PLASMA EJECTION FROM MAGNETIC FLARES AND THE X-RAY SPECTRUM OF CYGNUS X-1

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ABSTRACT

The hard X-rays in Cyg X-1 and similar black hole sources are possibly produced in an active corona atop an accretion disk. We suggest that the observed weakness of X-ray reflection from the disk is due to bulk motion of the emitting hot plasma away from the reflector. A mildly relativistic motion causes aberration, reducing X-ray emission toward the disk. This in turn reduces the reprocessed radiation from the disk and leads to a hard spectrum of the X-ray source. The resulting spectral index is $\Gamma \approx 1.9/\sqrt{B}$, where $B = \gamma (1 + \beta)$ is the aberration factor for a bulk velocity $\beta = v/c$. The observed $\Gamma \approx 1.6$ and the amount of reflection ($R \approx 0.3$) in Cyg X-1 in the hard state can both be explained assuming a bulk velocity $\beta \sim 0.3$. We discuss one possible scenario: the compact magnetic flares are dominated by $e^\pm$ pairs, which are ejected away from the reflector by the pressure of the reflected radiation. We also discuss physical constraints on the disk-corona model and argue that the magnetic flares are related to magnetorotational instabilities in the accretion disk.

Subject headings: accretion, accretion disks — black hole physics — gamma rays: theory — radiation mechanisms: thermal — stars: individual (Cyg X-1) — X-rays: general

1. INTRODUCTION

Galactic black holes (GBHs) in their hard state and radio-quiet active galactic nuclei (AGNs) have similar X-ray spectra which are well explained by Comptonization of seed soft photons in a hot cloud with Thompson optical depth $\tau_T \sim 1$ and temperature $kT \sim 100$ keV (see reviews by Zdziarski et al. 1997; Poutanen 1998). The Compton reflection feature in observed spectra and a fluorescent iron line indicate the presence of relatively cold gas in the vicinity of the X-ray source. The most likely reflector is an accretion disk, and the hard X-rays are possibly produced in a hot corona of the disk (e.g., Bisnovatyi-Kogan & Blinnikov 1977; Liang 1979). The dissipation of magnetic energy in the corona may feed the X-ray luminosity (Galeev, Rosner, & Vaiana 1979, hereafter GRV). According to GRV, the reemitted spectra and a fluorescent iron line indicate the presence of a cold reemitting horizon in the disk due to a combination of the differential Keplerian rotation and the turbulent convective motions. The amplification timescale at a radius $r$ is given by $t_a \sim r/3v$, where $v$ is the convective velocity. GRV showed that inside luminous disks the field is not able to dissipate at the rate of amplification, and buoyant magnetic loops must emerge above the disk surface where the magnetic field may annihilate quickly. The rate of magnetic energy production per unit area of the disk equals $F_m = 2hw_B/n_B$, where $w_B = B^2/8\pi$ is the magnetic energy density and $h$ is the half-thickness of the disk. Assuming that the magnetic stress $T_m = B^2/4\pi$ is responsible for the transfer of...
angular momentum in the disk, one can compare $F_\mu$ to the total surface dissipation rate $F_c = 3\pi c \sigma T$, where $c$ is the sound speed in the disk (Shakura & Sunyaev 1973). One then finds $F_\mu/F_c = h/r$ (taking into account that $2\pi v_b^a/m_e = B_\mu/B_c = c/h$ in the GRV model). Hence, the GRV mechanism is able to dissipate only a small fraction, $\sim h/r \ll 1$, of the total energy. This is in conflict with the spectra of GBHs in the hard state, where $h/r \sim 1$, i.e., fast enough to explain the bulk of energy release as magnetic dissipation. Combined with the GRV argument for magnetic buoyancy, it follows that a large fraction of $F_c$ may dissipate in the corona.

The corona thus may be the place to which magnetic stress driving the accretion is transported and released. For a standard radiation pressure-dominated disk, $t_T \approx m_e c^2/\sigma T$, where $\sigma T$ is the Keplerian angular velocity. The magnetic field in the corona is likely to be strongly inhomogeneous, and it is plausible that the main dissipation occurs in localized blobs where the magnetic energy density $w_b$ is much larger than $t_T$ given by the standard model. The accumulated magnetic stress may suddenly be released on a timescale $t_T \sim 10 r_c/c$ (the “discharge” timescale, see Haardt et al. 1994), where $r_c$ is the blob size. This produces a flare of luminosity $L \sim r_c^2 w_b t_T$, which may have a compactness parameter $l \equiv L r_c^2/m_e c^2 \sim 10^{-5}$.

The compact flares can get dominated by $e^\pm$ pairs created in $\gamma-\gamma$ reaction (e.g., Svensson 1986). Pairs would keep an optical depth $\tau_T \sim 1$ during the flare and upscatter soft radiation entering the $e^\pm$ blob. The blob temperature is typically in the range $50$–$200$ keV depending on $l$ and on whether there is a nonthermal $e^\pm$ tail. A strongly localized $e^\pm$ flare with $l \sim 10^{-2}$–$10^{-1}$ may have a low temperature, $kT < 100$ keV, observed in some GBHs and AGNs. In the presence of a nonthermal $e^\pm$ tail, the flare may have even $kT \ll 50$ keV, which is observed in the case of GX 339–4 (Zdziarski et al. 1998).

Consider now one possible mechanism of bulk acceleration that should operate in an $e^\pm$ blob located above a reflector. The blob luminosity is partly reflected, and hence the pair plasma is immersed in an anisotropic radiation field. The net radiation flux $F \sim L/\sigma_T^2$ must efficiently accelerate the light pairs. The timescale for acceleration to relativistic velocities is $t_a \sim m_e c f/L_{\nu}^2$, where $f \sim F_{\nu} \sqrt{m_e c}$ is the accelerating radiative force. The shortness of $t_a \ll r_c/c$ implies that the bulk velocity saturates at some equilibrium value limited by the radiation drag (e.g., Gurvich & Rumyantsev 1965; Sikora & Wilson 1981). We discuss the pair acceleration by reflected radiation in more detail in §4.

The bulk motion in the hot blob does not necessarily mean that the blob itself moves. Instead, the dissipation region can be static. Then the escaping pairs are replaced by newly created $e^\pm$. The escaping pairs are immediately cooled down to the Compton temperature $kT_C \sim 10$ keV. Note that in a bright $\gamma$-ray source, $\gamma-\gamma$ interactions would produce a lot of pairs above/below the active regions. An outflow of cold pairs covering the whole inner region of the disk may then be created (Beloborodov 1998a).

3. BULK MOTION AND REFLECTION

It was noted previously that bulk motion of the flaring plasma may reduce the reflected radiation from an accretion disk (Woźniak et al. 1998). Reynolds & Fabian (1997) calculated the impact of motion on the Fe K\alpha line. Here, we evaluate the amount of reflection and study how fast motion affects the coupling between the X-ray source and the reflector. We assume that the active blob luminosity is due to Compton amplification of the reflected (reprocessed) radiation, and the multiply upscattered photons get isotropized in the plasma comoving frame. The angular distribution of the blob luminosity in the lab frame is then given by

$$L(\mu) = \frac{L}{2\gamma^2(1-\beta\mu)^2},$$

where $\mu = \cos \theta$, $\theta$ is the angle between the ray and the plasma velocity, $\gamma = (1-\beta^2)^{-1/2}$, and $L = \frac{L_\nu}{L}$, $d\mu$ is the total luminosity of the blob. For a typical spectral index, $\Gamma \sim 2$, the relativistic transformation of the specific luminosity $L(\mu)$ is the same as that of the bolometric luminosity (Rybicki & Lightman 1979). In equation (1), we assume that the dissipation region is static and there is no retardation effect (see Rybicki & Lightman 1979). For a static blob with a lifetime $t_T \gg r_c/c$, one may consider a time-independent picture in a fixed geometry like the usual disk-corona model. The only difference is that now the hot plasma in the blob has a bulk velocity.

We will assume that the velocity is perpendicular to the disk. The X-ray luminosity striking the disk is then given by

$$L_\nu = \int_{-1}^{0} L(\mu) d\mu = \frac{L(1+\beta/2)}{2\gamma^2(1+\beta)^2}.$$ (2)

For a system inclination $\mu = \cos \theta$, the apparent strength of reflection $R = \Omega/2\pi$ is

$$R = \frac{L_\nu}{L(\mu)} = \frac{(1+\beta/2)(1-\beta\mu)}{(1+\beta)^2}.$$ (3)

The reprocessed radiation from the disk is partly intercepted by the active blob. The intercepted soft luminosity may be expressed in terms of an effective $\mu_c$:

$$L_\nu = \chi \int_{-1}^{\mu_c} L(\mu) d\mu.$$ (4)

The case $\mu_c = 0$ describes a slab geometry of the active region, while $\mu_c \sim 1/2$ roughly corresponds to a blob with radius of order its height. The factor $\chi = 1-a$ represents the efficiency of reprocessing, and $a \approx 0.1$–$0.2$ is the disk albedo (e.g., Magdziarz & Zdziarski 1995).

The fraction of $L$ that returns as soft radiation is a function of $\beta$ and $\mu_c$:

$$L_\nu = \frac{L_{\nu\nu} - (1+\mu_c)\beta/2}{L_\nu \gamma^2(1+\beta)^2(1+\mu_c\beta)^2},$$ (5)

where $L_{\nu\nu} = \chi(1-\mu_c)L/2$ corresponds to the usual case of a static corona ($\beta = 0$). With increasing $\beta$, the intercepted soft luminosity is suppressed.

The spectral index of the Comptonized radiation emerging from the blob is determined by the Compton amplification factor taken in the plasma comoving frame. A luminosity $dL = L(\mu) d\mu$ transforms into the comoving frame as $dL' = \gamma^2(1-\beta\mu)^2 dL = (1+\beta\mu)^{-1}dL$, where the index $c$ stands for
the comoving frame (see Rybicki & Lightman 1979). This gives \( L' = L \) and \( L' = \gamma^2(L' - \gamma \dot{\Phi}) \), where \( \dot{\Phi} = \frac{1}{|L'_{(\mu)} \mu | \, d\mu} \) is the total flux of the soft radiation through the blob. We approximate \( L_{(\mu)} = L_{(1 - \mu)} \) at \( \mu > \mu_s \), and \( L_{(\mu)} = 0 \) at \( \mu < \mu_s \). Then

\[
\Phi_s = \frac{L_{(1 + \mu_s)}}{2},
\]

which yields \( L_s/L_{(1 + \mu_s)} = \gamma^2[1 - \beta(1 + \mu_s)/2] \). We then get the amplification factor

\[
A = \frac{L}{L_{s}} = \frac{2\gamma^2(1 + \beta)^2(1 + \mu_s \beta)^2}{\chi(1 - \mu_s)[1 - (1 + \mu_s \beta)^2]/4}.
\]

For \( \mu_s \approx 0.3 \), this expression can be approximated as \( A \approx A_0 B^\beta \) within 1% accuracy in the range \(-0.1 < \beta < 0.7\). Here, \( A_0 = 2/\chi(1 - \mu_s) \), \( B = \gamma(1 + \beta) \), and \( \psi = 2.7, 2.9, \) and 3.1 for \( \mu_s = 0.3, 0.4, \) and 0.5, respectively.

4. BULK VELOCITY IN A PAIR-DOMINATED FLARE

We now estimate the expected bulk velocity for a pair-dominated flare using a simple toy model. The power consumed by the \( e^- \) blob is the sum of two parts: the primary injected power due to dissipation of magnetic energy \( L_{(\text{diss})} \approx L - L_{(b)} \) and the reflected radiation intercepted by the blob \( = L_{(b)} \). \( L_{(\text{diss})} \) is probably injected in the form of relativistic particles which are supposed to share their energy and momentum with the thermal \( e^- \) plasma, most likely due to collective effects or due to synchrotron self-absorption (Ghisellini, Guilbert, & Svensson 1988). Let \( \Phi_{(b)} \) be the total vertical flux of the injected energy. The injected particles may accelerate or decelerate the thermal plasma depending on the angular distribution of injection. The reflected radiation always tends to accelerate the coronal plasma away from the disk.

The energy streaming into the blob per unit time equals \( L_{(\text{diss})} + L_{(b)} = L \), and the corresponding net energy flux equals \( \Phi_{(\text{diss})} + \Phi_{(b)} = \Phi \). The energy consumed during time \( dt \) increases the blob momentum by \( dp = (\Phi_{(b)} c) \, dt \) and the inertial mass by \( dm = \left( \frac{L_{(b)}}{c^2} \right) \, dt \). The net acceleration is \( d\beta = d(p/mc) = (\Phi - BL) \, dt/mc^2 \). The \( e^- \) plasma accelerates on a short time-scale (see § 2), and the bulk velocity should relax to an equilibrium value for which the net acceleration vanishes. The equilibrium velocity is determined by the equation

\[
\Phi - BL = 0.
\]

This equation also expresses the condition that the total energy flux vanishes in the comoving frame, \( \Phi^c = \gamma^2(\Phi - BL) = 0 \).

We do not possess a detailed model of injection in the blob and do not know the feedback of the plasma bulk velocity on the angular distribution of the injection. Therefore, we consider two particular cases:

1. \( L_{(\text{diss})} \) is isotropic in the lab frame.—Then \( \Phi_{(\text{diss})} = 0 \), and the equilibrium equation (8) combined with equation (6) yields the relation

\[
\frac{2\beta}{1 + \mu_s} = \frac{L_{(1 + \mu_s)}}{L}.
\]

On the other hand, \( L_{(1 + \mu_s)} \) is determined by equation (5). Equating equations (5) and (9), we get an equation for the self-consistent bulk velocity and find \( \beta \sim 0.1 \) for \( 0 \leq \mu_s \leq 0.5 \).

2. \( L_{(\text{diss})} \) is isotropic in the comoving frame.—In this case, the energy flux associated with \( L_{(\text{diss})} \), vanishes in the comoving frame and the equilibrium velocity is determined by the soft radiation only, \( \beta = \Phi/L_{(b)} \). This yields \( \beta = 1/2 \) for \( \mu_s = 0 \), and \( \beta = 3/4 \) for \( \mu_s = 1/2 \).

A proton fraction exceeding \( m_p/m_e \) would increase the plasma inertia, and then the bulk velocity may fall below the equilibrium value. Pairs will tend to stream through the heavy proton component, and a plasma instability may be initiated. In the stationary case, the stream velocity establishes itself just at the instability threshold, which may be expected to be comparable with the thermal electron velocity, \( \sim c/2 \).

5. CONCLUSIONS AND DISCUSSION

The hard state of accreting black holes may be explained as a state in which a large fraction of the luminosity is released in a magnetic corona atop a cold accretion disk. The magnetic energy is probably generated by the magnetorotational instability in the disk and then transported to the corona due to buoyancy. The soft state may be observed when the coronal activity is suppressed and most of the energy is dissipated inside the disk.

The energy release in the corona is likely to proceed in compact bright flares, where the local radiation flux strongly exceeds the average surface flux from the disk. If the flaring plasma is comprised of \( e^- \) pairs, it should be accelerated away from the disk by the pressure of the reflected radiation. The bulk velocity is then expected to be in the range \( 0.1 \sim \beta \sim 0.7 \).

The radiative acceleration of a pair plasma is one possible
reason for bulk motion in the flares. Even if the flare is dominated by a normal proton plasma, the heating may be accompanied by pumping a net momentum into the hot plasma at a rate \( \sim L/c \). The transferred momentum per particle per light-crossing time \( r/c \) is \( \sim l m c \). Since the flare duration \( t_0 \gg r/c \), the protons may acquire a momentum comparable to \( m c \). Note that magnetic flares in which hot plasma is ejected toward the disk are possible. This case is formally described by equations (1)–(7) with \( \beta \leq 0 \).

The impact of the source velocity on the observed reflection \( R \) (eq. [3]) and the Compton amplification factor \( A \) (eq. [7]) is summarized in Figure 1. Comptonization in the source produces a power-law X-ray spectrum. The photon spectral index \( \Gamma \) is related to \( A \) by

\[
\Gamma \approx 2.33(A - 1)^{-1}
\]

(10)

where \( \delta \approx 1/6 \) for GBHs and \( \delta \approx 1/10 \) for AGNs. This simple formula approximates within a few percent the results of calculations we performed by using the code of Coppi (1992) (see Beloborodov 1998c for details). Combined with equation (7), this yields the dependence \( \Gamma(\beta) \). Assuming a typical \( \mu_1 \approx 0.5 \), one can approximate \( \Gamma(\beta) \) as \( \Gamma \approx 1.9B^{-0.5} \) for GBHs and \( \Gamma \approx 2B^{-0.3} \) for AGNs, where \( B = \gamma(1 + \beta) \).

From the amount of reflection in Cyg X-1 (\( R \sim 0.3 \)), we infer a bulk velocity of the emitting plasma \( \beta \approx 0.3 \) for a system inclination \( \sim 30^\circ \) (see Fig. 1, bottom). On the other hand, from the reported intrinsic slope \( \Gamma \approx 1.6 \) and the corresponding \( \lambda \sim 10 \), we again infer the same \( \beta \approx 0.3 \). Thus, \( \beta \sim 0.3 \) resolves the two problems of the disk-corona model mentioned in § 1. The beaming of the corona luminosity \( L_c \) also leads to the weakness of observed reprocessed blackbody component in the spectrum, \( L_{bb} \sim \chi RL_{c} \sim L_c/4 \). Note that the inferred bulk velocity is comparable with the thermal electron velocity, \( \sim c/2 \) for \( kT \sim 100 \) keV.

The ejection model may apply to both GBHs and AGNs. The velocity of ejected hot plasma \( \beta \) may vary (and in some cases it might be \( \beta < 0 \)). An increase in \( \beta \) leads to decreasing \( R \) and \( \Gamma \). A correlation between \( R \) and \( \Gamma \) is observed in the spectra of GBHs and AGNs (Zdziarski, Lubinski, & Smith 1998a), and it is well reproduced by our model (Beloborodov 1998c). Note also that fast outflows may manifest themselves in optical polarimetric observations of AGNs (Beloborodov 1998b).

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REFERENCES

Balbus, S. A., & Hawley, J. F. 1998, Rev. Mod. Phys., 70, 1
Beloborodov, A. M. 1998a, MNRAS, submitted
Beloborodov, A. M. 1998b, ApJ, 496, L105
Beloborodov, A. M. 1998c, in ASP Conf. Ser., High Energy Processes in Accreting Black Holes, ed. J. Poutanen & R. Svensson (San Francisco: ASP), in press
Bisnovatyi-Kogan, G. S., & Blinnikov, S. I. 1977, A&A, 59, 111
Coppi, P. S. 1992, MNRAS, 258, 657
Esin, A. A., Narayan, R., Cui, W., Grove, J. E., & Zhang, S.-N. 1998, ApJ, 505, 854
Galeev, A. A., Rosner, R., & Vaiana, G. S. 1979, ApJ, 229, 318
Ghisellini, G., Guilbert, P. W., & Svensson, R. 1988, ApJ, 334, L5
Gierliński, M., Zdziarski, A. A., Done, C., Johnson, W. N., Ebisawa, K., Ueda, Y., Haardt, F., & Pihlps, B. F. 1997, MNRAS, 288, 958
Gurevich, L. E., & Rumiantsev, A. A. 1965, Sov. Physics—JETP, 20, 1233
Haardt, F., & Maraschi, L. 1993, ApJ, 413, 507
Haardt, F., Maraschi, L., & Ghisellini, G. 1994, ApJ, 432, L95
Krolik, J. H. 1998, ApJ, 498, L13
Liang, E. P. 1979, ApJ, 231, L111
Magdziarz, P., & Zdziarski, A. A. 1995, MNRAS, 273, 837
Poutanen, J. 1998, in Theory of Black Hole Accretion Disks, ed. M. Abramowicz, G. Björnsson, & J. Pringle (Cambridge: Cambridge Univ. Press), 100
Poutanen, J., Krolik, J. H., & Ryde, F. 1997, MNRAS, 292, L21
Poutanen, J., & Svensson, R. 1996, ApJ, 470, 249
Reynolds, C. S., & Fabian, A. C. 1997, MNRAS, 290, L1
Ross, R. R., Fabian, A. C., & Young, A. J. 1998, MNRAS, submitted
Rybicki, G. B., & Lightman, A. P. 1979, Radiative Processes in Astrophysics (New York: Wiley)
Shakura, N. I., & Sunyaev, R. A. 1973, A&A, 24, 337
Shapiro, S. L., Lightman, A. P., & Eardley, D. M. 1976, ApJ, 204, 187
Sikora, M., & Wilson, D. B. 1981, MNRAS, 197, 529
Stern, B. E., Poutanen J., Svensson, R., Sikora, M., & Begelman, M. C. 1995, ApJ, 449, L13
Svensson, R. 1986, in IAU Colloq. 89, Radiation Hydrodynamics in Stars and Compact Objects, ed. D. Mihalas & K.-H. Winkler (New York: Springer), 325
———. 1996, A&AS, 120, 475
Woźniak, P. R., Zdziarski, A. A., Smith, D., Madejski, G. M., & Johnson, W. N. 1998, MNRAS, 299, 449
Zdziarski, A. A. 1998, MNRAS, 296, L51
Zdziarski, A. A., Coppi, P. S., & Lamb, D. Q. 1990, ApJ, 357, 149
Zdziarski, A. A., Johnson, W. N., Poutanen, J., Magdziarz, P., & Gierliński, M. 1997, in The Transparent Universe, ed. C. Winkler, T. J.-L. Courvoisier, & Ph. Durouchoux (ESA SP-382; Noordwijk: ESA), 373
Zdziarski, A. A., Lubinski, P., & Smith, D. A. 1998a, MNRAS, submitted
Zdziarski, A. A., Poutanen, J., Mikołajewska, J., Gierliński, M., Ebisawa, K., & Johnson, W. N. 1998b, MNRAS, 301, 435