larger than the equilibrium value by a factor $\approx 1 + \mathscr{R}$. Similar behavior is found for an initial field $B_z = B_0/R$ for which $\mathscr{R} \approx 9.0(\epsilon_i/10^{-6})$ for the same range of ϵ_i. Notice however that there are no absolute steady-state solutions to equations (2)–(11) when the radiation pressure is dominant.

A disk initialized with both a distributed B_z field and an enhancement or "bump" in B_z in the vicinity of some large radius r_0 evolves with the bump moving inward and diffusively spreading. For a larger amplitude of the initial bump in B_z and/or a larger background field, the bump moves inward more rapidly, it narrows in radial width, and it forms a soliton-like disturbance. The disturbance propagates inward while growing in strength up to the time when it is "swallowed," that is, when it passes within r_i. The disturbance produces almost simultaneous bursts in the power output of the winds and in the disk radiation. Figures 1 and 2 show the main quantities for sample cases.

For a fixed ratio of the peak amplitude of the bump to that of the background, the temporal width of the burst is a strongly decreasing function of the background field. On the other hand, for a fixed background field the burst width is a strongly decreasing function of the initial bump amplitude. Table 1 shows the main properties of the bursts resulting from bumps started at $R = 20$ for the reference value $M = 10^8\ M_\odot$, $\dot{M} = 0.1\ M_\odot\ \mathrm{yr}^{-1}$, and $\alpha = 0.1$, and a background field $B_z = B_0/R$ with $\epsilon_i = r_i(B_0)^2/2\pi\sigma_i v_{\mathrm{K}i}^2$). A is the ratio of the peak of the bump to the background field, and ΔT is the full-width at half-maxumum of the burst in the jet power. A dimensionless time $T = 10^6$ corresponds to 35 yr. For comparison, the viscous ($B = 0$) accretion time from $R = 20$ to $R = 1$ is $\sim$6000 yr. Figure 3 shows the disk behavior for a case where there is a gradual variation in the initial B_z field amplitude.

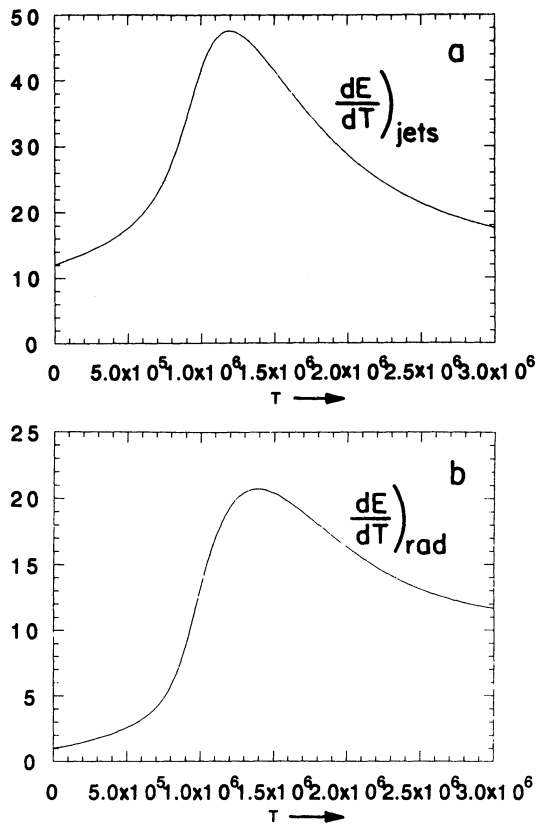
3. DISCUSSION

This work has developed a theoretical model of time-dependent magnetized accretion disks which exhibits short timescale outbursts of high-power winds or jets associated with bursts of disk radiation. Prediction of the time-dependent spectrum of an active galactic nucleus outburst is beyond the scope of the present work because it requires an analysis of the strongly time-dependent dynamics, radiation, and radiation transfer of an impulsively driven relativistic magnetized jet.

TABLE 1

BURST PROEPRTIES

Background Field	Bump Amplitude A	ΔT_{FWHM} (jets)	$\left[\frac{dE}{dT}\Big\vert_{\max} \Big/ \frac{dE}{dT}\Big\vert_{\mathrm{init}}\right]_{\mathrm{jets}}$	$\left[\frac{dE}{dT}\Big\vert_{\max} \Big/ \frac{dE}{dT}\Big\vert_{\mathrm{init}}\right]_{\mathrm{rad}}$
$\epsilon_i = 10^{-6}$	2	1.65×10^6	4.1	21.5
	3	0.46×10^6	9.8	47.1
$\epsilon = 3 \times 10^{-6}$	2	0.22×10^6	6.9	31.5
	3	0.08×10^6	14.4	66.8



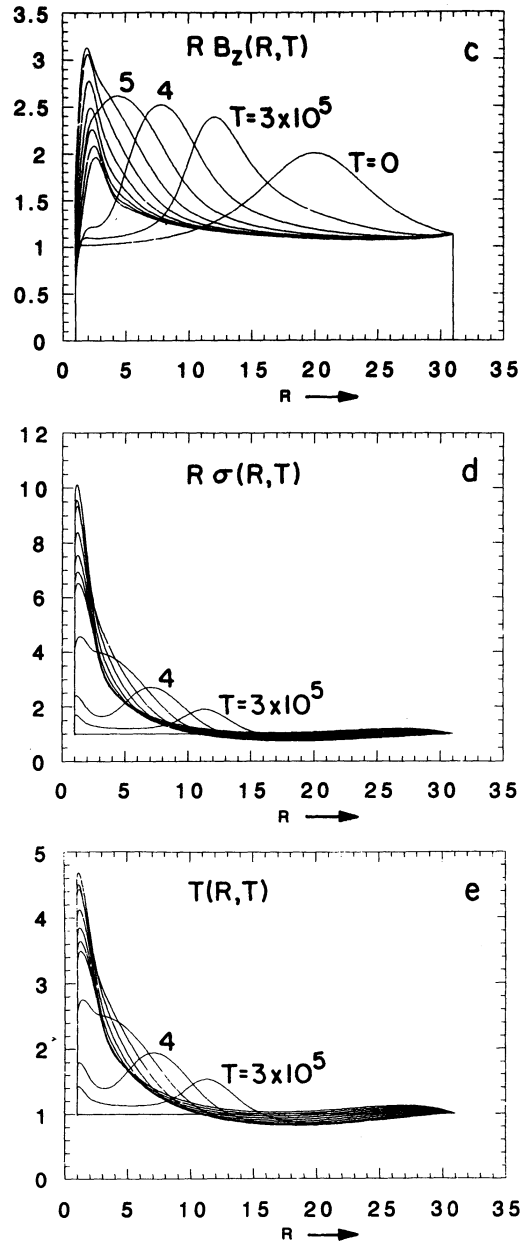
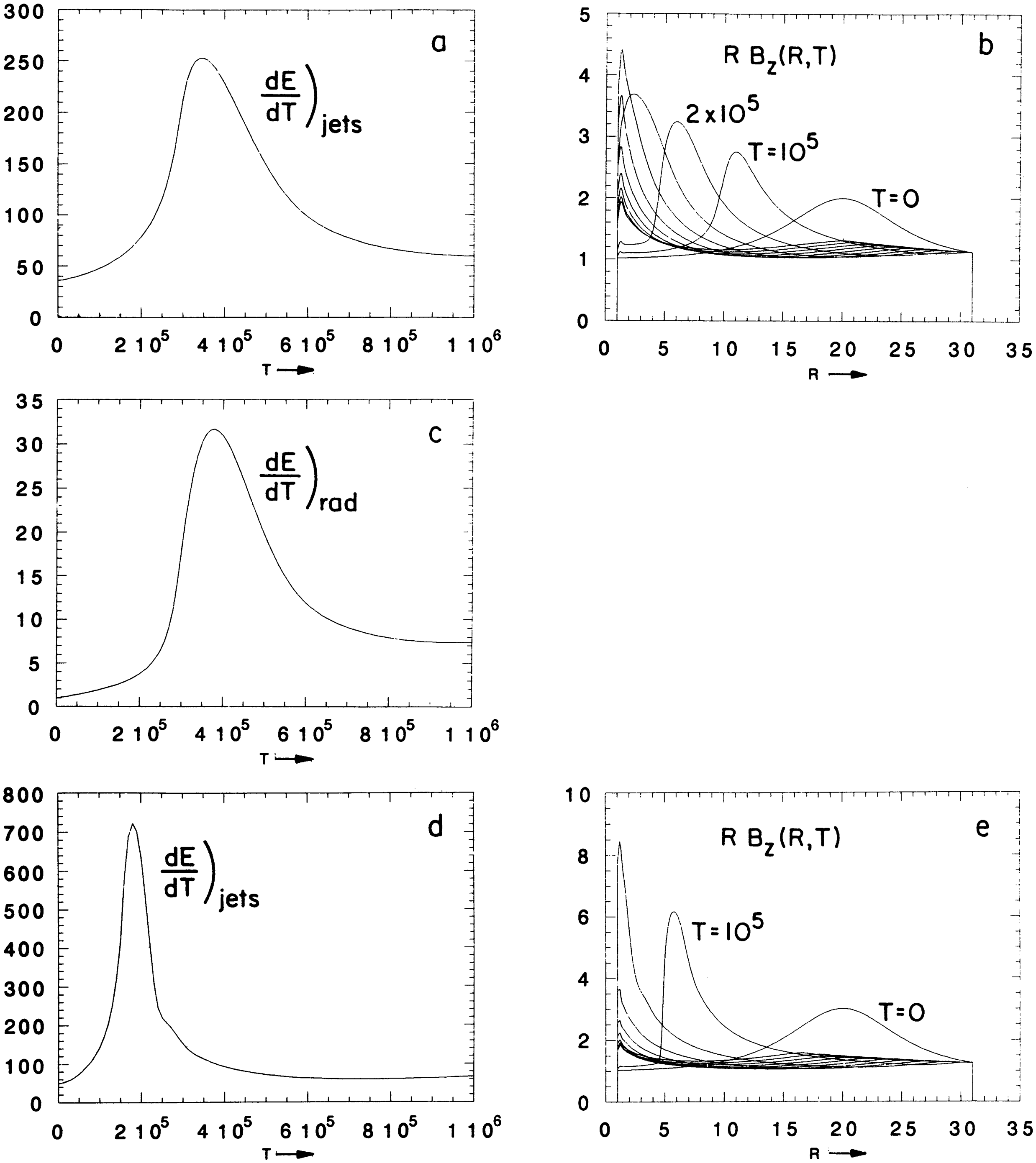


FIG. 1.—Disk behavior for a case where the initial field consists of a background component $B_z = B_0/R$, a bump in B_z centered at $R = 20$, of full-width to $1/e$ of $\Delta R = 8$, and of peak amplitude twice the background (see [*c*]). For this case, $\epsilon_i = r_i(B_0)^2/(2\pi\sigma_i v_{Ki}^2) = 10^{-6}$, $\alpha = 0.1$, $\epsilon_R = 0.05$, and $D = 1$ for the magnetic diffusivity. The power in the jets (*a*) and in the disk radiation (*b*) are measured in units of the radiative power of the nonmagnetized disk, L_0, discussed in the text. The disk surface density in (*d*) and temperature in (*e*) are measured in units of the Shakura-Sunyaev values. The B field is measured in units of B_0. In this simulation, the outer radius is $R_0 = 31$ and $\Delta R = 0.1$. The vertical lines at $R = 1$ and $R = 31$ are retrace artifacts of the plotting routine.

The burst of energy in jets found here will produce outward propagating disturbances in preexisting ($\pm z$) jet flows. The disturbances are strong and can be expected to steepen and form strong shock waves as assumed in the model of Blandford & Königl (1979) and Hughes, Aller, & Aller (1989a, b). The strong linear polarization (up to $\sim 30\%$) and large polarization position angle changes (90° and more) seen in the optical outbursts of some quasars (Kinman 1967; Kinman et al. 1968; Babadzhanyants & Hagen-Thorn 1975; Moore & Stockman 1981, 1984) clearly indicate synchrotron radiation, and this component of the radiation is probably associated with the burst in jet energy in the present model.

The origin and nature of the B field in disks of active galactic nuclei is clearly important to the present model. The mentioned high polarization of the optical radiation points to ordered, nonturbulent, fields in the nuclei of some objects. Evidence for appreciable ordered poloidal and toroidal B fields (10^{-4} to 10^{-3} G) on scales 1–10 pc in the center of our Galaxy comes from Zeeman splitting measurements of the H I radio line (Schwarz & Lasenby 1990), infrared dust polarization measurements (Hildebrand et al. 1990), and high-resolution radio mapping (Yusef-Zadeh, Morris, & Chance 1984; Morris 1990). On the theoretical side, the $\alpha - \omega$ dynamo has been studied intensively as a likely mechanism for generating both large and small-scale B fields (Pudritz 1981a, b; Stepinskii & Levy 1988; Pringle 1989; Geertsema & Achterberg 1992). The instability of Balbus & Hawley (1992) may be important in initiating the dynamo process. At large radial distances in a disk, the accretion speed is low if the B field is small. Thus a dynamo-generated field can grow with negligible convective losses. The lowest order dynamo mode has a dipole-like poloidal field which would give a positive bump in B_z at one radius and a negative bump at another. Once the dynamo field reaches a large enough amplitude, it can precipitate the enhanced accretion discussed in this work. Alternation of the polarity of the disk B_z field acts to limit the growth of the total B_z flux within r_i.

Recent observations (Reipurth 1989; Reipurth & Heathcote 1991) and theoretical work (Raga et al. 1990) support the idea that jets from T Tauri stars are time-varying. Changes in the jet



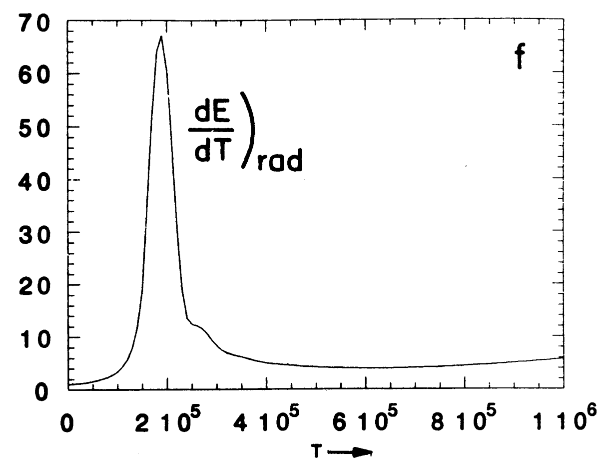
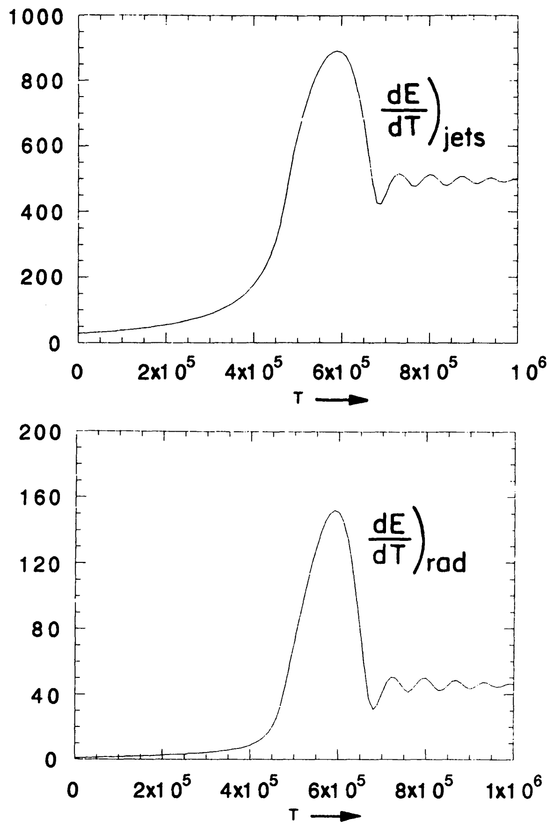


FIG. 2.—(*a*)–(*c*) Disk behavior for a case the same as Fig. 1, except that the background field is larger by $3^{1/2}$, that is, $\epsilon_i = 3 \times 10^{-6}$. (*d*)–(*f*) Disk behavior is shown for $\epsilon_i = 3 \times 10^{-6}$ and a peak bump amplitude 3 times the background. Our treatment of the disk requires that the radiative energy flux is less than the Eddington luminosity. However, the energy flux in the winds or jets can exceed the Eddington value without inconsistency.



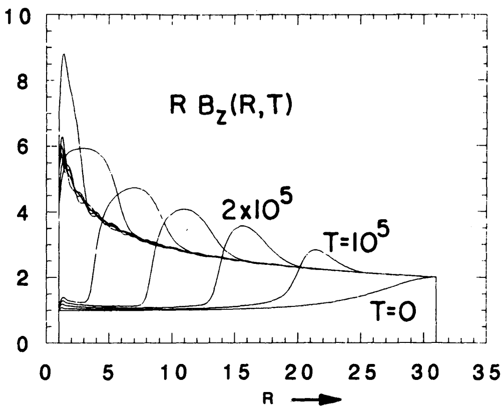


FIG. 3.—Disk behavior for a case similar to Fig. 2, except that in place of the initial bump there is a gradual, factor of 2 variation in the initial B_z field amplitude.

speed and energy flux are predicted to give rise to internal shocks in the jets which appear as bright knots referred to as Herbig-Haro objects (Raga et al. 1990). These jets are thought to be magnetically driven from circumstellar disks (see, for example, Königl & Ruden 1993; LBC; LRC). The time variations in stellar jets may be due to variations in the distribution of magnetic field threading an accretion disk analogous to the situation treated here. An analysis of the dynamics of a magnetized T Tauri disk needs to account for the more complicated opacity law (Bell & Lin 1993). This opacity law is found to give a thermal instability which greatly increases the accretion speed possibly explaining the spectacular outbursts of FU Orionis objects (Bell & Lin 1993). A magnetic field threading the disk may be greatly amplified by the increased accretion speed during the outburst. This increased field can in turn give strongly enhanced winds as observed for FU Orionis stars (Herbig 1966, 1977; Bastian and Mundt 1985).

Note added in manuscript.—Two recent studies by Lubow, Papaloizou, & Pringle (1994a, b) of the time-dependent behavior of magnetized disks have come to our attention. In essential respects these works appear to be compatible with the present.

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