ACCRETION DISK ELECTRODYNAMICS

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Accretion disk electrodynamic phenomenae are separable into two classes: 1) disks and coronae with turbulent magnetic fields; 2) disks and black holes which are connected to a large-scale external magnetic field. Turbulent fields may originate in an $\alpha - \omega$ dynamo, provide anomalous viscous transport, and sustain an active corona by magnetic buoyancy. The large-scale field can extract energy and angular momentum from the disk and black hole, and be dynamically configured into a collimated relativistic jet.

1.0 INTRODUCTION

This review emphasizes some recent developments on accretion disk and black hole electrodynamics. The broader subject of accretion disk structure has been excellently reviewed by Pringle (1981). Regrettably, we cannot discuss accretion onto magnetized neutron stars or white dwarfs.

Accretion disk electrodynamics can be divided into two broad points of view, which are not necessarily exclusive. The first view argues that the strong Keplerian differential rotation and convective turbulence within the disk produces a highly disordered magnetic field which is susceptible to Parker's (1972; 1979) topological dissipation - the rapid reconnection relaxation of localized regions of turbulently induced magnetic shear. Although it may be primordial, the disk field is probably maintained against buoyancy losses by a regenerative $\alpha - \omega$ dynamo. The alternative view is that the disk remains connected to a large-scale external magnetic field. The disk and the black hole possess a magnetosphere which can inertially extract angular momentum from the Keplerian flow and the hole's rotational energy, and establish outflowing jets aligned along the rotation axis.

In Section 2.0 we briefly review the standard accretion disk model. Section 3.0 discusses the turbulent view of disk electrodynamics.

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Section 4.0 discusses hydromagnetic disk winds, and in Section 5.0 we review black hole electrodynamics.

2.0 ACCRETION DISK HYDRODYNAMICS - THE STANDARD MODEL

Although earlier progenitors exist, the (so-called) standard model of a thin accretion disk surrounding a central black hole was developed by Shakura and Sunyaev (1973) and Novikov and Thorne (1973) who solved the conservation laws of steady state hydrodynamics in conjunction with radiation transport. The strong central gravity requires the disk material to be in circular Keplerian orbits. The vertical component of gravity is balanced by the gradient of the total pressure; in the inner region near the black hole, radiation pressure (P_r) usually greatly exceeds the thermal gas (P_g) pressure. A radially inward accretion flow results from the viscous dissipation of orbital angular momentum; viscosity diffuses angular momentum to larger radii thus allowing the gas to diffuse inward. For rapid accretion the viscous stress ($t_{r\phi}$) must be anomalous, arising from either hydrodynamic and/or hydromagnetic turbulence; a common ansatz is to scale $t_{r\phi}$ to one of the basic pressures - P_r, P_g, or the magnetic pressure $B^2/8\pi$.

Accretion liberates gravitational binding energy which must be radiated or transported vertically away from the disk. For a mass accretion rate $\dot{M}$, the total luminosity is $L = \varepsilon \dot{M} c^2$ where ε is determined by the radius of the last bound orbit; $\varepsilon = 0.06(0.42)$ for a Schwarzschild (maximally rotating Kerr) black hole. An accretion rate of 10 $M_\odot$/yr ($10^{-8} M_\odot$/yr) will produce a luminosity of $\sim 10^{47}$ (10^{38})ergs/sec which is comparable to the energy output of quasars and active galactic nuclei (galactic x-ray sources). This very efficient conversion of rest mass energy to luminosity remains the strongest argument that energetic astrophysical objects are powered by accretion onto black holes.

In the standard model, the energy output is radiative, and occurs as either black body radiation if the disk is optically thick, or as various forms of Comptonized radiation if electron scattering dominates free-free absorption. The effective radiation temperature of stellar mass ($\sim M_\odot$) scale disks corresponds to x-ray energies, whereas for supermassive ($\sim 10^9 M_\odot$) disks the effective temperature drops 10^5 to 10^6°K.

Critiques of the standard model abound in the literature (see Pringle, 1981). For our purposes, 3 points are of particular significance. The anomalous viscosity mechanism is highly uncertain. The vertical structure of a radiation-supported disk is quite likely to be thermally unstable (Bisnovatyi-Kogan and Blinnikov, 1977). A thermally convective disk might have an active magnetic dynamo, a mechanically or hydromagnetically excited corona, and an escaping wind or jet. If a significant fraction of the disk's energy output is deposited in a corona or a wind, the disk structure and predictions of the standard model would be fundamentally altered.

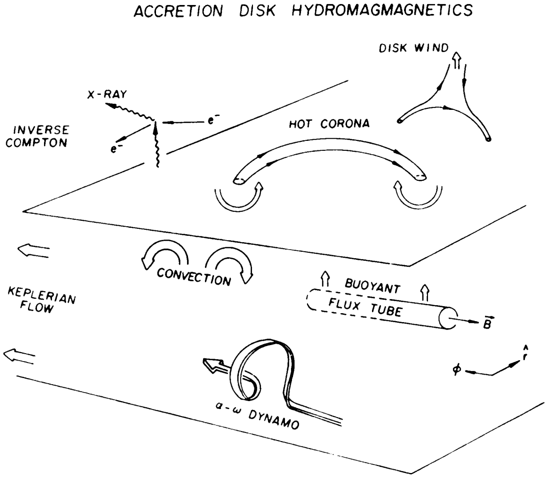

Figure 1. Hydromagnetic phenomenae in thin accretion disks.

3.0 ACCRETION DISK HYDROMAGNETICS

Since accretion disks crudely resemble flat stars, their hydromagnetic structure and dynamics possess the full panoply of difficulties associated with solar and stellar magnetism (Figure 1). For vertically thin disks, a basic question is whether magnetic buoyancy (Parker, 1979) will rapidly remove magnetic flux from the disk (Section 3a). We summarize the attempts to incorporate magnetic dynamos into the theory of disks in Section 3b. A model for the anomalous viscosity due to turbulent magnetic Maxwell stresses is presented in Section 3c. Magnetized disk coronae are discussed in Section 3d.

3.a. Magnetic Buoyancy

In a gravitationally bound atmosphere, a magnetic flux tube experiences a buoyancy force since pressure equilibrium requires that the tube's internal density be less than that in the surrounding atmosphere (Parker, 1979). For an isothermal, gas pressure-dominated atmosphere a flux tube with $B^2/8\pi \ll P_g$ is weakly buoyant, and rises (against gravity) slowly until it reaches density equilibrium. If $B^2/8\pi \gg P_g$, the buoyant rise is rapid, reaching a drag-limited velocity comparable to the Alfven speed. In an adiabatic atmosphere, the density contrast is preserved as the tube rises so that the atmospheric pressure gradient expels even weakly magnetized flux tubes (Moreno-Insertis, 1983).

For an accretion disk we have two additional considerations: 1) in the inner region ($P_r \gg P_g$) does strong buoyancy set in when $B^2/8\pi \sim P_g$ or P_r; 2) differential rotation can shear the flux tube thereby attempting to change the flux tube's volume and field strength. Coroniti (1981) and Stella and Rosner (1983) have argued that flux tubes become strongly buoyant when $B^2/8\pi \sim P_g$, independent of the ratio P_r/P_g. Recently Sakimoto (1983) has carried out extensive numerical calculations on buoyant flux tubes including the effects of shear, heat flow into the tube, inertia, and Newtonian drag. For $P_r >> P_g$, the general conclusions are: 1) heat flow maintains the tubes near temperature equilibrium; 2) for an adiabatic disk atmosphere, even $B^2/8\pi << P_g$ tubes are strongly buoyant, although shear can increase the magnetic pressure to above P_g; 3) for an isothermal atmosphere shear drives even $B^2/8\pi << P_g$ tubes out of the disk; in the absence of shear, $B^2/8\pi << P_g$ tubes rise until density equilibrium is reached, while $B^2/8\pi >> P_g$ tubes are rapidly lost.

Most flux tubes reach a drag-limited velocity $V \sim C_s\ (R_T\Delta\rho/H\rho)^{1/2}$ where C_s is the sound speed based on the total pressure ($C_s{}^2 \sim P_r/\rho$), R_T is the tube radius, H is the disk scale height ($H \approx C_s/\Omega$ where Ω is the Keplerian angular velocity) and $\Delta\rho/\rho$ is the density contrast. The buoyant flux tubes rise much faster than the Alfven speed ($<< C_s$), and escape from the disk in 1 to 10 Keplerian rotation periods. Hence buoyancy places the stringent limit $B^2/8\pi < P_g$ on the magnetic field strength which can be retained in an accretion disk.

3.b. Accretion Disk Dynamos

The strong differential rotation and likely thermal convective turbulence have tempted several authors to develop an $\alpha - \omega$ dynamo model (Parker, 1955; 1979; Moffatt, 1978) for accretion disks. As a maximal model Takahara (1979) assumed that all of the accretion energy was invested in the dynamo generation of a strong toroidal magnetic field. Since the disk is cold and thus magnetically supported against gravity, the flux is buoyantly lost on the Alfven time scale, and accumulates in a low density corona. A similar model was developed by Galeev, et al. (1979) who argued that meridional convective turbulence could lead to the exponential growth of the toroidal field.

In these models the dissipation of accretion energy into heat and radiation occurs in the corona where the high Alfven speed permits rapid reconnection of the field. Although a magnetically active corona is attractive (see next section), the underlying dynamo disk models remain incomplete since the source of the convective turbulence required for the regenerative α-effect is uncertain. If the disk has a finite temperature so that the convection is thermally driven, it may be difficult to avoid $P_r >> P_g$ in the disk's inner region; if so, the buoyancy limit $B^2/8\pi < P_g$ would severely reduce the strength of the coronal magnetic fields and reconnection dissipation. As an alternative, Stella and Rosner (1981) suggest that the buoyancy instability may generate the necessary convective turbulence in the disk; although promising, this "bootstrap" dynamo model requires further elaboration.

In a different approach, Pudritz (1981a,b) and Pudritz and Fahlman (1982) have applied the Mean Field Dynamo Theory (Steenbeck and Krause, 1969; Moffatt, 1978) to accretion disks. Pudritz (1981a) attempts to

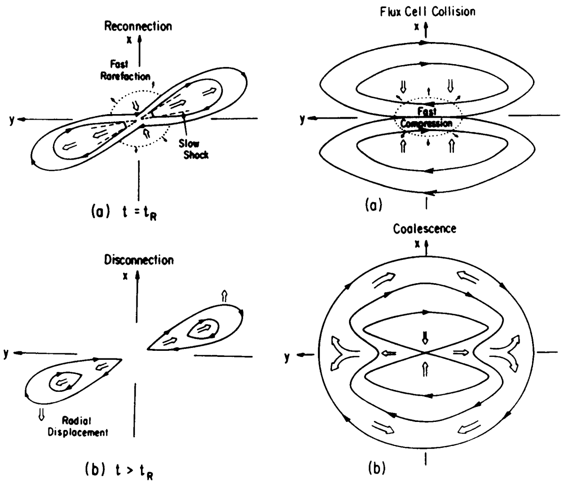

Figure 2. Flux cell dynamics and magnetic viscosity.

self-consistently determine the turbulent magnetic stresses which produce the required anomalous viscosity and to calculate the dynamo-generated large-scale magnetic field (Pudritz, 1981b). A potential difficulty is that the turbulent magnetic field is much larger than the spatially-averaged mean field, which is comparable to the radiation pressure in the inner disk (Pudritz and Fahlman, 1982). Buoyant losses will greatly reduce the strength of the turbulent field on much shorter time scales than the Ohmic diffusion time.

In conclusion, buoyancy severely limits the strength of the turbulent magnetic fields in accretion disks. Even the presence of a relatively weak magnetic field probably requires some form of dynamo activity to replace the buoyancy flux losses. However, as for the solar dynamo problem, a self-consistent accretion disk dynamo model and its coupling to a magnetically active corona still resides in the future.

3.c. Magnetic Viscosity

Early disk models (Lynden-Bell, 1969; Shakara and Sunyaev, 1973; Novikov and Thorne, 1973) recognized that a turbulent Maxwell stress ($B_r B_\phi / 4\pi$) would exert local torques, which could diffuse angular momentum. Eardley and Lightman (1975) suggested a model for this magnetic viscosity in which topological dissipation breaks up the initial magnetic field into localized flux cells; for simplicity, we assume that the cells are oriented parallel to the disk plane and extend roughly over the disk height. Figure 2 outlines the "life of a flux cell" in a later version of this model (Coroniti, 1981).

We assume that the initial magnetic field is weak ($B^2/8\pi << P_g$) in order to minimize buoyancy losses. The Keplerian differential rotation (y-direction) stretches the flux cell until the Maxwell stress builds up to stop the shear flow within the cell; the plasma within the cell shear increases the field strength (at the expense of the flow kinetic energy) and decreases the radial (x-direction) extent of the cell; both effects greatly reduce the time scale for the oppositely-directed magnetic fields to reconnect at the center. Fast reconnection bifurcates the cell, and dissipates the stored magnetic energy into heat across the slow shocks which stand in the incident flow (Petschek, 1964).

After disconnection, the plasma in the inner (outer) cell has too little (too much) angular momentum to be in local Keplerian equilibrium; the cells move radially until a new equilibrium is achieved. If the cells are closely packed, collisions between cells result in their reconnection coalescence; during coalescence the internal plasma mixes within the new cell. Averaging over many disconnection-coalescence events, the plasma undergoes a random walk in radial location.

The radial diffusion coefficient (ν) depends on the size of the flux cells (L) and the reconnection-coalescence time scale. If $L > C_g/\Omega$ where $C_g = (P_g/\rho)^{1/2}$ is the gas sound speed, the shear flow raises $B^2/8\pi$ above P_g, and buoyancy rapidly removes the cell from the disk; thus $L < C_g/\Omega$. The reconnection time depends on the shear distortion and compression of the flux cell, and is estimated to be $\lambda t_R \approx \approx M_A{}^{-1}(C_A/\lambda L)^{1/3}$ where $\lambda = 3/2\ \Omega$, C_A is the Alfven speed, and M_A is reconnection Alfven Mach number which is of order $M_A \sim 1/10$ to $1/20$ (Petschek, 1964; Parker, 1979). We then obtain $\nu \sim 9/4\ C_g{}^2 M_A{}^{2/3}/\lambda$, which yields an anomalous viscous stress of $t_{r\phi} = \rho\nu\lambda \approx 9/4\ M_A{}^{2/3} P_g$.

Although greatly over-simplified, the flux cell model demonstrates that a dynamic disk magnetic field can provide an anomalous viscous transport. The scaling of $t_{r\phi}$ with P_g is consistent with the conclusion of Stella and Rosner (1983) and leads to accretion disks which are much denser and somewhat hotter in their inner regions than predicted by the standard model with $t_{r\phi} \propto P_r$ (Sakimoto and Coroniti, 1981).

3.d. Accretion Disk Coronae

An inevitable consequence of convective turbulence and magnetic buoyancy will be to form a hot, magnetically active corona (Stella and Rosner, 1983). The significance of coronal dissipation is that the hard x-rays observed form Cygnus X-1 and the softer x-rays from quasars cannot originate in the disk, but must be emitted in the corona. The soft disk photons are Compton up-scattered to x-rays by the hot or mildly relativistic coronal electrons (Price and Thorne, 1975; Liang and Price, 1977; Eardley, et al., 1978).

A difficult problem is to determine what fraction of the disk's energy output goes into heating the corona, which, in turn, will bound the system's x-ray luminosity. Drawing on recent developments in solar physics, Galeev, et al. (1979) suggested that the corona is confined and heated in strong magnetic loops which have buoyantly emerged from

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the disk. The twisted magnetic fields of the loops are dissipated by rapid reconnection in the high Alfven speed coronal plasma. Instead of being cooled by thermal conduction or bremstrahlung, the hot electrons in the loops lose energy by Compton scattering the soft disk photons. Along similar lines, Ionson (1983) has scaled the solar coronal heating function to disks, and concluded that a large fraction of the disk's energy can be dissipated in the corona. He argues that the corona may be optically thick to Compton scattering so that the emergent coronal spectrum will have the power law form of unsaturated Compton scattering (Shapiro, et al., 1976; Rybicki and Lightman, 1979); a power law spectrum is observed for the hard x-ray component of Cygnus X-1 and in quasars.

4.0 HYDROMAGNETIC DISK WINDS

The existence of an active corona suggests that disks should have winds (Piran, 1977; Liang and Thompson, 1979). Since the gravitational binding energy is of order 450 R_s/r MeV ($R_s = 2GM/c^2$ is Schwarzschild radius), a thermally driven wind is unlikely. However, if the disk's luminosity exceeds the Eddington limit ($\sim 10^{38}$ $M/M_\odot$ ergs/sec), the disk's inner region becomes almost spherical and a radiatively driven wind will develop (Shakura and Sunyaev, 1973; Meir, 1979, 1982). In addition, if the outer disk flares appreciably, the absorption of x-rays emitted from the inner region can stimulate a thermal wind (Begelman, et al., 1983). Although thermal and radiative winds undoubtedly occur, we will limit discussion to electrodynamically driven winds.

Despite the cogent arguments that disk magnetic fields should be turbulent, the disk may remain connected to the large-scale field which originally threaded the infalling plasma. Blandford (1976) realized that the external field and plasma are coupled to the disk by a system of field-aligned currents which close by flowing across the disk's magnetic field. The closure currents exert a $\underline{J} \times \underline{B}$ spin-down torque on the disk, thus extracting angular momentum from the Keplerian flow and permitting the disk material to flow radially inwards. Disk angular momentum is carried away either by Alfven (or electromagnetic) waves, or, if the disk also loses mass, as part of a hydromagnetic wind.

Blandford and Payne (1982) have proposed a rotationally driven hydromagnetic wind which takes the form of a magnetically collimated jet aligned along the disk's rotation axis (Figure 3). Extragalactic radio sources have large-scale jets (10kpc to 2Mpc) which are often aligned with the VLBI jets (~ 1pc) emanating from the nuclei of galaxies and quasars (compact radio sources)(Miley, 1980). Their high luminosity implies that the jets are powered by gravitational accretion or by the central black hole (Blandford and Znajek, 1977; Macdonald and Thorne, 1982; next section).

Blandford and Payne (1982) assume that the magnetic pressure of the large-scale field dominates the coronal gas pressure. The coronal

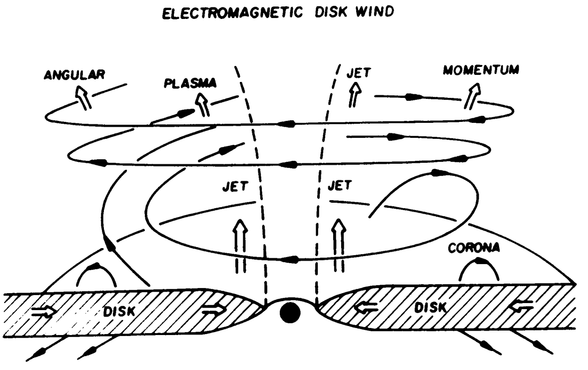

Figure 3. Centrifugally driven hydromagnetic disk wind.

magnetic field is assumed to rigidly connect to the disk so that the field lines rotate with the Keplerian velocity at the radial distance where they enter the disk; this requires that the Alfven speed in the corona must exceed the Keplerian velocity. Since the disk's field convects with the plasma, the magnetic pressure in the disk must be much less than the dynamic pressure of the Keplerian flow. Hence in order to obtain the required high coronal Alfven speed, the coronal density must be low so that the hydromagnetic wind transports only a relatively low mass and particle angular momentum flux.

If the coronal field lines make an angle of less that 60° to the disk plane, the cold coronal plasma is centrifugally slung outwards along the rigid lines of force. Near the wind's Alfven singular point, the plasma's inertia bends the lines of force forming a toroidally wound field which carries off the disk's angular momentum. As is typical with cold magnetic sling solutions, the wind does not pass through a fast critical point (which formally resides at infinity)(Mestel, et al. 1979; Michel, 1973). Since the asymptotic wind remains dominated by the magnetic Poynting flux, the hoop stress of the toroidal field collimates the flow into a jet. (The wind equations also have a solution in which the fast Mach number exceeds unity at infinity, but the jet is pinched-off by the magnetic field at a finite distance above the disk. These solutions probably correspond to the unphysical solutions of radial stellar wind theory (see Kennel, et al., 1983). The physical solutions with a fast Mach number of unity at infinity are not pinched; these will remain the correct physical solutions when finite pressure is included.)

Since the rate of gravitational energy liberation maximizes in the disk's inner region, the jet's velocity is largest near the rotational axis, and decreases with radius; consequently the jet has a high speed core which can approach relativistic velocities. The transport of angular momentum maximizes in the outer region of the jet. The efficiency with which the wind extracts energy and angular momentum from the disk depends on the strength of the large scale field. In the Blandford-Payne model, the critical field strength at which the entire disk luminosity (L_D) is carried off by the jet is roughly $B_c \simeq (L_D/\sqrt{GMr_o^3})^{1/2}$ where r_o corresponds to the inner disk radius; for $M = M_\odot (10^9 M_\odot)$, $B_c \sim 10^8 (10^4)$ Gauss which is comparable to the turbulent disk field strengths in the standard model if Maxwell stresses supply the dissipation torques. A field of this strength could easily be achieved by the accretion compression of a large scale stellar ($M \sim M_\odot$) or interstellar ($M \sim 10^9 M_\odot$) magnetic field.

The large-scale field is convected inward with the accretion flow, and must ultimately fall in toward the black hole. The electrodynamic interaction of this magnetic field with the black hole can also power an outflowing jet.

5.0 BLACK HOLE ELECTRODYNAMICS

In the late 1960's two distinct power sources were suggested to explain high energy astrophysical objects - accretion onto a compact gravitational object, and the extraction of the rotational kinetic energy of a compact object by electromagnetic torques. The accretion process was associated with galactic x-ray sources (Prendergast and Burbidge, 1968) and quasars (Lynden-Bell, 1969). Gold's (1969) hypothesis that pulsars are rapidly rotating, highly magnetized neutron stars remains unchallenged. Morrison (1969) generalized the pulsar concept to any object - the spinar - whose evolution, including slow gravitational collapse, is driven by the electromagnetic loss of rotational kinetic energy; supermassive spinars can have quasar-scale luminosities.

Blandford (1976) and Lovelace (1976) realized that an accretion disk which possessed a large-scale magnetic field would be a flat spinar. In direct analogy with the Goldreich-Julian (1969) theory of pulsars, Blandford (1976) developed a disk model in which the electromagnetic torque was exerted by a nearly charge-separated, force-free magnetosphere which extends beyond the light cylinder. The Blandford and Payne (1982) jet model is just the hydromagnetic version of Blandford's (1976) force-free model; when the Alfven speed is much less than the speed of light, plasma inertia influences the field and flow structure well inside the light cylinder.

In a fundamental conceptual advance, Blandford and Znajek (1977) demonstrated that the rotational energy or the reducible mass of a Kerr black hole could be extracted by electromagnetic torques if the hole was threaded by a large-scale magnetic field; black holes could also be spinars. Recently, MacDonald and Thorne (1982) reformulated the

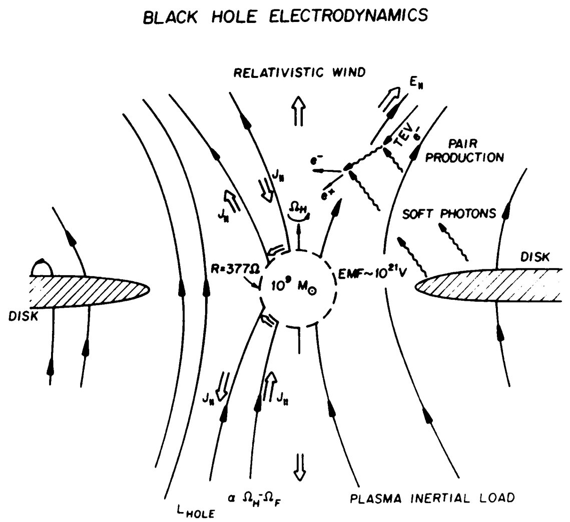

Figure 4. Coupling of black holes and accretion disks.

Blandford-Znajek concept into an absolute space - universal time description of General Relativity which, by using the familiar concepts of nonrelativistic hydromagnetics, completely illuminates the analogy between black hole and pulsar electrodynamics. In this section we pre sent the MacDonald and Thorne (1982) arguments which lead to the conclusion that relativistic jets can be powered by the rotational energy stored in Kerr black holes.

5.a. Magnetized Black Holes

Suppose that the plasma in the accretion disk is connected and frozen-in to a large-scale magnetic field. As the plasma falls into the black hole, the magnetic field becomes causally disconnected from and slips relative to the plasma. Due to the gravitational redshift, the magnetic field appears to concentrate just outside the event horizon (Figure 4). Price's (1972) "no hair" theorem states that the black hole cannot have an intrinsic field unless it carries a net charge, an unlikely possibility in the presence of external plasma. However, to an external observer, the flux concentration at the horizon will make the black hole appear to be threaded by a magnetic field.

Blandford and Znajek (1977), Znajek (1978), and Damour (1978) showed that the magnetized black hole could be mathematically represented as if it were a magnetized conducting body which satisfied conventional hydromagnetic boundary conditions at the event horizon: 1) electric field lines terminate on surface charges; 2) the tangential component of the magnetic field is shielded from the hole's interior by surface currents; 3) the resistivity of the horizon is $R_H = 4\pi/c =$ $= 377$ ohms; and 4) the tangential electric field (E_H) is related to the surface current density (j_H) by Ohm's law $E_H = R_H j_H$. In the absence of the highly conducting disk, the finite resistance of the horizon would cause the magnetic field to "diffuse" away at the speed of light. Thus the disk and its magnetic field are essential to maintaining a magnetized black hole.

5.b. Pulsar - Black Hole Electrodynamics

Suppose that the black hole's magnetosphere is initially a vacuum, is symmetric about the rotation axis, and rotates with the hole's angular frequency Ω_H; of course this aligned-rotator configuration does not emit magnetic dipole radiation. As in the Goldreich-Julian (1969) analysis for pulsars, the external solution to the vacuum Maxwell equations has a large electric field ($E_\parallel$) component which is parallel to the magnetic field; the "source" of $E_\parallel$ is the Goldreich-Julian space charge which polarizes the vacuum in the corotating frame. For pulsars, $E_\parallel$ causes field emission from the neutron star's surface. The extracted particles are accelerated to relativistic energies and radiate curvature photons by the synchrotron process. As the photons cross field lines, their perpendicular momentum is converted into electron-positron pairs in a quantum electrodynamic cascade. The dense pair plasma shorts-out $E_\parallel$, limiting the acceleration to a narrow vacuum gap (Sturrock, 1971; Ruderman and Southerland, 1975; Arons, 1979).

Blandford and Znajek (1977) suggested that $E_\parallel$ might also lead to pair creation on black hole field lines; for a 10^9 M , maximally rotating Kerr hole, a total emf of 10^{21} volts is possible, making acceleration to relativistic energies very likely. If a "seed" electron enters the vacuum $E_\parallel$ region, after acceleration it will Compton scatter soft disk photons to γ-ray energies. The γ-rays can then collide with soft photons to produce pairs provided that the product of the two photon energies exceeds $(mc^2)^2$.

Blandford and Znajek (1977) assumed that the pair plasma density would increase until the force-free condition $\eta \underline{E} + 1/c\, \underline{J} \times \underline{B} = 0$ was satisfied throughout most of the magnetosphere; η is the charge density. The parallel electric field would occur in a narrow gap somewhere between the horizon and the light cylinder. Some of the pair plasma falls into the black hole, while some flows outward along the expanding magnetospheric field lines. Near the light cylinder, the plasma inertially loads the magnetic field, causing the force-free condition to be violated and driving a system of field-aligned currents. The currents close by flow across the magnetic field at the event horizon and in the load region. The $\underline{J} \times \underline{B}$ stress in the load exerts a torque which attempts, but of course fails, to maintain the angular frequency of the plasma and magnetic field (Ω_F) equal to that of the hole. The plasma's inertia bends the poloidal lines of force into the toroidal direction, so that angular momentum flows outward. The integrated $\underline{J} \times \underline{B}$ stress in the horizon exerts a spin-down torque, thus extracting the rotational energy of the black hole.

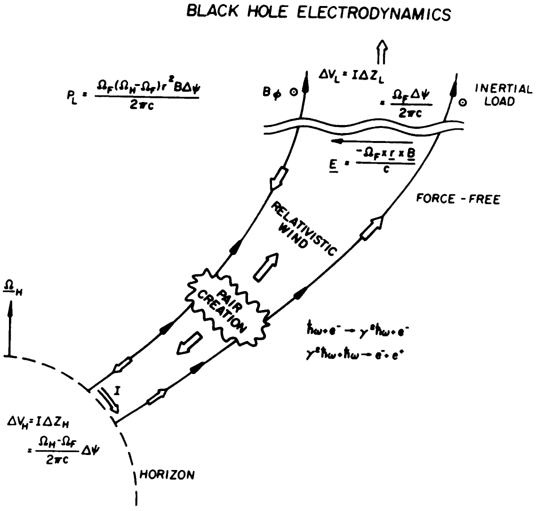

Figure 5. Equivalent circuit for black hole-magnetospheric coupling.

5.c. Equivalent Electrodynamic Circuit

MacDonald and Thorne (1982) developed a circuit analogy (first suggested by Blandford, 1979) which summarizes the magnetosphere-black hole interaction, and permits a simple estimate of the power extracted from the hole. The great advantage of the absolute space - universal time formulation is that we can forget about General Relativity, curved space-time, and the dragging of inertial frames, and treat the electrodynamics exactly as in Goldreich-Julian (1969).

Consider a ring of magnetic flux $\Delta\psi = 2\pi\ r\ B_H \Delta\theta$ which threads the black hole at a perpendicular distance r from the rotation axis and over a latitude $\Delta\theta$; ignore the fact that r is really the magnitude of a Killing vector. Just outside the horizon the magnetic field rotates at $\Omega_F \neq \Omega_H$ so that the electric field satisfies $c\underline{E} = -\underline{\Omega}_F \times \underline{r} \times \underline{B}_H$. In the horizon, frame dragging forces the magnetic field to corotate with the hole; hence in the "rest frame" of the hole, there is an electric field $c\underline{E}_H = (\underline{\Omega}_H - \underline{\Omega}_F) \times \underline{r} \times \underline{B}_H$ which is directed opposite to $\underline{E}$ if $\Omega_H > \Omega_F$ as expected from the inertial load. Across the latitude strip $\Delta\theta$, $\underline{E}_H$ provides a voltage drop $\Delta V_H = (\Omega_H - \Omega_F)\Delta\psi/2\pi c$ (Figure 5). The impedance across the strip is $\Delta Z_H = R_H \Delta\theta/2\pi r = \Delta\psi/\pi c r^2 B_H$. By Ohm's law, the total current (I) through the horizon satisfies $I\Delta Z_H = \Delta V_H$, yielding $I = 1/2(\Omega_H - \Omega_F) r^2 B_H$.

As a simple model of the load region, MacDonald and Thorne (1982) assume that the angular velocity of the plasma and field lines drops rapidly from Ω_F to zero. Hence the electric field in the magnetosphere $c\underline{E} = -\underline{\Omega}_F \times \underline{r} \times \underline{B}$ appears across the load, and corresponds to a voltage drop $\Delta V_L = \Omega_F \Delta\psi/2\pi c$. If the impedance of the load region is ΔZ_L, the current through the load satisfies $I\ \Delta Z_L = \Delta V_L$, and I must equal the

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current through the horizon; ΔZ_L is unspecified, but is presumably related to the plasma's inertia.

From the above circuit analysis, we can determine the power which is dissipated in the load as $\Delta P_L = I^2 \Delta Z_L$ or

$$\Delta P_L = \frac{(\Omega_H - \Omega_F)\Omega_F \, r^2 B_H \Delta\psi}{4\pi c}$$

Properly interpreted, the expression for ΔP_L corresponds to the exact General Relativistic result. The power vanishes if $\Omega_H = \Omega_F$ so that the magnetosphere exactly corotates with the black hole, and maximizes for $\Omega_F = 1/2\ \Omega_H$, which corresponds to matched impedances ($\Delta Z_H = Z_L$) between the hole and the load.

As a rough estimate for the total power extracted from the black hole, we take $\Delta\psi \sim \pi r^2 B_H$ and set $r \sim GM/c^2$ amd $r\Omega_H \sim c$ (maximally rotating Kerr hole) to obtain

$$L_H \sim \frac{(\Omega_H - \Omega_F)}{\Omega_F} \left(\frac{\Omega_F}{\Omega_H}\right)^2 \frac{B^2 G^2 M^2}{4c^3} \sim 4 \times 10^{45} B_4{}^2 M_9{}^2 \text{ ergs/sec}$$

with $\Omega_F = 1/2\ \Omega_H$, $B_4 = B/10^4$ Gauss and $M_9 = M/10^9 M$. Hence the luminosity of the black hole is comparable to the total power observed in the jets of extragalactic radio sources and even approaches the luminosities of quasars and active galactic nuclei. The partitioning of L_H between plasma thermal and flow energy, radiation, and even hydromagnetic waves depends on the (unknown) dissipation precesses in the load region. Presumably the plasma will be accelerated to relativistic energies, and at least part of L_H will flow outward in a jet along the rotation axis.

Although the black hole's luminosity does not depend on the structure of the accretion disk, except through its high conductivity, the infalling plasma can resupply the angular momentum lost by the hole and can add new magnetic flux to the hole. An interesting possibility would be if the disk possessed regions of opposite magnetic polarity to that of the black hole. Since all velocities near the hole are comparable to the light speed, upon falling in reconnection would rapidly dissipate the anti-parallel fields. The tension of the newly reconnected field lines would accelerate twin relativistic jets parallel to the rotation axis. These reconnection events would probably be sporadic, and may contribute to the optical, x-ray, or radio outbursts observed in quasars.

6.0 Discussion

In such a brief review we cannot do justice to the extensive efforts which have established accretion disk and black hole electrodynamics as a fundamental problem in high energy astrophysics. Although, in the astrophysical lexicon, the present models are all reasonable, they only represent a hydromagnetic skeleton which outlines the possible dynamics occurring in galactic x-ray sources, quasars, and active galactic nuclei. In order for those hydromagnetic models to acquire explicative and predictive powers, local plasma physical dissipation processes must be incorporated into the hydromagnetic flow solutions. As accretion disk and black hole models evolve, plasma physics advances in solar physics, planetary magnetospheres, and pulsars will continue to provide guidance and inspiration.

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DISCUSSION

Benford: Dissipation and particle production near black holes are, of course, messy problems. But can't we learn about the final wound-up field, B_o, by simply extrapolating the confining field in jets back from the extragalactic region?

Coroniti: Yes, this may be the best way.

Sturrock: As I understand it, a jet produced by a black hole will comprise an electron-positron plasma. Can this low density plasma be reconciled with the dynamics of fully developed jets?

Coroniti: The usual argument is that the positronic jet picks up hydrogenic as it propagates into the galactic and extragalactic environment, and is converted into an ordinary jet. I do not believe that this pick-up has been demonstrated, nor that the jet could survive the inertial drag.

Henriksen: The sub-parsec scale jets are the energy sources of the VLA jets, but they are probably not in the same physical state. The small scale jets will have highly ordered, high specific energy which will be degraded by mixing and entrainment as the jet expands to the VLA scale. See e.g. Henriksen, Bridle, and Chen, 1982.

D. Smith: In the case of pulsars, we have been working for years with very detailed observations, but there is no agreement on the dissipation mechanisms. Using this as a gauge, is there any hope that we could be in any better shape in the case of quasars?

Coroniti: Hope springs eternal in an astronomer's breast. In both cases, we need to construct a good numerical simulation in order to self-conistently study the basic dissipation. Even if the simulations are not realistic (parametrically), they will still be essential for sorting out which of the possible plasma dissipation processes are important.

Vasyliunas: Is the 377 ohm effective resistivity of the black hole low enough to enforce corotation of the field lines, as required for application of pulsar magnetospheric physics?

Coroniti: If the inertial load is large enough, the field lines would be decoupled from the hole. However, once decoupled, the source of emf which ultimately drives the pair production goes away, thus decreasing the inertial load. Hence the system may be self-regulating.

Uchida: You talked about the magnetic field hung up on the event horizon while the mass is absorbed into the black hole. Is the process of this defreezing of the magnetic field from the mass explored?

Coroniti: Frame-dragging at the Kerr hole's horizon causally decouples the plasma from the magnetic field. The same effect leads to the finite, free-space resistivity of the horizon.

Vlahos: Are there impulsive events observed in accretion coronae? What are the scales of the variability of the emitted energy?

Coroniti: Yes, both Cygnus X-1 and many quasars exhibit temporal variability in their x-ray and (for quasars) optical light curves; quasars also have radio outbursts. The smallest time scales are comparable to the light time across the Schwarzschild radius. Whether the observed variability is due to unsteady coronal processes or a major instability in the disk structure is not known.

A. Ray: How collimated are the jets from the black holes?

Coroniti: I do not know whether this analysis has been done yet. Sterling Phinney is attempting to calculate the jet's properties, so I suppose that the degree of collimation would come out of his relativistic wind analysis.

Davila: Is there any way to directly observe the rotation of the black holes at the center of these disks, perhaps gravity wave radiation for instance?

Coroniti: I do not think that the gravitational radiation is detectable.

Gilden: In accretion disks around black holes the efficiency of gravitational radiation is only 2%, while the photon efficiency may be $\sim$ 40%.

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