

STRUCTURED CORONAE OF ACCRETION DISKS

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ABSTRACT

A model for the fluctuating hard component of intense cosmic X-ray sources (such as Cyg X-1) is developed, based upon the amplification of magnetic fields by convective motions and differential rotation within a hot ($T \gtrsim 10^8$ K) accretion disk. Field reconnection within the inner portion of the disk is shown to be ineffective in limiting field amplification; magnetic fields may therefore attain strengths comparable to the equipartition value, leading to their emergence via buoyancy in the form of looplike structures and resulting in a very hot ($T > 10^8$ K) magnetically confined, structured corona analogous to the observed structure of the solar corona. The energy balance of these loop structures is examined, and it is shown that the disk soft X-ray luminosity determines the predominant energy loss mechanism in loops: at low disk luminosities, thermal bremsstrahlung from these loops dominates and contributes a steady, shot-noise-like hard X-ray component. At high disk luminosities the emerging loops are Compton-cooled; the soft X-ray flux from the disk is Comptonized by the emerged loops, forming a transient, flarelike hard X-ray component.

Subject headings: hydromagnetics — stars: accretion — stars: coronae — X-rays: binaries

I. INTRODUCTION

Observations of compact X-ray sources thought to be associated with close binary systems have led in the past decade to the development of theoretical models which account for the X-ray emission by invoking accretion of gas onto the compact, collapsed component (see Blumenthal and Tucker 1974). The most detailed observational and theoretical work has been done for the X-ray source Cyg X-1 (see Oda 1977; Eardley *et al.* 1978), and for this reason we shall focus upon this source in the following discussion. The basic observational results we shall address, whose details are to be found in the above-mentioned surveys, are as follows:

a) Both optical and X-ray observations show periodic variability characteristic of a binary system; in particular, the X-ray source has been identified with the spectroscopic binary HDE 226868.

b) The energy spectrum of X-ray emission shows evidence for the presence of two components. The hard ($E > 10$ keV) component is always present, whereas the soft ($E < 10$ keV) component has been reliably detected only during Cyg X-1's "high state." There is weak evidence that the intensity of the hard component is anticorrelated with the intensity of the soft component.

c) The X-ray emission is known to be highly variable, on time scales ranging over 10 orders of magnitude, down to fluctuations at the millisecond level. Monte Carlo simulation has shown much of the variability to be characteristic of that associated with shot noise.

The thrust of our discussion will be to develop a coherent model relating the two X-ray components and their temporal variability; as part of this development we suggest a reason why Cyg X-1 may be distinguished from other binary X-ray sources. We shall not dwell upon other proposed models for Cyg X-1 except where directly relevant to our discussion.

It is now thought that the observed hard X-ray component originates from a hot ($T_e \gtrsim 5 \times 10^8$ K), optically thin plasma near the compact soft X-ray source, the latter most likely an accretion disk surrounding a gravitationally collapsed object (cf. Eardley *et al.* 1978). A wide variety of specific models have been proposed to account for these observations; most recently it has been suggested that the hard component derives from "Comptonization" of soft X-rays in a uniform hot corona surrounding the cooler, soft X-ray emitting accretion disk (Liang and Price 1977, Bisnovatyi-Kogan and Blinnikov 1977). This suggestion, to some extent motivated by the fact that standard accretion disk models are convectively unstable (cf. Shakura and Sunyaev 1973), is based upon an analogy with the solar corona, which is conventionally thought to be fairly homogeneous, and

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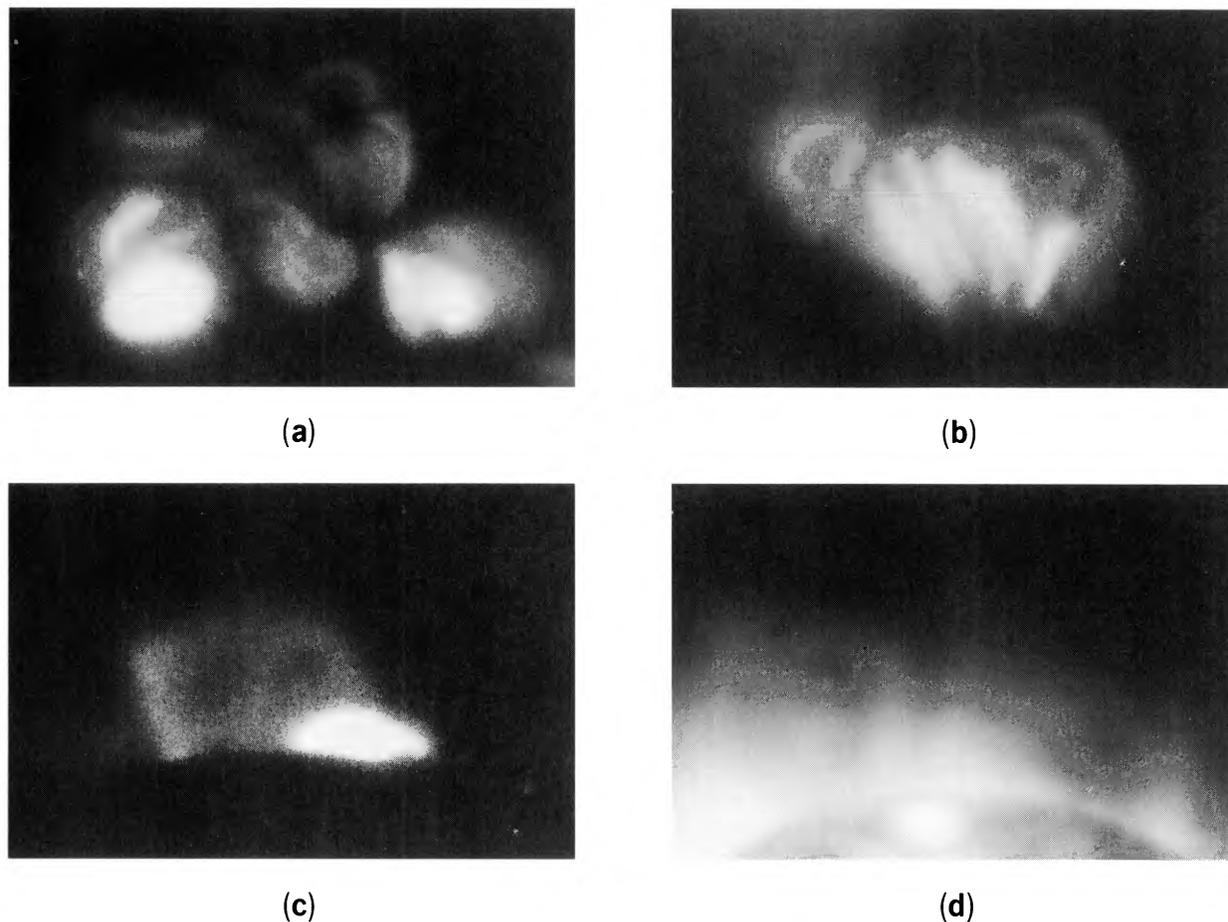


FIG. 1.—Soft (2–32, 44–54 Å) X-ray images of a variety of typical solar coronal regions obtained by the HCO/AS&E S-054 X-ray telescope on board *Skylab*: (a) emerging small active regions; (b) evolving large active region; (c) flare loop seen on solar limb; (d) large loop structures in “Quiet Sun.” The X-ray emission derives primarily from closed, looplike structures defined by coronal magnetic fields which emerge from the solar convective zone. Although the plasma regime envisaged here for the accretion disk corona is vastly different from that of the Sun, its geometric structure is likely to be quite similar to that pictured because in both cases the low- β coronal plasma is dynamically controlled by magnetic fields whose footpoints are embedded in a convectively unstable, $\beta \sim 1$ plasma.

heated by acoustic flux generated by turbulent convection motions in the outer solar envelope (cf. Athay 1976). In these models magnetic fields are invoked primarily to provide an efficient mechanism for angular momentum transfer within the accretion disk (Eardley and Lightman 1975) and are not involved directly in the maintenance of the hot corona.

This argument by analogy is, however, vitiated by recent EUV and X-ray observations of the solar corona, particularly by the instruments carried on board *Skylab*; these have shown the solar corona to be composed of a variety of complex loop structures generally associated with underlying photospheric and chromospheric magnetic field complexes, in contrast with the homogeneous geometry previously assumed. The bulk of solar X-ray emission derives instead from magnetically confined, looplike volumes; volumes which appear open to the interplanetary medium show relatively low X-ray brightness and correlate well with high-speed solar wind streams (Fig. 1). These observations suggest that spatial confinement is the key to obtaining a hot, dense corona (see Vaiana and Rosner 1978); in particular, it has been argued that the necessary element for the formation of a strong corona appears to be the presence of a confining magnetic field.

Magnetic fields have, of course, long been considered an important element in the dynamics of the accretion disk, primarily as a mechanism for supplying internal stresses required for efficient angular momentum transfer (Eardley and Lightman 1975, Ichimaru 1977). These calculations assume that the level of field fluctuations is primarily determined by dissipative processes *within* the disk itself. We shall show explicitly, however, that even under the most favorable conditions of field annihilation, internal dissipative processes are ineffective at limiting the growth of magnetic field fluctuations. This consideration leads us to reject these previous accretion disk models invoking relatively weak ($\beta \gg 1$) magnetic fields and suggests instead a new model in which the magnetic field plays a more central role. In this model strong magnetic fields ($\sim 10^8$ gauss) can be realistically generated within the inner portion of an accretion disk by the joint action of plasma thermal convection and differential rotation along Keplerian orbits. Field amplification will then be limited by nonlinear effects; as a consequence of buoyancy, magnetic flux will be expelled from the disk, leading to an accretion disk corona consisting of many magnetic loops where the energy is stored. The magnetic fields in these loops can provide an energy source for plasma heating, leading to a state in which the loops contain very hot, relatively low-density plasma as compared to the disk itself. Hard X-ray emission is the consequent result, due both to bremsstrahlung and to Comptonization of soft X-rays from the cooler underlying disk; which of these two loss processes dominates will be shown to be determined by the accretion disk luminosity, and both the luminosity and the time scale of fluctuations of the ensuing hard X-ray component are demonstrated to be consistent with observations.

II. MAGNETIC FIELD GENERATION AND THE FORMATION OF CORONAL LOOPS

Magnetic field generation due to differential motion of conductive media is usually described in the MHD limit with the aid of the induction equation for the magnetic field (cf. Parker 1955)

$$\frac{\partial \mathbf{B}}{\partial t} = \text{curl} [\mathbf{v} \times \mathbf{B}] + \frac{c^2}{4\pi\sigma} \nabla^2 \mathbf{B}, \quad (1)$$

where \mathbf{B} is the magnetic field induction, \mathbf{v} the velocity of the fluid, and σ the electrical conductivity. For typical parameters of the plasma in the accretion disk (plasma density $n \approx 10^{22} \text{ cm}^{-3}$, plasma temperature $T \sim 10^7 \text{ K}$; cf. Bisnovatyi-Kogan and Blinnikov 1977) the magnetic Reynolds number Re_m is large,

$$\text{Re}_m \approx 10^{12}; \quad (2)$$

therefore, magnetic field evolution via the plasma differential motion is well described by equation (1) if we neglect the (classical) collisional dissipation of the magnetic field implied by the last term on the right-hand side of equation (1). In the case of Keplerian differential rotation, equation (1) then describes (assuming an axisymmetric geometry) the amplification of the azimuthal magnetic field B_ϕ in the presence of a seed radial magnetic field B_r :

$$\frac{\partial B_\phi}{\partial t} = r \frac{\partial \Omega}{\partial r} \cdot B_r, \quad (3)$$

where $\Omega = (GM/r^3)^{1/2}$ is the Keplerian angular velocity. It has been proposed that reconnection of the azimuthal magnetic field then accounts for the generation of the radial magnetic field and that the magnetic field stress

$$t_{\phi r} = (1/4\pi) B_\phi B_r, \quad (4)$$

provides the primary mechanism for angular momentum transfer within the accretion disk (Eardley and Lightman 1975). We now show that even the fastest (Petschek-type) reconnection mechanism is insufficiently rapid to develop effectively in the inner portion of the accretion disk and that the buildup of magnetic fields within the disk is instead limited by nonlinear effects related to convection.

To this purpose we consider the thermal convection of the highly conducting plasma within the accretion disk. Since convection takes place primarily perpendicular to the plane of the disk, we shall assume differential rotation to remain the dominant mechanism for azimuthal magnetic field (B_ϕ) generation; the generation of the remaining field components (B_r and B_z) will then be dominated by convection-mediated effects. For convection cells whose aspect ratio is $\sim O(1)$, the radial magnetic field spatial scale will then be of the order of the convective cell size z_0 , equal to the half-thickness of the disk; in that case the right-hand side of equation (3) can be replaced by the

expression $z_0 \partial \Omega / \partial r \cdot B_r$, and equation (3) is modified to read

$$\dot{B}_\phi \approx -\frac{3}{2} \frac{z_0}{r} \Omega \cdot B_r. \quad (3')$$

To describe the generation of the radial (B_r) and vertical (B_z) magnetic field components, we invoke magnetic flux conservation within the magnetic field loop:

$$l B_r \approx z_0 B_\phi, \quad (5)$$

where l is the length of the loop in the azimuthal direction; furthermore, mass conservation applied to the convective motion and differential rotation of the loop yields the relation

$$\begin{aligned} \rho v l &\approx \rho \bar{\Delta} v_\phi z_0 \\ &\approx \rho \cdot \frac{1}{2} \Omega z_0 \cdot z_0, \end{aligned} \quad (6)$$

where v is the plasma convection velocity, assumed to be slower than the differential rotation speed v_ϕ . Figure 2 shows the structure of the resulting magnetic field. Combining equations (3'), (5), and (6), we finally obtain the rate of magnetic field generation:

$$\dot{B}_\phi = \tau_G^{-1} B_\phi, \quad \tau_G^{-1} = 3v/r. \quad (7)$$

We see that the magnetic field grows exponentially until either reconnection or nonlinear effects stop the growth; similar arguments for exponential magnetic field growth are given by Shields and Wheeler (1976). Note that convection is essential to the occurrence of exponential field amplification because it couples azimuthal fields generated by differential rotation back to the remaining field components from which the azimuthal field is produced.

Let us first estimate the limit imposed by nonlinear effects. We can expect magnetic field growth to saturate when the magnetic field tension suppresses convection, the primary source for magnetic field generation. Thus field amplification is bounded by the constraint

$$B_\phi^2 / 4\pi R_c \lesssim \rho v^2 / z_0, \quad (8)$$

where R_c is the radius of curvature of the magnetic field lines in the vertical (or radial) direction. For elongated loops with $l \gtrsim z_0$,

$$R_c \approx l^2 / z_0. \quad (9)$$

In order to estimate the length of a typical loop, we combine equation (6) with the sound speed c_s [$\approx (v_\phi/r) z_0$; cf. Shakura and Sunyaev 1973] to obtain

$$l \approx \frac{1}{2} (c_s/v) z_0 \approx z_0 / \alpha^{1/3}, \quad (10)$$

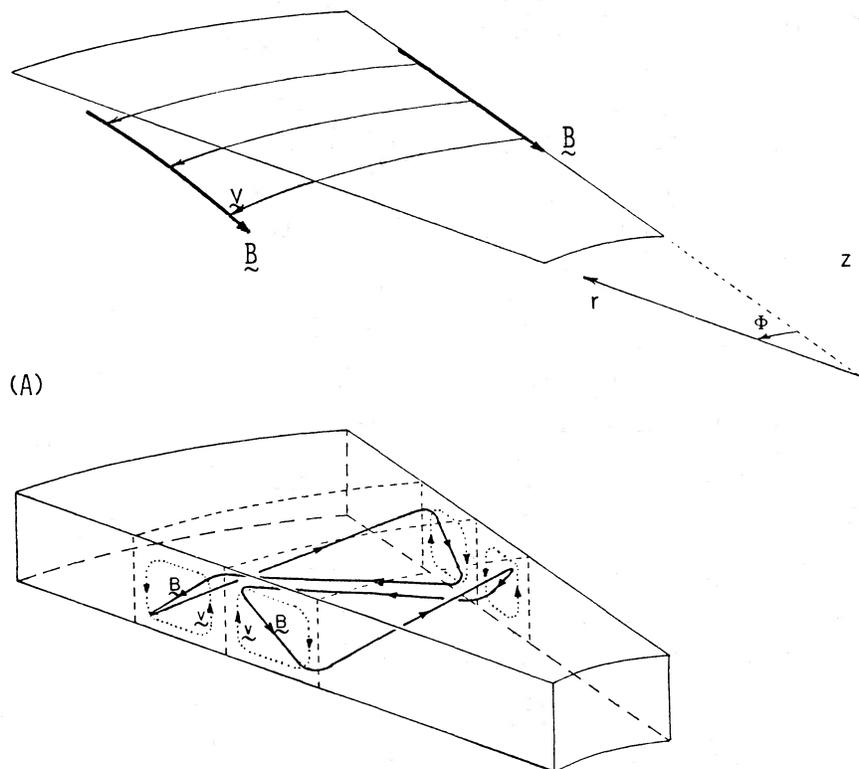


FIG. 2.—(a) Schematic presentation of azimuthal magnetic field (B_ϕ) generation via differential rotation within the accretion disk (cf. eq. [3]). (b) Schematic configuration of the action of convection within the accretion disk upon the internal magnetic field (flow indicated by short arrows), resulting in the generation of additional z and r field components from the azimuthal magnetic field; this process provides the feedback necessary for exponential growth of the magnetic field within the disk (cf. eq. [7]).

where α is a small parameter relating the plasma stress $t_{\phi r}$ to the pressure p (e.g., $t_{\phi r} = \alpha p$), and where we have used the estimate of Bisnovaty-Kogan and Blinnikov (1977) for the convection velocity $v[\approx \alpha^{1/3}c_s]$. Then using equations (6) and (8)–(10), we obtain the limiting strength of the magnetic field:

$$B_{\phi}^2/4\pi \approx \rho c_s^2. \quad (11)$$

Thus, the magnetic field pressure becomes comparable to the ambient gas pressure, as one would expect from simple energy equipartition arguments; however, our conclusion cannot be based upon equipartition alone because the latter implicitly assumes that field reconnection does not come into play.

Our argument therefore requires an explicit comparison of the time scales associated with field growth and reconnection, and a demonstration that—for some significant portion of the accretion disk—no reconnection occurs which may stop magnetic field growth before the nonlinear limit is reached. In a plasma with $n \approx 10^{22} \text{ cm}^{-3}$, $T = 10^7 \text{ K}$, $B \approx 7.5 \times 10^7 \text{ gauss}$, $l \sim 3 \times 10^6 \text{ cm}$, the Reynolds number is $\text{Re}_m \approx 10^{12}$, and the upper limit for the Petschek-type reconnection time is (Priest and Soward 1976):

$$\tau_R^{-1} = \frac{\pi v_A l^{-1}}{4(\ln \text{Re}_m + 0.74)} \approx 2.1 \times 10^{-2} \frac{v}{z_0}, \quad (12)$$

where $v_A = B/(4\pi\rho)^{1/2}$ is the Alfvén speed and is of the order of the sound speed (eq. [11]). Reconnection is therefore unimportant if $\tau_R/\tau_G > 1$, or (using eqs. [7] and [12])

$$r/z_0 \lesssim 140. \quad (13)$$

Now, within the inner portion of the accretion disk $z_0 = 3r_0(L/L_c)[1 - (r_0/r)^{1/2}]$ (Bisnovaty-Kogan and Blinnikov 1977); therefore, for $1.2r_0 \lesssim r \lesssim 10r_0$, we obtain

$$2.25 \lesssim (r/z_0)(L/L_c) \lesssim 5, \quad (14)$$

where L/L_c is the luminosity in units of the Eddington limit $L_c = 4\pi cGMm_p/\sigma_T$. Comparing equations (13) and (14), we see that for $L \sim 0.1L_c$ —in the most radiative part of the disk ($r \approx 3r_0$)—equation (13) is easily satisfied and therefore reconnection is unable to suppress magnetic field generation. In that case the argument leading to equation (11) applies. Magnetic flux tubes with such strong magnetic field strengths will contain less plasma than their ambient surroundings in order to maintain pressure balance; therefore they are subject to buoyancy forces and will penetrate the accretion disk to form “coronal” loops. As soon as these field lines emerge from the disk, reconnection becomes faster because the coronal density is much lower. Reconnection can therefore provide an efficient mechanism for plasma heating in the emerged loops; this possibility is explored below (§§ III and IV). A sketch of the envisaged geometry is shown in Figure 3.

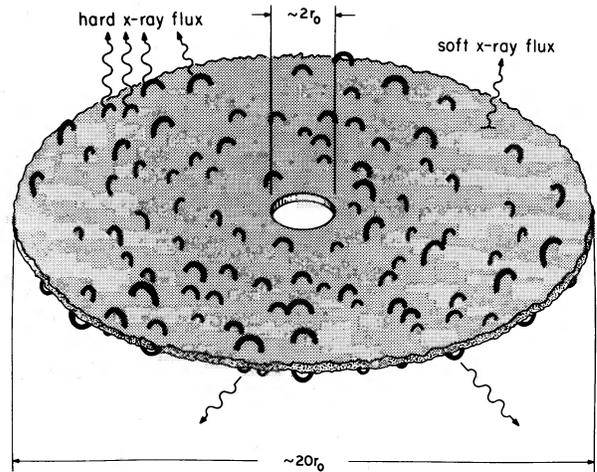


FIG. 3.—Schematic drawing of the inner accretion disk coronal geometry, with $r_0 \equiv 6GM/c^2 \sim 10^7 \text{ cm}$ in the case of Cyg X-1; only the inner portion of the disk is shown. The soft X-ray component derives from the relatively cool disk (including the outer portion not shown here) while the hard X-ray component is emitted by the ensemble of hot ($T_e \gtrsim 5 \times 10^8 \text{ K}$) plasma loop structures which have emerged from the inner disk. The length of typical loop structures is of the order of 10^6 cm , but can be expected to vary considerably as the loops evolve and their magnetic fields decay.

III. PLASMA THERMAL BALANCE IN CORONAL LOOPS: LOW LUMINOSITY STATE OF THE ACCRETION DISK

We now consider the thermal balance of plasma entrained within the magnetic loops formed above the accretion disk (§ II). At the temperatures of interest ($\gtrsim 10^8 \text{ K}$), and for the loop scale sizes involved ($\lesssim 10^6 \text{ cm}$), typical dynamic relaxation times are less than 10^{-2} s ; furthermore, the pressure scale height substantially exceeds the relevant physical dimensions. Therefore, on the time scale of $\sim 1 \text{ s}$, we can reasonably expect quasi-stationary conditions, with $p \approx$ constant along the magnetic field. This precludes modeling of the impulsive phase of loop eruption from the accretion disk, during which the loop is filled with hot plasma. In the present case, thermal balance is governed by the energy conservation equation (Rosner, Tucker, and Vaiana 1978)

$$E_H - \frac{\partial}{\partial s} F - E_R = 0, \quad (15)$$

where F is the thermal conduction flux, E_H is the heating rate, E_R is the relevant radiation rate (see below), and s is the coordinate along the field lines. We assume that Compton cooling (Illarionov and Sunyaev 1972; Felten and Rees 1972) of loops is relatively ineffective in the low-luminosity state; this assumption will be verified below.

For sufficiently hot plasma in the loop ($T \gtrsim 5 \times 10^8 \text{ K}$), the stationary thermal flux is strongly inhibited by the ion-sound instability of the thermal flux (Forslund 1970); calculations of the thermal flux based upon a quasi-linear theory of the ion-sound instability

yield the following expression for F (Galeev *et al.* 1979):

$$F = -\kappa_{QL}(kT_e/m_p)^{1/2}nkT_e \text{ sign } dT/ds \text{ ergs cm}^{-2} \text{ s}^{-1}. \quad (16)$$

The numerical coefficient κ_{QL} depends strongly upon the ratio of the ion and electron temperatures. If the loop heating is due to current dissipation by means of reconnection of fields within loops (Galeev *et al.* 1978), then

$$\pi R^2 l E_H = [\pi v_A / 4 (\ln \text{Re}_m + 0.74)] \frac{B_J^2}{4\pi} 2\pi R l$$

or

$$E_H = 1.4 \times 10^{-2} v_{AJ} (B_J^2 / 4\pi) (2/R), \quad (17)$$

where R is the minor radius of the coronal loop, B_J the loop current-generated magnetic field, and v_{AJ} the corresponding Alfvén speed. This mechanism preferentially heats electrons; we therefore assume that $T_e \gg T_i$. In that case the numerical coefficient κ_{QL} in the expression for the conductive heat flux can be calculated analytically by neglecting, for simplicity, ion-ion collisions; this gives $\kappa_{QL} \approx 19$. However, ion-ion collisions play an important role as electron-ion collisions and are known to reduce the classical heat conduction coefficient by a factor of 4 (Braginskii 1963). Therefore, we can expect that the quasi-linear limit for the heat flux, given by equation (16), will be reduced by a similar factor if ion-ion collisions are properly taken into account. In our numerical estimate we set $\kappa_{QL} \approx 5$, in accordance with experimental measurements of the limiting heat flux for the case of laser plasma heating (Gray, Kilkenney, and White 1977).

Radiative loss mechanisms to be considered include bremsstrahlung, and cyclotron or synchrotron radiation; we postpone discussion of Compton cooling until later, and note that cooling via e^-e^+ pair production is unimportant below $\sim 6 \times 10^9$ K. The relevant expressions for these losses are given by (cf. Novikov and Thorne 1973; Tucker 1975)

$$(E_R)_{\text{brems}} \approx 1.3 \times 10^5 p^2 T_e^{-3/2} \text{ ergs cm}^{-3} \text{ s}^{-1}, \quad (18a)$$

$$(E_R)_{\text{cyclo}} \approx 3.9 \times 10^{-9} p B^2 \text{ ergs cm}^{-3} \text{ s}^{-1}, \quad (18b)$$

$$(E_R)_{\text{synch}} \approx 2.7 \times 10^{-18} p T_e B^2 \text{ ergs cm}^{-3} \text{ s}^{-1}, \quad (18c)$$

where $p \approx nkT_e$, T_e [K] is the electron temperature, and B the loop magnetic field. If we define the parameter $\beta \equiv 8\pi p / B^2$ [$\sim O(1)$], we can judge the relative importance of these losses by computing their ratios; thus we obtain

$$\begin{aligned} \text{I. } (E_R)_{\text{brems}} / (E_R)_{\text{cyclo/synch}} \\ &\approx 20\beta [T_e / 10^7 \text{ K}]^{-3/2}, \quad T_e \ll 6 \times 10^9 \text{ K} \\ &\approx 0.06\beta [T_e / 10^9 \text{ K}]^{-5/2}, \quad T_e \gg 6 \times 10^9 \text{ K}. \end{aligned}$$

$$\text{II. } F / [(E_R)_{\text{brems}} l_H]$$

$$\approx 420 [n / 10^{17} \text{ cm}^{-3}]^{-1} [T_e / 10^9 \text{ K}] [l_H / 10^5 \text{ cm}]^{-1}.$$

$$\text{III. } F / [(E_R)_{\text{synch}} l_H] \approx 25\beta [n / 10^{17} \text{ cm}^{-3}]^{-1}$$

$$\times [T_e / 10^9 \text{ K}]^{-3/2} [l_H / 10^5 \text{ cm}]^{-1}.$$

In the above, F is the maximum thermal conductive flux in the loop (e.g., F where $\text{div } F = 0$; eq. [16]), and l_H is the half-length of the region where $\text{div } F > 0$, and hence thermal conduction acts as an energy loss mechanism (cf. eqs. [20]–[22] below). These ratios suggest that, although synchrotron losses dominate bremsstrahlung losses for loop temperatures above 10^9 K, the dominant loop energy loss mechanism at these temperatures is actually thermal conduction to the far cooler, underlying accretion disk. For this reason, we shall solve the thermal energy balance equation (15) separately in the heating-dominated (e.g., $|\text{div } F| \gg E_R$, or $E_H \gg E_R$) hot coronal portion and the radiation-dominated (e.g., $E_H \ll E_R$) transition region at the loop footpoints. In the latter region, the divergence of the thermal conductive flux supplies the energy required to balance local radiative losses; furthermore, because of the relatively low temperatures in this transition boundary region, bremsstrahlung losses dominate, and therefore E_R will be given by equation (18a). Note that the bulk of the loop radiative losses occurs in the transition region between the accretion disk and the loop “heating” region, via bremsstrahlung, and that losses due to other emission mechanisms—such as synchrotron emission from the heating region—constitute less than $\sim 10\%$ of the total loop radiative loss.

Thus, in the radiation-dominated region, which is located near the footpoints of the loop, E_H is neglected in equation (15); the temperature in this region then increases rapidly with height s :

$$T = [(S/k^2) / \kappa_{QL} (k/m_p)^{1/2}]^{1/2} (4ps)^{1/2},$$

$$S \equiv 2.4 \times 10^{-27}. \quad (19)$$

In the heating-dominated region, heat conduction can be considered as classical ($F = -10^{-6} T^{5/2} \partial T / \partial s$); this results in a comparatively low temperature gradient ($T \approx \text{constant}$). The transition from classical to inhibited (by ion-sound turbulence) conduction takes place at the point of maximum heat flux (or zero divergence of the heat flux). Thermal balance at this point requires

$$E_H l_H = F(T_R, p), \quad (20)$$

where T_R is defined by equation (19) at $s = l_R$, and l_R and l_H are the lengths of the radiation- and heating-dominated portions of the loop, respectively. The relation between the latter lengths is obtained from the equation

$$E_H = E_R(T_R), \quad (21)$$

which is satisfied at the boundary between the two parts of the loop. Combining equations (19)–(21), we obtain

$$l_H = 4l_R = (4/5)l, \quad (22)$$

$$T_e \approx 2.15\kappa_{QL}^{-1/2}(pl)^{1/2}, \quad (23)$$

$$p = 4.2 \times 10^{-4}\kappa_{QL}^{-3/5}E_H^{4/5}l^{3/5}, \quad (24)$$

where $p = nkT_e$ and T_e are the plasma pressure and electron temperature within the loop.

We now assume that the heating is due to reconnection-type current dissipation, as described by equation (17). Using the plasma β defined by the poloidal field B_J ,

$$\beta_J \equiv 8\pi p/B_J^2,$$

equation (24) can be rewritten in the form

$$\beta_J \approx 0.16(l/\kappa_{QL}R)^{2/3}. \quad (25)$$

In order to obtain numerical estimates, we use the plasma parameters calculated for a standard accretion disk by Bisnovatyi-Kogan and Blinnikov (1977):

$$\begin{aligned} B_0^2/4\pi &\approx [m_p c^2/(2^{1/2}\alpha\sigma_T r_0)](r_0/r)^{3/2} \\ &\approx 1.76 \times 10^{15}(M_\odot/\alpha M)(r_0/r)^{3/2}, \end{aligned} \quad (26)$$

$$z_0 = 3(L/L_c)r_0[1 - (r_0/r)^{1/2}], \quad (27)$$

where M is the mass of the massive central body (in Cyg X-1 probably a black hole with mass $M \approx 10 M_\odot$; Eardley *et al.* 1978), and L the luminosity of the disk (here $L \approx 0.1L_c$).¹

From stability arguments for the Kruskal-Shafranov mode in a toroidal pinch, the poloidal magnetic field cannot exceed the stability limit

$$B_J < (\pi z_0/2l)B, \quad (28)$$

where we estimated the major radius of the toroidal loop as $2l \approx \pi R_{\text{major}}$. With the above loop parameters, we find (from eqs. [23]–[24]) the loop plasma temperature and density:

$$\begin{aligned} \frac{kT_e}{m_e c^2} &\approx 7.07 \left(\frac{b^2}{\alpha}\right)^{1/2} \\ &\times \left\{ \frac{L}{L_c} \left(\frac{r_0}{r}\right)^{3/2} \left[1 - \left(\frac{r_0}{r}\right)^{1/2}\right] \right\}^{1/2} \left(\frac{l}{\kappa_{QL}R}\right)^{5/6}, \end{aligned} \quad (29)$$

$$\begin{aligned} n &\approx 2.42 \times 10^{19}(M_\odot/M)(r_0/r)^{3/2}(b^2/\alpha)^{1/2} \\ &\times [(L/L_c)(r_0/r)^{3/2}\{1 - (r_0/r)^{1/2}\}]^{-1/2} \\ &\times (l/\kappa_{QL}R)^{-1/6}, \end{aligned} \quad (30)$$

¹ Note that for disk luminosities of the order of $L \sim 0.06 L_c$ and lower, the accretion disk model based upon the assumption that the radiation pressure dominates the gas pressure is invalid. The above results are therefore to be considered as order-of-magnitude estimates only, with quantitatively more accurate results to be based upon more sophisticated models for the accretion disk.

where $b = B_J/B$. For low disk luminosities ($L \approx 0.02L_c$), the loops therefore attain temperatures of a few hundred keV.

It is now straightforward to calculate the Comptonization parameter y with the help of equations (29) and (30):

$$\begin{aligned} y &\equiv (4kT_e/m_e c^2)(n\sigma_T R)^2|_{R=z_0} \\ &\approx 5.44 \times 10^4 \left\{ \frac{b^2}{\alpha} \frac{L}{L_c} \left(\frac{r_0}{r}\right)^{3/2} \left[1 - \left(\frac{r_0}{r}\right)^{1/2}\right] \right\}^{3/2} \\ &\times (l/\kappa_{QL}R)^{1/2}, \end{aligned} \quad (31)$$

or $y < 1$ for typical values of the accretion disk parameters. Thus, while Comptonization cannot be ignored, equation (31) shows that its primary effect is to slightly modify the dominant bremsstrahlung spectrum from the coronal loop plasma (by increasing the energy of a small fraction of the bremsstrahlung photons), but not to affect the total magnitude of the bremsstrahlung losses (Illarionov and Sunyaev 1972; Felten and Rees 1972).

The above calculations allow one to estimate the total radiative emission from a typical accretion disk coronal loop; thus, the bremsstrahlung losses result in hard X-ray emission, with typical power

$$1.44 \times 10^{-2} v_{Aj}(B_J^2/4\pi)(2l)(2\pi R) \approx 3 \times 10^{35} \text{ ergs s}^{-1}. \quad (32)$$

The fluctuation time scale characterizing emission from a single loop is then of the order of the reconnection time

$$\tau = l/(1.44 \times 10^{-2} v_{Aj}) \approx 1 \text{ s}. \quad (33)$$

These results suggest that the stochastic emergence and heating of coronal loops is the physical process underlying the short-term temporal behavior of Cyg X-1, and provides a physical interpretation of the shot-noise model used by Terrell (1972; see also Boldt *et al.* 1975; Oda 1977) in the low disk-luminosity state.

IV. COMPTON COOLING OF CORONAL LOOPS: HIGH-LUMINOSITY STATE OF THE ACCRETION DISK

We now seek to determine the conditions under which electron cooling due to Compton processes (Illarionov and Sunyaev 1972; Felten and Rees 1972) dominates energy loss from coronal loops. We proceed by assuming the previously defined (§ III) heating-dominated region within a given coronal loop to be Compton-cooled (rather than by thermal conduction to the underlying disk) and, by comparing the resulting loop temperature with that obtained in the absence of Compton processes, determine the critical disk luminosity required to ensure the dominance of Compton cooling. The critical result to be obtained is that under conditions appropriate to the high-luminosity state of Cyg X-1, the coronal loops are primarily Compton-cooled.

Under the assumption of unsaturated Comptonization [$y \sim O(1)$; cf. Shapiro, Lightman, and Eardley 1976 and below], and using the heating function given by equation (17), the energy balance equation within the heating-dominated portion of the loop takes the form

$$1.4 \times 10^{-2} v_{\Lambda} \frac{B_j^2}{4\pi} \frac{2}{r} = \frac{4kT_e}{m_e c^2} n \sigma_T (\epsilon c), \quad (34)$$

where σ_T is the Thomson scattering cross-section and (ϵc) the soft X-ray radiation flux per unit area from one face of the disk; thus, we assume the accretion disk to be the dominant source of soft X-ray photons. If the total soft X-ray (< 10 keV) flux is of the same order as the total hard X-ray (> 10 keV) flux (see Oda 1977), then (ϵc) is given by (Novikov and Thorne 1973)

$$\epsilon c \approx (3L/8\pi)(r_0/r^3)[1 - (r_0/r)^{1/2}]. \quad (35)$$

Note that in the present model, the fraction of the total energy made available by the disk accretion which heats the corona is not a free parameter, but is observationally determined by the ratio of soft to hard X-ray emission. By setting $R = z_0$ and using equations (27) and (35) we can write the energy balance equation in the form

$$kT_e/m_e c^2 = 9.4 \times 10^{-3} (B_j^2/4\pi n m_p c^2)^{3/2} \times \{(L/L_c)(r_0/r)^{3/2}[1 - (r_0/r)^{1/2}]\}^{-2}. \quad (36)$$

To find the plasma density, we invoke energy balance within the radiation-dominated portion of the loop (e.g., at the loop footpoints) between bremsstrahlung radiation and the heat flux from the overlying heating region (cf. eq. [19]):

$$n = [5\kappa_{QL}(k/m_p)^{1/2}/4SI]kT_e. \quad (37)$$

With equations (36) and (37) we finally obtain the electron temperature within the coronal portion of the loop:

$$kT_e/m_e c^2 = 0.21(b^2/\alpha)^{3/5} \times \{(L/L_c)(r_0/r)^{3/2}[1 - (r_0/r)^{1/2}]\}^{-1/5} \times (l/\kappa_{QL}z_0)^{3/5}. \quad (38)$$

Comparing equation (38) with the equivalent result (e.g., eq. [29]) obtained in the absence of Compton processes, we conclude that Compton cooling results in lower loop temperatures, and hence dominates the coronal energy balance, for accretion disk luminosities L_* larger than that found from the relation

$$(L_*/L_c)(r_0/r)^{3/2}[1 - (r_0/r)^{1/2}] = 6.6 \times 10^{-3}(b^2/\alpha)^{1/7} \times (l/\kappa_{QL}z_0)^{-1/3}. \quad (39)$$

For loops situated in the region of maximum disk luminosity (i.e., for $r \sim 3r_0$) this requires that $L_* \gtrsim 0.055L_c$. Electron temperatures in the coronal loops

reach the maximum value for this luminosity independent of their position on the disk²

$$kT_e/m_e c^2 \approx 0.57(b^2/\alpha)^{4/7}(l/\kappa_{QL}z_0)^{8/15}. \quad (40)$$

The Comptonization parameter is then

$$y_* = 29.2(b^2/\alpha)^{12/7}(l/\kappa_{QL}z_0); \quad (41)$$

and, using equations (4), (25), and the relation $t_{\phi r} = (1/4\pi)B_\phi B_r$, is of order unity under these loop conditions. An emerging coronal loop is therefore strongly cooled—with a consequent collapse of the loop atmosphere—by the Comptonization of the soft X-ray flux emitted by the disk itself; this results in hard X-ray emission (e.g., the Comptonized soft X-ray photons) whenever a magnetic flux loop emerges from the disk. Whereas in the low disk-luminosity state, the intensity of the hard X-ray flux reflects the rate of energy deposition within loops, its intensity in the high state is determined primarily by the cooling rate of the entire loop plasma by the Compton processes. Consequently, the high disk-luminosity state should also be characterized by a highly variable hard X-ray component, whose intensity is anticorrelated with the underlying disk luminosity (cf. eqs. [37] and [38]); in addition, the frequency of large, flarelike transients should be higher than would be expected on the basis of a shot-noise model (as in the case of the low state).

V. CONCLUSION

We have shown that by taking into account (a) convection and (b) the ineffectiveness of magnetic field reconnection within the inner portion of an accretion disk, a structured magnetically confined accretion disk corona can form, consisting of many small-scale, extremely hot coronal loops. For the low-luminosity state of the accretion disk, plasma temperatures in these loops can be as high as $T_e \gtrsim 500$ keV; bremsstrahlung from the transition region of such loops gives rise to bursts of hard X-ray radiation, with maximum power in a single burst of the order of 3×10^{35} ergs s⁻¹ and a temporal duration of the order of ~ 1 s. Such fine-scale temporal variability has been observed for the Cyg X-1 source (see discussion by Eardley *et al.* 1978).

In the high-luminosity state of the disk, the coronal loops are efficiently cooled due to Compton electron scattering of the soft X-rays from the disk. This result is consistent with the observed dominance of the soft X-ray component during the high-luminosity state of the disk. Because the rate of magnetic field emergence can be expected to be relatively independent of the disk luminosity, the accretion disk corona will in this limit be characterized by a continual process of coronal loop emergence and consequent Compton cooling; a significant fraction of the soft X-rays emitted by the

² Note that when the loop temperature is sufficiently low, classical thermal conduction prevails, and the above results based upon plasma turbulence-inhibited heat conduction fail (cf. Galeev *et al.* 1978).

disk will thereby be Comptonized, leading to flarelike emission of hard X-rays, each burst the result of Compton-cooling of an entire emerged coronal loop. Our model does not encompass the mechanism responsible for the transitions of the soft X-ray component; however, among the several mechanisms proposed, changes in the accretion rate from the massive stellar companion of Cyg X-1 seem most reasonable (cf. Oda 1977).

A further crucial result for the dynamics of accretion disks is our finding that reconnection is ineffective at limiting magnetic field amplification within the inner disk. The ineffectiveness of field reconnection modifies the process responsible for "magnetic viscosity" as an angular momentum transfer mechanism (Eardley and Lightman 1975; Ichimaru 1977), a process which in standard models relies upon field dissipation within the disk; thus, we show that the ultimate dissipation of inner disk magnetic fields may well occur outside the disk, after the amplified, strong fields emerge via the action of buoyancy forces.

An important aspect of the present model is that the seed magnetic field acted upon by differential rotation and convection be internal to the accretion disk; this will occur if the seed field is either carried into the disk by the accreting plasma or generated locally via a homopolar (Biermann) dynamo mechanism (Bisnovatyi-Kogan and Blinnikov 1977). If the field threading the disk is externally connected—as would

be the case if the accretion disk penetrates the magnetosphere of the accreting central object—the topology of the atmosphere above the disk is forced to be open, and confinement of hot coronal plasma is unlikely; this case would correspond to regions in the solar corona ("coronal holes") which are topologically connected to the interplanetary medium (open) and are observed to be substantially less dense (and somewhat cooler) than the ambient closed corona and associated with vigorous plasma outflows (e.g., high-speed solar wind streams). This effect is independent of the possible suppression of convection within the disk by the external fields (to which the discussion of § II and the estimated limiting field strength given by eq. [11] apply). The connectivity of the seed magnetic field or, alternatively, its source, therefore may be a key element in determining the emission processes of accretion disks.

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