

On the Anti-Correlation Between High Accretion Luminosity and Radio Jet Ejection In GRO J1 655-40 and Other Objects

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ABSTRACT

A model is described in which radio jet formation in accreting objects is suppressed by processes which occur when the accretion rate approaches the Eddington limit. This was motivated by GRO J1655-40 and other objects which show an anticorrelation between high luminosity and the onset of a radio jet.

The jet-production mechanism employed is the Blandford-Payne MHD acceleration process, seeded by an e^+e^- pair wind. Observations indicate that all key ingredients of this mechanism are, or should be, present in these sources. Observed jet velocities and total powers are consistent with theoretical and numerical predictions of this model.

The primary jet-suppression mechanism proposed is the Papaloizou-Pringle compressible shear instability, which should disrupt the jet-producing region of the disk when the accretion rate approaches about one-third Eddington. The turn-on of the jet in GRO J1655-40 is consistent with this estimate. When *super-Eddington*, the disk should also drive an optically-thick, sub-relativistic wind, which may be a secondary jet-suppression mechanism. The presence of such a wind possibly is seen in the early spectral evolution of GRO J1655-40 and in the broad absorption lines of certain QSOs. Important tests of the model would be independent measurements of the compact object masses in these sources, a comparison of normal and BAL QSO X-ray spectra to see if the latter objects are significantly cooler, and a low frequency search around BAL QSOs to see if at least some have fossil radio sources.

Subject headings: accretion, accretion disks X-rays: stars quasars

1. Introduction

On 28 July, 1994 GRO J1655-40 underwent a strong hard X-ray outburst, reaching a luminosity of $6-12 \times 10^{37}$ ergs $(d/4\text{kpc})^2$ in 0.1-430 keV X-rays and remained at roughly that luminosity for about two weeks (Greiner 1994; Inoue et al. 1994; Sazonov & Sunyaev 1994; Harmon et al. 1995a). During this time the object was radio quiet. The subsequent drop in hard X-ray intensity beginning on ~ 8 August was accompanied on 13 August by a dramatic rise in radio flux density in the entire centimeter wavelength spectrum (Campbell-Wilson & Hunstead 1994) and the development of an expanding radio source with superluminal motion at an apparent β of 1.5 ± 0.4 and minimum actual β_{\min} of $0.84-0.1$ (Tingay et al. 1995), and a $\beta \cos \theta$ value of ~ 0.33 (Hjellming & Rupen 1995a), where β is the jet, component velocity in units of the speed of light and θ is the line-of-sight angle.

GRO J1655-40 outburst again in hard X-rays on 6 September 1994, but at only half the intensity of the first outburst. Again it was accompanied by a brightening of the central radio component as the X-rays decreased in intensity with a shorter delay between the beginning of the X-ray outburst and the radio outburst (Hjellming & Rupen 1994; Campbell-Wilson, McKay, & Lovell 1994).

The apparent anti-correlation between the periods of high X-ray luminosity and the radio jet ejection events, coupled with the general belief that X-rays are produced by accretion, indicates that the production (and suppression) of the radio jet is probably accretion related. This prompted Tingay et al. 1995 to suggest that accretion at greater than the Eddington limit,

$$\dot{M}_{Edd} \equiv \frac{L_{Edd}}{c^2} = 1.9 \times 10^{18} \text{ g s}^{-1} (\epsilon/0.1)^{-1} m \quad (1)$$

destroyed the conditions necessary for jet production by the accretion disk. Here $m \equiv (M/M_{\odot})$ is the dimensionless compact object mass, $L_{Edd} = 1.2 \times 10^{38} \text{ erg s}^{-1} m$ is the Eddington luminosity limit (above which radiation pressure due to electron scattering exceeds the force of gravity), and ϵ is the accretion efficiency. This paper examines the processes at work in high-luminosity accretion disks which might enable, as well as suppress, MHD jet mechanisms, and compares predictions with recent observations of GRO J165540 and other objects.

2. Jet Production in Accretion Disks

2.1. MHD Jet Production

One of the more successful classes of model for producing jets from accretion disks is that of magneto-hydrodynamic (MHD) acceleration and collimation, first outlined for central objects without magnetospheres (i. e., uncharged black holes and bare neutron stars) by Blandford & Payne 1982. If poloidal magnetic field lines, threading vertically through a geometrically thin disk and its corona, make an angle greater than 30 degrees with the rotation axis, then material trapped in the field will be accelerated centrifugally outward along those lines. Eventually, when the velocity of the plasma relative to the differentially-rotating field lines exceeds the Alfvén velocity, the field is wound helically around the axis, collimating the plasma. If the material in the field lines is continuously replenished from below, a continuous jet results.

Recently Meier, Payne, & Lind 1995 (see also Ustyugova et al. 1995) have performed a parameter study of the Blandford-Payne mechanism using detailed two dimensional MHD simulations. They find that well-collimated outflows occur over a large region of parameter space, and that at least mildly-relativistic flow (bulk $V_j \sim 3$) is possible (see Figure 1). Unlike the semi-analytic model of Blandford and Payne, however, the simulations show that jet ejection is far from self-similar. Rather, the jet emanates from a small region $R_{in} < R < R_j$, where R_{in} is the disk inner radius and $R_j \sim \text{few} \times R_{in}$.

The means by which the disk achieves the proper magnetic field configuration is still unknown (see, e.g., Begelman 1993). We will not address that issue here. Instead, we point out that necessary conditions for the MHD disk mechanism to work are

1. a stable inner disk in which to anchor the field lines
2. a poloidal magnetic field extending into the disk corona
3. a supply of material from below to replace the jet ejecta i.e., a "seed" wind

2.2. M1 II) Jet Production in High Luminosity Accretion Disks

2.2.1. *Why Jet Power*

To study the behavior of the Blandford-Payne mechanism in more detail, we treat it in the context of the “ α ” accretion disk model, where the viscous stresses, presumed to be magnetic and/or turbulent processes, scale with the disk pressure. A general expression for the magnetic field is obtained from conservation of angular momentum (Shakura & Sunyaev 1973)

$$\frac{B^2}{8\pi} \approx \alpha p = \frac{M\Omega}{4\pi H} \mathcal{F}(R) \quad (2)$$

where α is the viscous stress scaling parameter, Ω is the angular velocity (assumed to be Keplerian), and $\mathcal{F}(R)$ is of order unity except very near R_{in} . The disk half-thickness H contains all the physics of heating and cooling balance and varies greatly from model to model. We restrict our models to high luminosity disks only

$$\dot{m} > \dot{m}_I \equiv 1.6 \times 10^2 (\alpha/0.1)^{-\frac{1}{2}} \dot{m}^{-\frac{1}{2}} \quad (3)$$

where $\dot{m} \equiv M/\dot{M}_{Edd}$ is the dimensionless accretion rate and \dot{m}_I is the rate above which the most models predict radiation pressure effects to be important in the inner disk region. (This is essentially equation 2.18 of Shakura & Sunyaev 1973 after correcting for a factor of 2 error.) Recent models show that at least a portion of this region should be optically thin, with electrons at relativistic temperatures ($> 10^9$) and pressure dominated by the even hotter ions (Shapiro, Lightman, & Eardley 1976; Liang 1979; Misra & Melia 1995).

In spite of more than 20 years of effort, $H(\dot{M}, R)$ is still a poorly-known function in this regime. We therefore have parameterize our model in terms of the disk thickness at the radius $R_e \equiv R_g c$ where most of the luminosity is produced

$$h \equiv H_e/R_e \equiv H_e c/R_g \quad (4)$$

which is generally a function of \dot{m} only, and

$$H(\dot{M}, R) = H_e (R/R_e)^s \quad (5)$$

where $R_g \equiv GM/c^2$ is the gravitational radius. For most disk models which do not include the effects of advective cooling, $h \sim \dot{m}$, with $s = 0$ for the original (Shakura & Sunyaev 1973) optically thick disk models

and $s \approx 1$ for the optically thin models (Misra & Melia 1995). For transonic accretion models which do not include other forms of advective energy transport (convexion, turbulence, wind) $h \rightarrow 1$ for $\dot{m} \rightarrow \infty$ and $s \approx 1$ (Abramowicz et al. 1988).

Using equation (2) and the expression for MHD jet power from Blandford & Payne 1982, we derive an *instantaneous* jet power and a *non-luminous* accretion rate induced by the jet acceleration process of

$$\frac{P_{BP}^{jet}}{c M c^2} = \frac{\dot{m}_{BP}}{\dot{m}} = h^{-1} > 1 \quad (6)$$

That is, in the Blandford-Payne mechanism the jet-producing region of the disk generates more power than is available in the accretion itself and a higher accretion rate than the steady-state value. Therefore, in the absence of additional matter and energy sources, jet ejection must be very unsteady, depleting the central disk region on a time scale of

$$\tau_{BP, drift} = h \tau_{drift} = \alpha^{-1} h^2 r_e^2 \tau_{dyn} \quad (7)$$

where $r_e \equiv R_e/R_g = c^{-1}$ and $\tau_{dyn} = 5 \times 10^{-6} s r_e^{\frac{3}{2}} \dot{m}$ is the orbital timescale at R_e . Afterward, the jet shuts off until accretion from the outer disk replenishes the central region and the jet begins again. For compact objects of 1 - 10 M_\odot , $\alpha \sim \dot{m} \sim 0.1$, and $r_e \sim 10$, these time scales are quite short compared to the hours it takes for radio observations of the jet: $\tau_{dyn} \sim 10^{-4} - 10^{-3} s$; $\tau_{BP, drift} \sim 1 - 10 s$; and $\tau_{drift} \sim 10 - 10^2 s$. Therefore, VLBI observations measure the time-averaged jet power over the duty cycle h , or

$$\langle P_{BP}^{jet} \rangle \approx h P_{BP}^{jet} = c M c^2 = L \quad (8)$$

which is of order the accretion luminosity. This is a straightforward result: if the jet is to be powered by the disk, then its time-averaged power (over a few hours for stellar systems) is limited to the time average of the accretion power itself.

2.2.2. *The Seed Wind*

A variety of winds have been studied in the context of accretion disks, but most are not good candidates for the seed wind. One which may be relevant for $\dot{m} > \dot{m}_I$ is the hot pair wind. Misra & Melia 1995 show that for moderately high accretion rates,

$$\dot{m} > \dot{m}_{pw} \approx 4 \times 10^2 (c/0.1) (\alpha/0.1)^{0.55} \quad (9)$$

¹This fit to their numerical results is accurate in the range $\alpha \in 0.04 - 0.5$.

the disk becomes so hot that e^+e^- pair-production and escape becomes an important cooling mechanism, with up to half the energy emitted by the disk in the form of pairs. The pair wind therefore is present for accretion rates slightly above \dot{m}_J . The wind should be driven by radiation pressure on the electrons and positrons since the disk is already greatly super-Eddington for e^+e^- pairs by a factor of $\dot{m}(m_p + m_e)/2m_e \gtrsim 40$.

Even for a relativistic pair wind ($\Gamma_{pw} \sim 2$), the ram pressure of this wind of light particles is usually much less than $2m_e c^2 \sigma_T^{-1} R_g^{-1}$ (Misra & Melia 1995), which is far smaller than the expected magnetic field pressure just above the disk (equation 2)

$$\frac{\rho_{pw} v_{pw}^2}{P_{mag}} \ll 0.02 \dot{m}^{-1} h \text{ (c/O.J.) } (r/r_e)^{s-\frac{1}{2}} \quad (10)$$

Because the pair wind is such a copious source of escaping particles, but not disruptive to the magnetic field, it is an ideal candidate for seeding the Blandford-Pringle jet mechanism.

2.3. Suppression of MHD Jets for Accretion Rates Approaching the Eddington Limit

2.3.1. By Papaloizou-Pringle Instabilities

Accretion disks with $\dot{m} > \dot{m}_J$ are subject to a number of known instabilities. In particular, the compressible Papaloizou-Pringle instability (Papaloizou & Pringle 1984; Papaloizou & Pringle 1985) is expected to induce strong turbulent motions on a dynamical time scale in differentially-rotating ('accretion tori' when the radial shear in a region $\Delta R \sim 2H$ exceeds the sound speed. The length scale of motions generated are expected to be $\gtrsim c_s/\Omega \sim H$. If $H \ll R$, an overall disk structure should nevertheless obtain. Also the existence of a coronal magnetic field should not be affected by this process, although the field will be perturbed by Alfvén waves generated by eddies occurring where the field lines are anchored.

However, as $\dot{m} \rightarrow 1$ ($L \rightarrow L_{Edd}$), $H \ll R$ is no longer valid. Compressible Papaloizou-Pringle instabilities, then, have the potential for completely disrupting the central region between R_{in} and R_j when the induced motions are of order $R_{je} \equiv (R_{in} R_j)^{\frac{1}{2}}$. To investigate this we replace c_s in the compressible Papaloizou-Pringle instability criterion with the orbital velocity of the jet-producing region $v_{\phi j} = R_{je} \Omega(R_{je})$, which is greater than the sound speed by a factor of order h^{-1} . Disruption then occurs when

$(\Omega(R_{in}) - \Omega(R_j))^2 > \Omega^2(R_{je})$, which, for angular velocity laws with $\Omega \propto R^{-q}$, gives a condition on h

$$h \sim \frac{H(R_{je})}{R_{je}} \approx \frac{(R_j - R_{in})}{(R_j + R_{in})} > 0.2 - 0.3 \quad (11)$$

for values of q ranging from 2 (uniform angular momentum) to $\frac{3}{2}$ (Keplerian rotation), respectively. For accretion disks where $h \sim \dot{m}$, this corresponds to a critical accretion rate and luminosity of $\dot{m}_{PP} \approx 0.2 - 0.3$. Above this accretion rate the central portion of the disk cannot produce a jet because Papaloizou-Pringle instabilities destroy the disk interior to $R_{PP} \equiv \dot{m} R_g / c \dot{m}_{PP}$. Furthermore, if $H \propto R$ ($s \approx 1$), the instability could engulf much of the hot pair-producing region outside R_{PP} as well.

2.3.2. By a Strong Super-Eddington Wind

Many models of accretion disks with super-Eddington accretion ($\dot{m} > 1$) have been constructed in which no radiation-pressure-driven outflow of electrons and protons occurs. To do so, these models maintain a delicate balance between the local surface radiation pressure and effective gravity and/or they maintain a transonic radiation-trapped advective flow toward the black hole in the disk interior (Maraschi, Reina, & Treves 1976; Paczynski & Wiita 1980; Begelman & Meier 1982; Abramowicz et al. 1988). In the advective case photons generated above L_{Edd} are carried into the black hole before they have a chance to radiate.

In light of the disruptive compressible Papaloizou-Pringle instability, it is unclear if such delicate radiation pressure balance or laminar trapped flow can be maintained. The region interior to R_{PP} is likely to be characterized by violent sonic or supersonic turbulent motions, shocks (pushing α to unity or above), and gross inhomogeneities. Under such circumstances photons generated in the torus interior are more likely to be exposed to low optical depth regions, either in between the inhomogeneities themselves or by being carried to the disk surface by the violent turbulent convection. The release of a fraction of these photons in excess of the Eddington limit will drastically change the character of the inflow, driving a significant portion of it out in a strong, optically thick wind (Shakura & Sunyaev 1973; Meier 1982a,c). Indeed, in the case of an accreting neutron star, outflow is the *only* possibility, as there exists no horizon into which photons can be lost.

Models of this wind with $\dot{m} \sim 1$ and $\alpha \sim 1$ (more appropriate here than $\alpha \sim 0.1$) are as hot as the thin accretion disk models (6×10^9 I) with $\dot{m} \sim 1$, but for super-Eddington accretion rates, e.g., $\dot{m} \sim 10$, the disk is enshrouded in a cooler optically thick wind of only 10^8 K (Meier 1982b,c). Its ram pressure is far higher than that of the pair wind:

$$\frac{\rho_{scw} v_{scw}^2}{p_{mag}} = 0.7 f(\dot{m}) \alpha^{1/2} (r/r_i)^{s-1/2} \quad (12)$$

where $f(\dot{m})$ is of order unity for $\dot{m} > 2$ and $r_i \approx \dot{m} c^{-1}$ is the radius at which the wind is ejected. This expression is of order unity and possibly larger if the field drops off as $B \propto r^{-1}$ or faster in the corona ($s \geq \frac{1}{2}$). A super-Eddington wind, therefore, has the potential for disrupting the field in the disk corona.

3. Application to GRO J 1655-40

Our model for GRO J 1655-40 characterizes the source as going through all the previously-discussed stages in reverse order. The first outburst undergoes accretion from a finite reservoir at a value above the Eddington rate. MHD jet-production mechanisms remain suppressed until the accretion rate declines to the point where Papaloizou-Pringle instabilities no longer disrupt the central disk region. MHD mechanisms then "kick in", considerably after the initial rise in hard X-ray intensity, and eject an unsteady jet until the reservoir is completely depleted.

3.1. The First Outburst

3.1.1. The Super-Eddington Phase

The clearest indication that during the first outburst (28 July - 16 August 1994) GRO J1655-40 may have undergone a super-Eddington event with a strong wind is the evolution of its spectrum (Harmon 1995b), which is quite different from that expected of normal sub-Eddington accretion disks. Shortly after the initial (< 1 day) rise to its peak 40-430 keV luminosity, the source's X-ray spectrum softened sharply over a ~ 1 -day period from a photon spectral index of -3.0 to -3.5. During that time the X-ray luminosity remained steady. The spectrum then began to slowly harden to -2.5 over a ~ 9 -day period, again while the hard X-ray luminosity remained within $\sim 15\%$ of the peak.

This behavior can be explained by a super-Eddington event with a steadily decreasing accretion rate.

The initial softening of the spectrum during a period of constant luminosity signalled the initial increase of the accretion rate to super-Eddington values and the formation of a large, cooler, optically thick super-Eddington wind. Once \dot{m} exceeded unity the steady hard X-ray luminosity (dominated by the photons at the lower energies of the band) remained roughly fixed at a fraction

$$f_{HX} \equiv L_{HX}^{max}/L_{Edd} \quad (13)$$

of the Eddington limit with the remainder coming out in lower energy X-rays. The subsequent hardening of the spectrum was due to the shrinkage of the wind photosphere as the accretion rate slowly dropped. However, because \dot{m} was still above unity, the bolometric luminosity remained at L_{Edd} , and L_{HX} remained at $f_{HX} L_{Edd}$.

With detailed radiative transfer models of the wind, one should be able to deduce the run of temperature and density with spherical radius and, therefore, the variation of the mass loss and accretion rates with time. In lieu of those models, as an *ad hoc* example, we shall assume a linearly-falling accretion rate of the form

$$\dot{m} = \dot{m}_0 - \ddot{m}_0(t - t_0) \quad (14)$$

where the 0 subscript refers to the beginning of the outburst. Taking 8 August as the time when $\dot{m} = 1$ and 16 August when $\dot{m} = 0$, we fit the data to an initial accretion rate of $\dot{m}_0 \approx 2.9$. For $\alpha \sim 1$, simple thermal models of super-Eddington winds (Meier 1982c) predict temperatures of 1.7×10^9 K (150 keV) during the first few days of the outburst (after the initial softening) and about 5.9×10^9 K (500 keV) on 8 August, when $\dot{m} \approx 1$. While GRO J1655-40 has an excess above the power law spectrum near ~ 500 keV, the spectral resolution of the observations is too low to distinguish any possible annihilation line (Harmon et al. 1995a). The model wind velocity of $\dots 2 - 3 \times 10^4$ km s $^{-1}$ applies mainly to the hard X-ray emitting gas. (This velocity is much greater than the optical line widths seen during the outbursts [Della Valle 1994; Wagner & Bertram 1994] which we attribute to material rotating in the disk at $\sim 10^4$ Schwarzschild radii, where the effective temperature is $\sim \text{few} \times 10^4$ K.)

While the 40-430 keV luminosity is about 2.2×10^{37} ergs $^{-1}$ ($d/4$ kpc) 2 during outburst (Harmon et al. 1995a), there is a substantial amount of flux at lower energies as well. In between the outbursts, there is

a similar amount of luminosity in the ROSAT 0.1 – 2.4 keV band (Greiner 1994) and in the ASCA 2–10 keV band (Inoue et al. 1994), the latter with an $E^{-2.1}$ spectrum, but there is not yet information on these values *during* the first outburst. Depending on whether or not the 2 – 20 keV X-rays during outburst had their inter-outburst value or increased along with the 40–430 keV band (and had an $E^{-2.1}$ spectrum), we estimate a bolometric luminosity of $L_{bol} \sim 6 - 12 \times 10^{37} \text{ erg s}^{-1}$. This implies a compact object mass of $m = L_{bol}/1.2 \times 10^{38} \text{ erg s}^{-1}$ or

$$M_{1655-40} \sim 0.5 - 1 M_{\odot} \quad (15)$$

3.1.2. Sub-Eddington Accretion Phase

After 8 August L_{HX} began to decrease steadily. This was accompanied by another softening for the first 4 days, but on 13 August the *radio flux* increased dramatically (Campbell-Wilson & Inunstead 1994), a jet was ejected (Tingay et al. 1995), and the 40 – 430 keV spectrum once again hardened (to -2.2 ; Harmon 1995b).

This behavior can be explained by a continuation of the decreasing accretion rate postulated above, but now to values below $\dot{m} = 1$, where the spectral hardness is *directly* related to accretion rate. The final hardening, however, is not explainable by standard accretion theory. It is probably related to the jet itself, e.g., Comptonization of disk and jet photons by the relativistic radio plasma.

The fact that the jet waits *another four days* to form), when the accretion rate and hard X-ray luminosity are of order

$$\dot{m} = \frac{L_{bol}}{L_{Edd}} \approx \frac{L_{HX}}{L_{HX}^{max}} = 0.35 \pm 0.15 \quad (16)$$

indicates that the mechanism for suppressing jet production operates even when the accretion is sub-Eddington. If this is due to Papaloizou-Pringle instabilities disrupting the central disk region, then our rough estimate for when this occurs (\dot{m}_{JP}) is consistent with the numbers above. Better theoretical estimates must await very detailed near-1-Eddington disk models, but the observations indicate that it should be well below the Eddington value (unity).

3.1.3. The Jet Ejection Phase

We estimate GRO J 1655-40's total jet power in two different ways. The first is similar to that used by

Mirabel & Rodriguez 1994 for GRS 1915+105. However, here we assume that the ejected radio plasma is composed of e^+e^- pairs rather than ions and electrons, use minimum radio source energy arguments (Pacholczyk 1970) to determine the lobe internal energy and mass, and assume that both electrons and positrons radiate and that the low frequency cutoff for this compact source is 1 GHz. For the 23 August lobe flux values of $S_{2.3 \text{ GHz}} = 0.42 \text{ Jy}$, source dimensions of 130 mas x 320 mas, a distance of 4 kpc (Tingay et al. 1995), and a maximum Doppler factor of 0.89, we find that in the rest frame

$$E_{lobe} \gtrsim 5.0 \times 10^{42} \text{ erg} \quad (17)$$

If the pair plasma is relativistic ($\gamma_{e^+} \sim \gamma_{e^-} \gg 1$), the lobe mass is just E_{lobe}/c^2 , and for a bulk $\Gamma_{j,min} \sim 1.7$, the minimum total lobe energy (including kinetic and internal) is $\Gamma_{j,min} M c^2 \gtrsim 8.3 \times 10^{42} \text{ erg}$. If formed during the initial 2-day rise in radio flux, the jet power must be at least

$$\langle P_{1655-40}^{jet} \rangle \gtrsim 4.8 \times 10^{37} \text{ erg s}^{-1} \quad (18)$$

Note that this value is three orders of magnitude smaller than that found by Mirabel & Rodriguez 1994, chiefly because their assumption of a mono-energetic electron distribution implicitly assumes a particular ratio of particle to field energy in the plasma which is several orders of magnitude out of equipartition.

The second method assumes that the observed lobes are not moving blobs but rather radiating cocoons surrounding the working surface of a much narrower jet as it drills through the interstellar medium (Norman et al. 1982; Lind, J'sync, Meier, & Blandford 1989). This assumption is more accurate if the lobe is low density compared to the interstellar medium, since moving low density blobs will be quickly halted by a higher density medium (Meier, Sadun, & Lind 1991). Indeed, in the limit of only relativistic particles in the lobe, its mass density should be just the relativistic inertia of the energy density, or $6 \times 10^{-27} \text{ g cm}^{-3}$, which is well below even intercloud densities ($3 \times 10^{-25} \text{ g cm}^{-3}$; Allen 1973). We therefore estimate the instantaneous jet power by equating it to the energy loss rate in the lobe. These losses are kinetic (essentially that calculated in equation 18), radiative (only $\sim 1032 \text{ erg s}^{-1}$), and adiabatic. Note that equation 17 yields a lobe pressure of $\sim 2 \times 10^{-6} \text{ dynes cm}^{-2}$, orders of magnitude above the typical interstellar pressure of $10^{-13} - 10^{-12} \text{ dynes cm}^{-2}$,

so adiabatic expansion should occur in a sound travel time or so, by which time the jet working surface has travelled further away from the source. Since the lobe length does not shrink or lengthen appreciably with time, we assume that all lobe internal energy is lost adiabatically within the time it takes the jet to expand one lobe length (~ 4.9 days), with an estimated instantaneous internal energy loss rate of $\sim 1.2 \times 10^{37} \text{ erg s}^{-1}$. Presumably this energy loss occurs during much of the jet *propagation* phase, which we take to be 2-3 c-folding decay times for the radio flux (Harmon et al. 1995a) or ~ 13 -19 days. Integrating over this time period, dividing as we did above by the estimated jet ejection time of ~ 2 days, and adding the kinetic power derived above, we get an initial power in the jet *ejection phase* of

$$\langle P_{1655-40}^{jet} \rangle \gtrsim (12-18) \times 10^{37} \text{ erg s}^{-1} \quad (19)$$

We conclude that either estimate of the jet power is of the same order as the bolometric X-ray luminosity when the jet was ejected ($4 - 6 \times 10^{37} \text{ erg s}^{-1}$), in agreement with the average power predicted by the Blandford-Pringle process (equation 8).

3.2. Subsequent Outbursts

Since the first hard X-ray outburst of GRO J1655-40, there have been several others occurring at intervals of a few months. Near the end of each the radio luminosity increases briefly, in qualitative agreement with the nature of the anticorrelation between jet ejection and high X-ray luminosity discussed here (Harmon et al. 1995a).

However, a surprising and unexpected behavior also has occurred during this time: the strength of each radio outburst has decreased approximately exponentially with time (Harmon et al. 1995a) until it became undetectable, while the strength of the X-ray outbursts remained approximately the same. GRO J1655-40 has evolved from a radio loud to a radio quiet "mini-quasar"!

In light of this behavior we note the following points. As discussed herein, the suppression of jet formation by Papaloizou-Pringle instabilities and/or super-Eddington winds may be responsible for short-term delays and anticorrelations between high accretion luminosity and radio emission. However, this mechanism cannot be the long-sought-after explanation for the difference between radio loud and radio quiet objects. If it were, each subsequent radio out-

burst would have had about the same radio luminosity and GRO J1655-40 always would have become a radio loud object after the accretion becomes sub-Eddington. The exponential decrease in radio outburst strength indicates that the overall ability of the object to produce any jet, even when accretion is sub-Eddington, has been decaying slowly with time.

Candidate mechanisms which could cause GRO J1655-40 to evolve into a radio quiet object could be intrinsic (e.g., evolution in a black hole property such as angular momentum) or extrinsic (e.g., evolution in properties of the accreting matter, such as its magnetic properties) without affecting our conclusions about short-term anticorrelations. A full understanding of GRO J1655-40's long-term evolution is beyond the scope of this paper, but may be very important for unifying radio loud and radio quiet objects, particularly quasars.

4. Application to Other Objects

4.1. GRS 19154105

GRS 1915-105 is a hard X-ray transient similar to GRO J1655-40 but lying at a somewhat greater distance of 12.54 ± 1.5 kpc (Mirabel & Rodriguez 1994) with a somewhat higher outburst X-ray luminosity of $3 \times 10^{38} \text{ erg s}^{-1}$. It also ejected (on 19 March, 1994) relativistically-moving radio components with apparent β_s up to 1.25 ± 0.15 . Recently it was reported that, like GRO J1655-40, there was a delay between the beginning of the X-ray and radio outbursts, with the radio occurring near or during the X-ray decline (Deal et al. 1995).

We therefore propose a model for this source which is nearly identical to that for GRO J1655-40, except that the central object mass ($m = L_{bol}/1.2 \times 10^{38} \text{ erg s}^{-1}$) must be larger

$$M_{19154105} \gtrsim 2.4 M_{\odot} \quad (20)$$

Again we identify the jet ejecta as a pair plasma, use minimum energy arguments to estimate the internal lobe energy, and compute jet power values both with and without taking cocoon adiabatic losses into account. With the Mirabel & Rodriguez 1994 lobe parameters 15 days after the initial radio outburst (similar to the 10 day age used for GRO J1655-40 above) and the $\lesssim 3$ day acceleration time used by these authors, we derive jet powers of

$$\langle P_{19154105}^{jet} \rangle \gtrsim (1-5) \times 10^{38} \text{ erg s}^{-1} \quad (21)$$

which again is comparable to the photon luminosity at the time. These estimates are three orders of magnitude *smaller* than that obtained by Mirabel & Rodriguez 1994, chiefly because of the assumption of minimum energy rather than an arbitrary monoenergetic beam energy.

4.2. 1E 1740.7-2942 and Cygnus X-1

We only briefly mention 1E 1740.7-2942, because it also is a hard X-ray source with a radio jet, and Cyg X-1 because of its similar X-ray continuum. Since independent information indicates a mass for Cyg X-1 of $\sim 10M_{\odot}$, and therefore a similar mass for 1E 1740.7-2942 by analogy, the 5×10^{37} erg S^{-1} X-ray luminosities for these objects yields a dimensionless accretion rate of only $\dot{m} \sim 0.04$. This is far below the Eddington limit, well below the threshold $\dot{m}_{th} \sim \dot{m}_{PP}$ for jet, suppression try Papaloizou-Pringle instabilities, and only barely at the Misra & Melia 1995 threshold for pair-wind production] (equation 9). As suggested by these authors) the lack of a strong radio jet in Cyg X-1 undoubtedly has more to do with the lack of large numbers of pairs than with disruption of a thick central torus. 1E 1740.7-2942, however, seems to be in a state very similar to that in which GRO J1655-40 and GRS1915+105 form their jets, but on a more permanent basis. The accretion rate is in the range $0.04 \lesssim \dot{m} \lesssim 0.35$, where the disk is thin, relatively stable, and hot. The copious quantities of pairs available (Misra & Melia 1993) and the hospitable environment for MHD mechanisms make an e^+e^- jet very likely. We favor M111) jet acceleration over the "hotplate" mechanism (Misra & Melia 1993) because of the former's ability to accelerate jets to relativistic velocities near the black hole and its potential for a high degree of collimation.

4.3. Broad Absorption Line and other QSOs

Broad absorption line quasi-stellar objects (BAL QSOs) are quasars which show outflow with velocities up to 0.1c. Models of these objects (Weymann et al. 1991; Goodrich & Miller 1995) usually depict them as normal quasars but viewed approximately edge-on to a flared accretion disk surface (not in the equatorial plane). Material stripped from the disk surface by an ever-present quasar wind, and blown toward the observer, form the broad absorption lines.

However, these objects are rather peculiar in that *all* are radio quiet (Stoček et al. 1992), while 10 -

15% of normal quasars are radio loud (Kellerman et al. 1989; Miller, Peacock, & Mead 1990). We therefore suggest that these objects are in a high-luminosity state similar to that in GRO J1655-40 or GRS1915+105 where radio jet production is suppressed. In this scenario BAL QSOs are those few quasars which have been undergoing accretion near or above the Eddington limit ($\dot{m} > 0.35?$). The presence of P-Cygni profiles with $v \sim 0.1c$ may indicate a super-Eddington accretion event in progress, while detached broad lines may indicate a recent event in the past with the current rate still large but sub-Eddington. Although high polarization [I] AL QSOs may indeed be those viewed when the accretion disk is nearly edge-on (Goodrich & Miller 1995), low polarization ones could be those viewed when it is more face-on.

Such models for BAL QSOs with large covering factors have been discounted in the past, not because of the small numbers of QSOs compared with the total QSO population, but because of a discrepancy between absorbed and emitted resonance line photons (e.g., Lyman α ; Hamann, Korista, & Morris 1993). While the broad and deep absorption troughs imply that resonance photons have been absorbed, because resonance scattering preserves photon number, these should eventually be re-emitted from the BAL cloud surfaces in a bright emission line centered near zero rest velocity. If the BAL cloud covering factor is large, the number of photons seen in emission should be close to that seen absorbed. However, in BAL QSOs the number of resonance photons seen in emission is small compared to that absorbed, leading one to conclude that the covering factor is small, with only a portion of the [I] AL cloud surfaces facing the observer.

There is, however, a better explanation for the lack of resonance photon re-emission - dust absorption within the BAL clouds themselves (Voit, Weymann, & Korista 1993; Turnshek 1995). It has been shown (Meier & Terlevich 1981) that even a small amount of dust in an ionized region can reduce the number of emitted resonance line photons by factors of several, and normal amounts of dust can destroy re-emission entirely. BAL clouds tend to be *overabundant* in metals (Turnshek 1995) and, therefore, may be quite dusty. If so, it is possible that they have little or no resonance line re-emission and that the observed emission line comes only from the broad line region (BELR) interior to the BAL region. Since in this case the covering factor cannot be determined by

counting photons, its value is unconstrained by the observations.

In the super-Eddington model for 11A], QSOs, the small number of these objects compared to normal QSOs is interpreted as being due to the scarcity of very high accretion rates rather than to viewing angle. As the wind photosphere only covers the X-ray emitting region of the disk, one would not expect normal and BAL QSO optical properties to differ appreciably. However, these objects may differ dramatically in their X-ray properties, with the BAL QSOs being significantly cooler. While the super-Eddington model seems somewhat unsatisfactory in that it requires two separate explanations for the resonance line emission/absorption discrepancy and the small numbers of 11A], QSOs, it offers the distinct advantage of explaining the lack of radio loud BAL objects as a GRO J1655-40-type anticorrelation between high luminosity, windy objects and strong radio emission.

For massive compact objects ($m \gtrsim 10^6$) the onset of a hot inner region should occur at a much lower accretion rate than for stellar objects ($\dot{m}_I \sim 1.6 \times 10^{-3}$). If jets are suppressed only near $\dot{m}_{th} \sim 0.35$ in these objects also, then there is a broad range of accretion rate in which e^+e^- jets can form by the pair-fed Blandford-Payne mechanism. The fact that only a small percentage of QSOs are radio sources indicates that either all objects identified as quasars are accreting near \dot{M}_{Edd} (which we consider unlikely) or there is an additional process at work in determining whether or not a QSO is a radio source.

We therefore conclude, as we did with GRO J1655-40, that while super-Eddington accretion can explain the anticorrelation between high luminosity, windy objects and strong radio source emission, it is not a good candidate for explaining why most (*non*-BAL) QSOs are radio quiet. In our super-Eddington model BAL QSOs should be considered as being drawn from *both* the radio quiet and radio loud populations, but in the latter case jet ejection is temporarily arrested by super- or near-super-Eddington accretion. When the rapid accretion subsides, some BAL objects should then return to being radio loud. High resolution, low frequency radio searches around 11A], QSOs should reveal at least some of them (up to 10-15%) to have fossil radio sources.

5. Conclusions

We have outlined a model in which production of relativistic radio jets by accreting objects is suppressed by processes which occur when the accretion rate approaches the Eddington limit. This was motivated by the observation that GRO J 165540 and other objects show an anticorrelation between high luminosity and onset of a radio jet. From the observations our best estimate for the jet-suppression accretion rate threshold is about one-third of the Eddington value.

Our jet-production model invokes the Blandford-Payne (M II) acceleration process, seeded by an e^+e^- pair wind. All the key ingredients of this mechanism should be present in accretion disks with m within 2-3 orders of magnitude of the Eddington limit (thin, relatively stable disk; magnetic field; hot, 6×10^6 K corona). Numerical simulations of this mechanism indicate that for a wide range of the input parameters, at least mildly relativistic jets with a high degree of collimation are ejected from a small region of the disk near the compact object. Total jet powers derived from the observations are consistent with those predicted theoretically.

We suggest the Papaloizou-Pringle instability as the primary jet-suppression mechanism. This is expected to be especially destructive to the jet-producing region when the accretion rate is near the Eddington limit and the inner disk is thick. Because of the instabilities, the onset of an optically thick, sub-relativistic ($v \lesssim 0.1c$), super-Eddington wind, which carries off power in excess of L_{Edd} , seems inevitable, even when the compact object is a black hole with the potential for swallowing a portion of the power through transonic advective accretion. The presence of such a wind can be inferred indirectly from the spectral evolution of GRO J1655-40 at high luminosity and directly from the absorption lines in BAL QSOs—another class of objects known to be radio quiet. The wind also may be a secondary jet-suppressing mechanism.

The main weakness of the model lies in the still very poor theoretical understanding of the details of the accretion processes when $\dot{m} > \dot{m}_I$. The variation of coronal magnetic field strength with disk radius is still poorly known (let alone the details of exactly how the field configuration required for MHD jet acceleration is maintained). We also have only poor estimates of the accretion rate at which the Papaloizou-

Pringle instabilities in *thick* disks have a destructive effect on the conditions for M1 II) jet acceleration. Least known are the structure of the inner regions of super-Eddington disks, the dominant process which carries off most of the excess power (wind or advective accretion) in the black hole case, and the effect the disk instabilities have on this situation.

Independent measurements of compact object masses, and therefore L/Edd , in these objects would be a first step toward testing this model. Very high spatial resolution Space VLBI observations during jet production may constrain some collimation mechanisms. High spectral resolution observations of X-ray lines during the high luminosity phase would help identify a super-Eddington wind. X-ray continuum observations of normal and BAL QSOs could test for a cooler optically thick wind in the latter class of objects. High resolution, low frequency radio observations of BAL objects should reveal at least some (up to 10-15%) to have fossil radio sources and therefore have been radio loud in the past.

While we believe the mechanism proposed herein to be responsible for the short-term anticorrelations between high luminosity/windy phenomena and radio jet ejection seen in some objects, we do not believe it to be the explanation for radio quietness in all objects. Some more fundamental, long-term mechanism must be at work which would, for example, allow GRO J 1655-40 to evolve over several months from a radio loud to a radio quiet "mini-quasar". Understanding this latter mechanism may help unify virtually all QSOs into a single model.

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6. FIGURE CAPTIONS

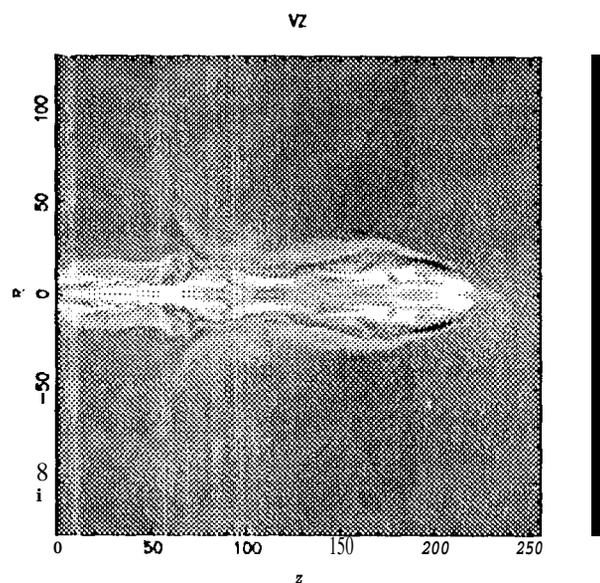


Fig. 1. Typical simulation of the Illanctford-1'sync jet-procluctio~r mechanism evolved to 250 inner disk dynamical times (from Meier, J'sync, and Lind 1995), showing a grey-scale plot of the axial velocity. The accretion disk is applied as a boundary condition at left in the equatorial plane, with a low-velocity wind blowing along disk-corona magnetic field lines main tained at a polar angle of 54° . Length units are in grid points, with an inner disk radius of $R_{in} \sim 5$.