ADVECTION-DOMINATED FLOWS AROUND BLACK HOLES AND THE X-RAY DELAY IN THE OUTBURST OF GRO J1655—40

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ABSTRACT

We show that the time delay between the optical and X-ray outbursts of the black hole soft X-ray transient source GRO J1655—40, observed in 1996 April, suggests that the accretion flow in this object must consist of two components: a cold outer accretion disk and an extremely hot inner advection-dominated accretion flow (ADAF). In quiescence, the model predicts a spectrum that is in good agreement with observations, with most of the observed radiation coming from the ADAF. By fitting the observed spectrum, we estimate the mass accretion rate of the quiescent system and the transition radius between the disk and the ADAF. We present a detailed numerical simulation of a dwarf nova-type instability in the outer disk. The resulting heat front reaches the ADAF cavity promptly; however, it must then propagate inward slowly on a viscous timescale, thereby delaying the onset of the X-ray flux. The model reproduces the observed optical and X-ray light curves of the 1996 April outburst, as well as the 6 day X-ray delay. Further, the model gives an independent estimate of the quiescent mass accretion rate that is in very good agreement with the rate estimated from fitting the quiescent spectrum. We show that a pure thin disk model without an ADAF zone requires significant tuning to explain the X-ray delay; moreover, such a model does not explain the quiescent X-ray emission of GRO J1655—40.

Subject headings: accretion, accretion disks — black hole physics — stars: individual (GRO J1655—40) — X-rays: stars

1. INTRODUCTION

The binary X-ray source GRO J1655—40 (also called X-ray Nova Scorpii 1994) is a member of the class of so-called soft X-ray transients (SXTs) or X-ray novae. In these systems, a low-mass Roche lobe-filling secondary star transfers mass through an accretion disk onto an compact object, a neutron star or a black hole. Compared to a neutron star transient, a black hole transient (BHT) generally has a larger X-ray outburst amplitude and a lower quiescent luminosity, which is a signature of the black hole's event horizon (Narayan, Garcia, & McClintock 1997b). The mass of the black hole primary in GRO J1655—40 is \( \approx 7.0 \ M_\odot \) (Orosz & Bailyn 1997, hereafter OB).

GRO J1655—40 is an exceptional BHT in view of its frequent outbursts in recent years. Most BHTs have recurrence times of decades or longer, whereas GRO J1655—40 has gone into outburst several times since its discovery on 1994 July 27 by BATSE on the Compton Gamma Ray Observatory (CGRO) (Zhang et al. 1994). Two subsequent outbursts occurred in 1995 late March (Wilson et al. 1995) and in 1995 July (Harmon et al. 1995). Following an extended period of X-ray quiescence, the source again went into outburst in 1996 April, as discovered by the All-Sky Monitor (ASM) on the Rossi X-Ray Timing Explorer (RXTE) (Remillard et al. 1996; Levine et al. 1996).

Thus, GRO J1655—40 has remained active off and on for nearly three years. Other BHTs have shown X-ray and optical activity several months after an outburst; however, none of them have sustained their activity for more than about a year (e.g., Tanaka & Shibazaki 1996). The frequent outbursts of GRO J1655—40 in recent years may be due to an enhancement of the mass transfer rate, estimated by OB to be relatively high—\( 2.2 \times 10^{17} \) g s\(^{-1}\). One should note, however, (H. Ritter 1997, private communication) that OB assume that the secondary is a giant, whereas GRO J1655—40 had not yet reached the giant branch. Systems with companions in the Hertzsprung gap should transfer mass at a very high rate, \( \geq 10^{18} \) g s\(^{-1}\) (Kolb et al. 1997), since the secondary expands on a thermal timescale; this is most probably not the case now in GRO J1655—40, and no observational determination of the mass transfer rate is available at present. As it happens, the value given by OB is plausible, and we will use it in the following discussion. In any case, on longer timescales GRO J1655—40 behaves more like other BHTs, since no outburst of the source has been reported in the previous 25 years. GRO J1655—40 is also distinguished by its radio outbursts, which are associated with superluminal expansion events (Tingay et al. 1995; Hjellming & Rupen 1995).

About six days prior to the most recent X-ray outburst of GRO J1655—40 (1996 April), a remarkable optical precursor was observed (Orosz et al. 1997, hereafter ORBM). As shown in Figure 1, starting from an initially quiescent state, the optical intensities (\( BVRI \)) were observed to rise grad-
mechanism successfully explains dwarf nova outbursts in the framework of the disk instability model (DIM) (see Cannizzo 1993a and references therein). The dwarf nova DIM has been extended to SXTs by Mineshige & Wheel er (1989; see also Cannizzo, Chen, & Livio 1995).

The DIM requires the quiescent accretion disk to be in a cold state; the accretion rate therefore must be everywhere lower than the critical accretion rate \( M_{\text{crit}}(R) \propto R_a \), where \( a \approx 2.6 \) (see, e.g., Ludwig, Meyer-Hofmeister, & Ritter 1994). This in turn implies that the accretion rate onto the compact object required by the DIM is extremely low, \( \sim 10^{-6} \text{ g s}^{-1} \) (Mineshige & Wheel er 1989; Lasota 1996a). However, observations of X-ray emission from quiescent BHTs imply accretion rates that are several orders of magnitude higher (McClintock, Horne, & Remillard 1995; Verbunt 1996; Narayan, McClintock, & Yi 1996 [NMY]; Narayan, Barret, & McClintock 1997 [NBM]; Robinson et al. 1997). A similar problem is encountered in some quiescent dwarf novae (see, e.g., Lasota 1997). It is clear that the standard DIM cannot apply to BHTs (Lasota 1996a, 1996b); in some cases it must be modified even to describe dwarf nova outbursts (Meyer & Meyer-Hofmeister 1994; Livio & Pringle 1992).

A model for SXTs in quiescence was proposed by NMY in which the accretion flow occurs as a thin disk only outside a transition radius \( R_{tr} \sim 10^3 \) Schwarzschild radii, while for \( R < R_{tr} \) the flow forms an advection-dominated accretion flow (ADAF) (Abramowicz et al. 1995; Narayan & Yi 1994, 1995; for a recent review see Narayan 1997). In the NMY model, observed X-rays are emitted with a very low efficiency by the ADAF, while UV and optical luminosity is produced by the outer disk. However, Lasota, Narayan, & Yi (1996) have pointed out that the NMY model is not self-consistent, as the outer disk is relatively hot and therefore subject to a thermal instability, contrary to the assumed stationarity (see also Wheel er 1997). Independently, Mineshige (1996) proposed that the SXT outbursts are due to an instability in the outer disk. More recently, NBM have shown that self-consistent models for the spectra of V404 Cyg and A0620—00 can be obtained by an ADAF that extends outward to \( R_{tr} \sim 10^5 \) Schwarzschild radii. In these models, the accretion disk is cool and the optical-UV flux is mostly supplied by synchrotron emission from the ADAF. In such a model, the transient outburst originates in the outer cold disk and is due to a dwarf nova-type instability.

In § 2, we discuss the consequences of the delay between optical and X-ray; this is intended to give, via an order of magnitude estimate, physical insight into the numerical results described in § 4. In § 3 we develop an ADAF model for GRO J1655—40 in quiescence. In § 4 we demonstrate that a time-dependent model of the outburst of GRO J1655—40 implies the existence of a two-component accretion flow. The parameters determined for this flow agree with the parameters independently determined in § 3.

2. INTERPRETATION OF THE X-RAY DELAY: EVIDENCE FOR A TWO-COMPONENT DISK

2.1. The UV Delay in Dwarf Nova Outbursts

The X-ray delay observed in the outburst of GRO J1655—40 is analogous to the well-known UV delay observed for dwarf novae (e.g., Warner 1995 and references therein). For dwarf novae, the rise in the UV flux starts

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Fig. 1.—Observed optical and 2–12 keV X-ray light curves during the initial phase of the 1996 April outburst of GRO J1655—40 (from ORBM). For the sake of clarity, only one average data point per night is plotted. \( B, V, \) and \( I \) magnitudes are represented by triangles, squares, and stars respectively. Time has been set to zero at an arbitrary point close to the onset of the outburst.

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usually for several days and brighten by about 30% before the onset of the X-ray outburst. In this article, we examine only one aspect of the complex behavior of GRO J1655—40 in outburst, namely, the properties of the optical precursor and the X-ray outburst and what they imply for models of quiescent BHTs and the outburst mechanism. As argued by ORBM, the substantial delay between the optical eruption and the X-ray outburst, the “X-ray delay,” may provide support for the advection-dominated accretion flow (ADAF) model of the inner regions of the quiescent accretion disk.

One outburst mechanism that has been developed for SXTs is mass transfer instability (Hameury, King, & Lasota 1986); however, in general it cannot reproduce the characteristic timescales of SXTs, and it has therefore been rejected (Gontikakis & Hameury 1993). It is now clear that the outburst mechanism must result from a disk instability. A natural candidate for such an instability is the thermal (and viscous) instability that arises from abrupt changes in opacity when hydrogen becomes partially ionized. Such a
about 5–15 hr after the beginning of the optical outburst. In the framework of the standard disk instability model (DIM), one can interpret the UV delay as due to an “outside-in” (or type A) outburst. According to the DIM (e.g., Cannizzo 1993a and references therein), a thermal instability in the outer disk creates an inward-propagating heat front. This front transforms the disk from a cold (quiescent) state to a hot state. Because the UV flux is mainly emitted close to the white dwarf, one expects a delay in its rise equal to the time it takes the front to travel from the outer disk to the white dwarf. In the DIM, however, the calculated travel time of the front is much shorter than the observed UV delay time (Pringle, Verbunt, & Wade 1986; Cannizzo & Kenyon 1987). Thus in its standard form the model fails to explain the UV delay.

Two solutions have been proposed in order to rescue the DIM; both of them invoke a central “hole” in the accretion disk. At the edge of such a hole, an inward moving heat front would have to stop. The hole would then fill up on a viscous timescale, much longer than the heat-front propagation time. Livio & Pringle (1992) suggested a mechanism for creating such a hole: they argued that at quiescent mass accretion rates, the magnetic field of a weakly magnetized white dwarf could disrupt the inner accretion disk. They showed that such a model can reproduce the UV delay observed in dwarf nova outbursts. This model cannot apply to systems in which the accreting object is a black hole, since a black hole cannot support a permanent magnetic field.

Meyer & Meyer-Hofmeister (1994) proposed a different scheme for quiescent accretion onto a white dwarf that also results in a central hole. They invoke inefficient cooling in the disk’s upper layers, which leads to the formation of a hot corona and ultimately to the evaporation of the inner disk. As a result, the inner accretion flow consists of a pure coronal plasma. A similar solution for quiescent SXTs has been independently proposed by NMY. In both cases, the hot inner flow provides a natural explanation for the hard X-ray emission observed in quiescent dwarf novae and SXTs.

Whether a hole is created by magnetic fields or by evaporation, the effect on the outburst of a dwarf nova is similar: when the heat front arrives at the inner edge of the truncated disk it cannot propagate any farther, and the (surface) density contrast slowly fills up the hole on a viscous timescale, thereby producing the required delay of the UV outburst.

Below, we show that observations of the 1996 April outburst of GRO J1655 − 40 imply the presence of a two-component accretion flow in this system, and that the parameters deduced from observations agree very well with a model comprising an outer cold disk and an inner hot ADAF, as proposed by NBM.

2.2. The X-Ray Delay in the Outburst of GRO J1655 − 40

We assume that the 1996 April outburst of GRO J1655 − 40 was triggered by a dwarf nova–type instability in a standard cold disk, extending from the transition radius \( R_{tr} \) to some outer radius \( R_{out} \). This disk is most probably marginally stable with respect to this instability (e.g., Lasota et al. 1996), or it may be globally unstable (see below). Once the instability is triggered, the resulting heat front propagates with speed \( v_f \approx \alpha c_s \) (Meyer 1984), where \( c_s \) is the equatorial-plane sound speed in the hot phase. Thus, the time it takes for the front to travel a distance \( R_{out} \) is

\[
t_f \approx \frac{R_{out}}{\alpha c_s} \sim 2.8 \times 10^{-1} R_{10} T_4^{-1/2} \text{ hr},
\]

where \( T_4 = (T/10^4 \text{ K}) \) is the central disk temperature, and \( R_{10} \) the radius in units of \( 10^{10} \text{ cm} \). If the instability starts at a sufficiently large radius, the front may take up to a day to reach the inner edge of the disk; this has been proposed by ORBM as the origin of the observed delay between the \( I, \bar{R}, V, \) and \( B \) light curves. Although very tempting, this explanation suffers from the fact that for any reasonable values of the radius and the primary mass, the effective temperature jumps to \( 10,000–11,000 \text{ K} \) on a thermal timescale at the point where the instability sets in. At these temperatures, \( (B − V) \sim 0, (V − I) \sim 0 \), and the \( B, V, \) and \( I \) magnitudes should increase simultaneously. Dilution of the disk light by the ADAF and secondary light does not alter this conclusion. Thus, the disk instability model cannot easily account for a 1 day delay at optical/IR wavelengths.

It is interesting to note that the observations of ORBM do not quite cover the very initial rise of the outburst, since their first observation shows that the \( R, V, \) and \( I \) fluxes have already risen, whereas the \( B \) flux remains at its quiescent value. This indeed implies the existence of a delay; however, the value obtained by ORBM assumes a linear extrapolation, which may overestimate the delay for two reasons. The initial rise could be very nonlinear, with a sharp increase in the optical light from the disk, in which case the delay could be only a few hours. Another possibility is that the mass transfer rate from the secondary may have significantly increased since the previous observation of the source in quiescence, one month prior to the outburst. In fact, it is quite plausible that such an increase could have triggered the thermal instability that caused the 1996 April outburst. In any case, such an increase in the mass transfer rate would not show up in the \( B \) band, because the quiescent temperature of the system does not exceed 6500 K. We also note that variations of the mass transfer rate from the secondary might help to account for the observed irregular outburst patterns of GRO J1655 − 40.

It is finally worth noting that a slow rise in brightness just before an outburst has been observed in WX Hyi (Kuulkers et al. 1991). There, the rise in the \( V \) band starts first, followed by the rise at higher frequencies. The rise in \( B \) is faster than in \( V \). The delay times are rather short, from 0.01 to 0.09 days. Such an observation shows that the rise to outburst in dwarf nova–type of eruptions is more complex than predicted by the simple version of the DIM.

When the heat front arrives at the transition radius \( R_{tr} \) where the dense (cold) disk ends, it cannot propagate any farther toward the black hole; however, the resulting (surface) density contrast will propagate inward as a result of viscosity. The speed at which the density front propagates is

\[
v_{visc} = \frac{v}{\alpha},
\]

where \( v \) is the kinematical viscosity \( (v = \alpha c_s H) \), and \( w \) is the scale of the density gradient. The density contrast will travel a distance \( R_{tr} \) in a time

\[
t_{visc} = \frac{R_{tr}}{v_{visc}}.
\]
The width $w$ can be written as (see, e.g., Cannizzo 1996)

$$w = \beta (HR)^{1/2}.$$  

(4)

For the cooling front during the decay phase, the width corresponds to $\beta \sim 1$ in this formula, while for the heating front the width is smaller by a factor of a few at most (see, e.g., Ludwig et al. 1994). We therefore expect $\beta$ to be $\lesssim 1$ during the rise phase considered in this paper.

If we identify the observed X-ray delay with $t_{\text{vis}}$, we can estimate the transition radius as

$$r_{\text{tr}} \approx 3.6 \times 10^4 \frac{R_{\odot}}{R_{*}} \frac{R_S}{r_{\text{tr}}} \frac{m_1 M}{M_\odot} \frac{\beta^{-4/3}}{T_4},$$  

where $r_{\text{tr}} = R_{\odot}/R_*$, $R_S = 2GM/c^2$, $m_1 = M/M_\odot$ is the mass of the central black hole, and $t_{\text{vis}}$ is the X-ray delay time in days. (Here and elsewhere we use the symbols $R$ and $r$ to refer to the radius in physical units and Schwarzschild units, respectively.) For $m_1 = 7$, $t_{\text{vis}} \approx 5$, and $\beta = 0.3$, one obtains $r_{\text{tr}} \approx 4.6 \times 10^3 \frac{\beta^{-4/3}}{T_4}$, which shows that the transition between the hot ADAF and the cold outer disk occurs at $r_{\text{tr}} \sim 10^4$. Remarkably, this same value for the transition radius was deduced for models of two BHTs by NBM in a completely independent way. Furthermore, in the next section we show that $r_{\text{tr}} \sim 10^4$ is close to the radius at which the outer disk becomes unstable to the dwarf nova instability.

The outer disk radius in GRO J1655−40 is given by

$$r_{\text{out}} \approx 7.3 \times 10^5 m_1^{-2/3} P_{60}^{2/3} \approx 2 \times 10^5,$$  

where $P_{60} = P_{\text{orb}}/60$ hr is the orbital period.

3. ADAF PLUS COLD DISK MODEL FOR GRO J1655−40 IN QUIESCENCE

In the previous section we have seen, using order of magnitudes estimates, that the X-ray delay in the outburst of GRO J1655−40 indicates that the quiescent state of the system should consist of a two-zone flow. In this section and in §4 we show that this in indeed the case. The thin accretion disk can extend only down to a transition radius $r_{\text{tr}}$, which must be greater than a few thousand Schwarzschild radii. Inside this radius, the flow must either be absent or have a much lower density than in the thin disk. This picture is very similar to the two-zone model proposed by NMY and NBM for fitting the spectral data of V404 Cyg and A0620−00 in quiescence; in that model, the flow inside the transition radius consists of an extremely hot two-temperature ADAF. Here we use the ADAF plus thin disk model to fit the spectral data on GRO J1655−40 in quiescence and thereby estimate some key parameters of the quiescent accretion flow.

We first select a set of system parameters for GRO J1655−40 that we use as inputs for our models of the source. These include the black hole mass, the binary inclination, the distance, and the velocity at the inner edge of the outer thin accretion disk. Second, we summarize the multiwavelength data (X-ray, optical, and radio) used to constrain our models of the quiescent state. Finally, we use the spectral data to constrain the remaining parameters of the model.

The mass of the black hole in GRO J1655−40 and the inclination of the system are very well determined; we adopt $M_1 = 7 M_\odot$ and $i = 70^\circ$ (OB). In §2.2, we estimated the outer radius of the thin accretion disk to be $r_{\text{out}} \approx 2 \times 10^5$; in the spectral models we take $\log (r_{\text{out}}) = 5.0$. This parameter does not need to be determined very accurately, as it has very little effect on the results. Based on studies of the radio jets, we adopt a distance of $D = 3.2 \text{ kpc}$ (Hjellming & Rupen 1995).

Dynamical and geometrical information about the thin accretion disk can be obtained from studies of the broad, double-peaked Balmer lines (Smak 1981; Horne & Marsh 1986); from this the transition radius $r_{\text{tr}}$ can be constrained. Of interest here is $v_{\text{in}}$, the projected velocity at the inner edge of the thin accretion disk. Estimates of this velocity have been inferred for several SXTs from orbit-averaged profiles of the H\alpha emission line (see NMY and references therein).

GRO J1655−40 is a difficult case because of the strong H\alpha absorption line present in the spectrum of the F subgiant secondary, and because of the relative brightness of the secondary. In order to obtain a useful spectrum at H\alpha, we formed a sum of 73 spectra that had been collected over a wide range of orbital phases in 1995 April–May, when the system was near quiescence (Bailyn et al. 1995), and then subtracted the spectrum of the best-fitting F5 IV star (J. A. Orosz 1996, private communication). In this way we measured the width of the H\alpha emission line and estimated $v_{\text{in}} \geq 1045 \text{ km s}^{-1}$. The inner edge of the thin disk, or equivalently the transition radius $r_{\text{tr}}$ between the thin disk and the ADAF, is then estimated to be $r_{\text{tr}} = (c \sin i/v_{\text{in}})^2/2 \leq 3.6 \times 10^4$ (NBM). We present below models corresponding to a range of values of $\log (r_{\text{tr}})$ consistent with this constraint.

The multiwavelength data, $v F_r$ vs. $v$, are summarized in Table 1. To derive the optical flux (entry 5), it was necessary to subtract a large stellar component. We assumed that the residual (nonstellar) component contributed 5% ± 2% of the total light (OB); for the total optical flux from the system we used the apparent magnitude and reddening given in OB. The optical flux in the $V$ band and its error bar are shown in Figure 2.

Apart from the optical and BATSE data (entry 2), all the remaining data in Table 1 were obtained during an intensive campaign of observations of GRO J1655−40 in quiescence conducted in 1996 March by Robinson and his collaborators (Robinson et al. 1997). All of the upper limits in Table 1 and below are quoted at the 3 $\sigma$ level of confidence. The OSSE flux limit (entry 1) corresponds to an intensity upper limit of 40 mcrab (100–600 keV). This is off-scale in Figure 2 and is not plotted. GRO J1655−40 was detected by ASCA (entry 3) at an (unabsorbed) flux level of $F_X (2–10 \text{ keV}) = (1.6 \pm 0.7) \times 10^{-13} \text{ erg cm}^{-2} \text{ s}^{-1}$; the photon power-law index ($\alpha$) is estimated to be 1.5 ± 0.6 (Robinson et al. 1997). Both the flux (with its error bar) and the allowed range of spectral slope are indicated in Figure 2. A stringent flux limit was obtained using the ROSAT HRI detector (entry 4): $F_X (0.2–2 \text{ keV}) < 6.1 \times 10^{-13} \text{ cm}^{-2} \text{ s}^{-1}$ (3 $\sigma$), assuming $\alpha = 2.1$ and $N_H = 5 \times 10^{21} \text{ cm}^{-2}$ (Robinson et al. 1997). This upper limit is indicated by the arrow in Figure 2. Finally, the VLA limits on GRO J1655−40 (entries 6 and 7) correspond to a flux limit of 0.5 mJy at both 4.9 GHz and 8.4 GHz (not shown on Fig. 2).

The BATSE limit (entry 2) is an average over the period 1996 April 19–April 30, and corresponds to a photon flux (20–200 keV) of $-0.0001 \pm 0.0002$ photons cm$^{-2}$ s$^{-1}$ (assuming $\alpha = 2.8$), or an intensity limit of 26 mcrab. This is consistent with the OSSE limit. Recall, however, that the inferred start time of the X-ray rise at 2–12 keV is $25.4 \pm 0.8$, 1996 April (Orosz et al. 1997). So, the BATSE
limit includes an ~ 5 day period when the 2–12 keV X-ray outburst was underway.

We have attempted to fit the quiescent spectral data on GRO J1655 – 40 using an ADAF plus thin disk model, analogous to the models described in NBM. The solid lines in Figure 2 represent four models, each model consisting of a pure ADAF for radii $r > r_{tr}$, and a thin accretion disk in the radius range $r_{tr} \leq r \leq r_{min}$. The models correspond to $\log (r_{tr}) = 4.5, 4.0, 3.5, 3.0$, respectively. The transition from the thin disk to the ADAF occurs via evaporation into a corona, as described in NBM.

The models assume equipartition between gas and magnetic pressure ($\beta = 0.5$ in the notation of NMY and NBM), and the viscosity parameter is taken to be $\alpha = 0.3$ in the ADAF region. It is assumed that a fraction (0.001) of the viscous energy directly heats the electrons in the ADAF (and the corona), while the remaining 0.999 of the energy goes initially into the ions (i.e., $\delta = 0.001$; see NBM for details). The shape and normalization of the computed spectra are quite insensitive to values chosen for $r_{tr}$, $\beta$, $\alpha$, and $\delta$ (see Figs. 3–5 in NBM).

The models shown here assume that the outer disk is steady, with a mass transfer rate equal to that feeding the ADAF. However, in a transient system, the local mass transfer rate in the quiescent outer disk varies with radius and increases from the ADAF value at the transition radius to the rate at which mass is transferred from the secondary at the outer radius. Therefore, the models shown here give only a zeroth order approximation of the contribution of the outer disk; in fact, these models provide a lower limit to the thin disk luminosity. It turns out, however, that the ADAF luminosity always dominates the outer disk, so the approximation is justified. This is demonstrated by the dashed line in Figure 2, which shows the spectrum of the outer disk calculated using the exact results of the time-dependent calculations presented in § 4; it can be seen here that the outer disk contributes only a minute fraction of the total flux.

In each model, only one parameter has been adjusted, namely the mass accretion rate. This has been optimized so that the model flux in the ASCA band agrees with the observed flux. Despite the large range of $r_{tr}$ covered by the four models, the mass accretion rates vary little from one model to another; in Eddington units, the accretion rates range from 0.0034 to 0.0037, corresponding to physical accretion rates of $M = (3.4–3.7) \times 10^{16}$ g s$^{-1}$. Thus, the spectral models constrain the $M$ of GRO J1655 – 40 in quiescence quite well. Technically, the models determine only the parameter combination $M/\alpha$, so $M$ depends on a knowledge of $\alpha$. However, the value of $\alpha$ in ADAFs is fairly well constrained by the various studies done to date (Narayan 1997), and is unlikely to vary by more than a factor of ~ 3 either way from the value we have assumed, $\alpha = 0.3$. This suggests that the above estimate of $M$ in the ADAF is good to about a factor of 3.

The models shown in Figure 2 are consistent with all the measurements available at present, including the OSSE and VLA limits (not shown in Fig. 2). Note in particular that the

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**TABLE 1**

Quiescent Spectrum of the Nonstellar Component of GRO J1655 − 40

| Entry Number | Wavelength (Å) | $\log v$ (Hz) | $\log (\nu F_\nu)^a$ (ergs cm$^{-2}$ s$^{-1}$ Hz$^{-1}$) | Observatory | Reference |
|--------------|---------------|---------------|-------------------------------------------------|-------------|-----------|
| 1........... | 1.05          | 19.818b       | $< -9.49$                                       | CGRO/OSSE   | 1         |
| 2........... | 0.17          | 19.258b       | $< -9.44$                                       | CGRO/BATSE  | 1         |
| 3........... | 2.3           | 18.123        | $-12.96$                                        | ASCA        | 1         |
| 4........... | 16.6          | 17.258        | $< -12.59$                                      | ROSAT/HRI   | 1         |
| 5........... | 5000          | 14.736        | $-11.24$                                        | CTIO        | 2         |
| 6........... | $3.6 \times 10^8$ | 9.924     | $< -16.38$                                      | VLA         | 1         |
| 7........... | $6.1 \times 10^8$ | 9.690     | $< -16.61$                                      | VLA         | 1         |

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*a Flux limits are at the 3 σ level of confidence.

*b Central frequency computed assuming a Crab-like spectrum with energy index $a_\nu = 1.1$ (see NBM).

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* Central frequency computed assuming $a_\nu = 0.5$ (see § 3).

References—(1) Robinson et al. 1997; (2) Orosz & Bailyn 1997.
models fit the observed optical flux, predict the correct slope in the ASCA band, and lie below the ROSAT flux limit.

A rather obvious point is that the quiescent data are incompatible with any model based only on a thin accretion disk. The ROSAT and ASCA data clearly indicate that (1) the X-ray flux of GRO J1655–40 in quiescence lies below the optical flux, and (2) the X-ray spectrum is quite hard, with a photon index of less than 2.7 at 2 $\sigma$. A thin accretion disk model, with either a constant or variable $M$ as a function of radius, cannot possibly reproduce such a spectrum. Thus, GRO J1655–40 is similar to A0620–00 and V404 Cyg (see NMY and NBM), where again the quiescent spectra are found to be inconsistent with a pure thin disk model but are well fitted with an ADAF plus thin disk model.

Although the spectral fit does not help us to determine $r_{tr}$, it is possible to obtain a fairly strong constraint on $r_{tr}$ by considering the stability of the outer thin disk (Lasota et al. 1996; NBM). Specifically, the outer disk will be unstable to dwarf nova instability if it has an effective temperature greater than about 5000 K; therefore, the quiescent disk cannot exceed this temperature at any radius. The four models presented in Figure 2, with $\log (r_{tr}) = 4.5, 4.0, 3.5,$ and 3.0, have effective temperatures at the transition radius of $T_{max} = 1700, 3700, 8400,$ and 20,000 K, respectively. These correspond in almost all cases to the maximum temperature in the outer disk; in fact, they are lower limits to $T_{max}$ in each case, owing to our simplifying assumption of a constant $M$ in the outer disk. The requirement that $T_{max} < 5000$ K thus provides the constraint $\log (r_{tr}) > 3.7$, or $r_{tr} > 5000$. Just prior to outburst, we expect the thin disk to be very close to the limiting value of $T_{max}$. We therefore estimate that GRO J1655–40 had its transition radius at $r_{tr} \approx 5000$, or $R_{tr} \sim 10^{10}$ cm, at the time of the 1996 April outburst.

4. INSTABILITY OF THE OUTER DISK

In this section, we present numerical simulations of the dwarf nova instability in the outer disk of GRO J1655–40. We compare the results with observations of the early stages of the outburst, paying particular attention to the 6 day delay between the optical and X-ray outbursts. We also compare the mass accretion rate implied by the outburst calculations with the quiescent $M$ determined independently in the previous section. The calculations presented here have been done with the code described in the appendix.

In the following discussion, we assume that the mass transfer rate from the companion star, i.e., the accretion rate at the outer rim of the accretion disk, has the value given by OB, i.e., $\dot{M}_{transfer} = 2 \times 10^{-7}$ g s$^{-1}$. For this value of $\dot{M}_{transfer}$, the outer cold disk is unstable to dwarf nova instability (see, e.g., Ludwig et al. 1994), so we are guaranteed that the code will produce an outburst.

We assume that the transition between the outer thin accretion disk and the ADAF is due to evaporation into a corona that gradually erodes matter in the disk as the cold inflowing material approaches the transition radius; such a model has been proposed for dwarf novae by Meyer & Meyer-Hofmeister (1994; see also NMY; NBM). We have estimated the evaporation rate by the following approximate method. Narayan & Yi (1995) have derived the maximum allowable mass transfer rate in the ADAF at small radii to be $\dot{M}_{ADAF, max} = 0.3 \pi \alpha^2$ (in Eddington units), and have found that $\dot{M}_{ADAF, max}$ decreases at large radii ($r > 10^3$) (see also Abramowicz et al. 1995). Assuming that the mass transfer rate within the inner ADAF is equal to the maximum, and using $M_1 = 7 M_\odot$, we adopt the following approximate prescription for evaporation:

$$\dot{M}_{ev} = \frac{2.8 \times 10^3}{(1 + K R_{tr,10}^2)} g s^{-1}, \quad R_{tr} \geq R_{tr,10},$$

where $K$ is a constant adjusted to give the required value of the transition radius, and $R_{tr,10}$ is the transition radius in units of $10^{10}$ cm. The local surface density evaporation rate is then related simply to the derivative of $\dot{M}_{ev}$ with respect to $R$,

$$\Sigma_{ev} = \frac{1}{2\pi R} \frac{d\dot{M}_{ev}}{dR} = \frac{9 \times 10^{-4}K}{(1 + K R_{tr,10}^2)} g s^{-1} \text{ cm}^{-2}.$$  

This prescription for the evaporation is numerically close to the formula given by Meyer & Meyer-Hofmeister (1994).

We also do not take into account any time-dependent effect that may occur in the ADAF (see Mineshige 1996). These could be important for explaining spectral variations that are observed during outbursts, but they will have almost no effect on the outer disk, since the ADAF contains such a small amount of mass that it has essentially no dynamical or thermal effect on the much more massive outer disk.

According to the dwarf nova DIM, the accretion rate in a quiescent disk must satisfy the condition

$$\dot{M}(r) < \dot{M}_{crit}(r) = 9.6 \times 10^3 m_1^{1.73} r_2^{-2.6} g s^{-1},$$

where we have taken $\dot{M}_{crit}$ from Ludwig et al. (1994). The disk first becomes unstable at its inner edge when $\dot{M}$ in the disk reaches a critical value near the transition radius. This triggers an inside-out outburst. Since most of the mass evaporation occurs close to the transition radius, the condition for the outburst is equivalent to the requirement $\dot{M}_{ev} = \dot{M}_{crit}$. For $K R_{tr,10}^2 > 1$, we then find

$$R_{tr} = 3.3 \times 10^{10} K^{-0.43} m_1^{0.19} \text{ cm}.$$  

As shown in the previous section, GRO J1655–40 in quiescence requires a transition radius $\sim 10^{10}$ cm, which means that we require a value of $K$ of order a few. In the detailed calculations presented below, we have selected $K = 5$, which gives $R_{tr} = 10^{10}$ cm for the quiescent model just before outburst. In this model, the mass transfer rate feeding the ADAF prior to the onset of the instability is found to be $4.6 \times 10^{16}$ g s$^{-1}$, which is in excellent agreement with the $\dot{M}$ in the ADAF estimated in § 3 by fitting the spectrum of GRO J1655–40 in quiescence ($\dot{M} \sim 3.5 \times 10^{16}$ g s$^{-1}$ for $\Sigma_{ADAF} = 0.3$).

In the presence of evaporation, the usual disk equation for mass conservation must be modified as follows:

$$\frac{\partial \Sigma}{\partial t} + \dot{\Sigma}_{ev} = -\frac{1}{R} \frac{\partial}{\partial R} (R \Sigma v_R),$$

where $\Sigma$ is the surface column density in the disk, and $v_R$ is the radial velocity. Because the evaporation law is independent of $\Sigma$, evaporation results in a disk sharply cut off at the transition radius $R_{tr}$. We thus use as an inner boundary condition the relation

$$\dot{M}_{disk}(R_{tr}) = 0.$$  

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Once the outburst begins, the transition radius moves in and reaches $R_{\text{in}}$, the inner edge of the grid. In order to avoid numerical complications, when this happens we switch to the standard inner boundary condition, $R = R_{\text{in}}$.

In the calculations presented here, the disk inclination was taken to be $70^\circ$, the outer radius of the disk was taken to be $4 \times 10^{11}$ cm, and the inner boundary of the grid was set at $R_{\text{in}} = 4 \times 10^{8}$ cm. Our choice of the inner boundary corresponds to $R_{\text{in}} = 194 R_g$ rather than $3 R_g$, but this is merely for numerical convenience and does not affect any of the results presented here. Once the outburst gets under way and the thin disk extends inward to $R_{\text{in}}$, the time it needs to travel the additional distance to the black hole is quite short compared to the time it took to move from $R = R_t = 10^{10}$ cm down to $R = R_{\text{in}} = 4 \times 10^{8}$ cm. Therefore, very little error is introduced by truncating the numerical simulation at $R_{\text{in}}$. The disk viscosity parameter $\alpha$ in the simulations varies from 0.035 on the cool branch to 0.15 on the hot branch.

Figure 3 displays the initial phases of an outburst seen in the numerical calculations. In this run, several outbursts have already occurred, so that the initial assumed density profile in the disc has been relaxed. The bottom panel shows the magnitude variations $-\Delta m = -m + m_0$ in the $B$, $V$, and $I$ bands, where $m_0$ accounts for the presence of diluting light originating essentially from the secondary (which dominates over the ADAF). For simplicity, we have assumed that $m_0$ is constant and corresponds to a 48 $L_\odot$ secondary with an effective temperature of 6500 K (OB).

During quiescence, the disk is extremely faint—fainter than both the secondary and the ADAF—and contributes less than 1% of the total light; however, in outburst its optical luminosity, although still smaller than that of the secondary, dominates the ADAF, justifying our assumption that $m_0$ is constant.

The $B$, $V$, and $I$ magnitudes decrease simultaneously in the calculations; this is independent of the magnitude of the diluting light $m_0$. The slopes, however, are directly related to $m_0$; for large diluting fluxes, the logarithm appearing in the definition of the magnitude can be linearized, giving $(L(t) \ll L_0)$,

$$-\Delta m = 1.09 \frac{L(t)}{L_0},$$

where $L(t)$ is the disk luminosity and $L_0$ is the luminosity of the companion plus the ADAF. The faster rise in the $B$ band is thus simply due to the fact that $(B - V) = 0$ for a disk in the hot state, whereas most of the diluting light comes from the secondary with $(B - V) \sim 0.5$. This gives a factor of $\sim 1.5$ between the slopes of the $B$ and $V$ magnitudes, as observed.

The middle panel in Figure 3 shows the variation of the transition radius with time. Once the outburst begins, the mass accretion rate increases significantly, and the evaporation is unable to keep up. The transition radius therefore decreases, slowly at first and then more rapidly as the characteristic viscous time decreases with decreasing radius. About five days after the start of the outburst, $R_t$ moves down to $R_{\text{in}}$, the inner boundary of the numerical grid. At later times, we assume that whatever accreted matter reaches $R_{\text{in}}$ will continue down to the black hole in the form of a thin disk, and we calculate the X-ray luminosity accordingly.

The top panel in Figure 3 shows the calculated X-ray light curve. This has been computed assuming that emission from the ADAF has an efficiency of 0.1%, while the matter flowing through $R_{\text{in}}$ in the thin disk has a standard efficiency of 10%. We assumed a conversion factor of $5.3 \times 10^{33}$ ergs count$^{-1}$ to relate ASM counts s$^{-1}$ (2–12 keV) to X-ray luminosity (ergs s$^{-1}$). This simulation shows that it takes about five days for the transition radius to move from its initial value of $10^{10}$ cm to values small enough for X-rays to be emitted, in excellent agreement
The simulations predict that the various optical bands should go into outburst simultaneously, whereas in the observations the outburst occurred first in the $I$ band, followed by the other bands in the order $R$, $V$, and $B$, spread over about a day. The model does not reproduce the delays in the optical bands, but as explained in § 2.2, this delay could result from an earlier increase in the mass transfer rate that occurred less than a month before the outburst, and that might well have caused it. It is also worth noting that our optical light curves deviate significantly from linearity during the first two days, whereas the subsequent evolution shows an almost linear variation of the optical magnitudes. This suggests that it may be somewhat difficult observationally to identify the relative time of outburst in various optical bands.

Both the quiescent mass transfer rate into the ADAF and the inner disk radius are solely determined by the evaporation law. On the other hand, the X-ray delay corresponds to the time it takes to rebuild a standard inner disk, and is thus proportional to the mass of the disk and inversely proportional to the mass transfer rate at the transition radius. This mass transfer rate depends essentially on $a/R$, the ratio of the Shakura-Sunyaev viscosity parameters in the hot and cold states of the disk. Therefore, an increase in $K$ results in a smaller transition radius and a smaller ADAF luminosity; the corresponding shortening of the X-ray delay can in turn be compensated for by taking a smaller $a/R$, i.e., by increasing the viscous time.

The calculations described so far show that a two-component accretion flow model consisting of an outer thin disk and an inner ADAF explains most of the key observations of GRO J1655—40. It explains the quiescent spectrum of the source as well as the characteristics of the outburst, notably the X-ray delay. We argued in § 3 that the quiescent spectrum of GRO J1655—40 cannot be explained by a pure thin disk. We now show that a pure thin disk model has difficulty reproducing the outburst observations. We consider two models for which the thin disk extends down to $R_{\text{in}} = 4 \times 10^8$ cm.

The first model describes an inside-out outburst (Fig. 4, solid lines). The surface density in the outer parts of the accretion disk has been adjusted so as to reproduce the correct slope of the X-ray and optical light curves, the viscosity remaining the same as in the previous model. For this reason, the agreement between the predicted and observed light curves is very good, better in fact than in Figure 3, for which such a fitting procedure was not performed, since there we consider a relaxed case (i.e., after several outbursts) in order to minimize the number of free parameters. The X-ray intensity, which increases simultaneously with the optical flux, is initially quite faint. Thus, the X-ray delay depends on the sensitivity of the X-ray detector; for example, the ASM would be unable to detect the X-rays during the first 1–2 days, and therefore the model predicts an X-ray delay of this order. The delay could be increased further by decreasing the viscosity parameter, which would slow the rise of the X-ray intensity. However, this would cause the rise in the optical flux to be unacceptably slow.

The second model (Fig. 4, dashed lines) describes an outside-in outburst. In order to obtain the longest possible characteristic times, and therefore the most optimistic scenario for the pure disk model, we started the outburst as far out in the accretion disk as possible. In long-period systems like GRO J1655—40, the outer disk is always cold and stable; therefore to trigger the outburst we added some matter at $R = 8 \times 10^{10}$ cm. The outburst then began at $R = 7 \times 10^{10}$ cm. The viscosity was chosen so as to reproduce the observed X-ray delay ($a/R = 0.10$); it takes 2.6 days for the heat front to reach the inner edge of the disk, and an additional 3.1 days for the mass accretion rate to reach $10^{16}$ g s$^{-1}$, the level at which the X-ray flux becomes detectable. This model does reproduce the observed X-ray delay. However, the optical light curves are not in agreement with observations, since most of the disk reaches a hot state before the heat front reaches the inner edge of the disk. Consequently, the optical flux increases too rapidly, on a thermal timescale, and then more slowly, on a viscous timescale. The observed optical light curve does not exhibit such a prominent two-phase behavior, nor does it show such a rapid optical increase, features that are the signatures of an outside-in outburst.

It therefore appears difficult to reconcile the observations with a pure disk model. This might not be an insuperable
difficulty, but it would most probably require making ad hoc assumptions about the density profile and viscosity in the disk. Even if a candidate model could be contrived, it would also be necessary to find an explanation for the quiescent X-ray flux that does not invoke accretion, since the whole disk must be in the low state, and therefore the mass transfer rate must be less than $5 \times 10^{36}$ g s$^{-1}$ at the last stable orbit. Such a small rate of mass transfer corresponds to a luminosity of less than $5 \times 10^{46}$ ergs s$^{-1}$, almost 6 orders of magnitude below the observed quiescent luminosity. Together, the difficulty of building a viable model and the near impossibility of explaining the X-ray flux in quiescence rule strongly against the pure disk model.

5. CONCLUSIONS

We have shown that both the X-ray spectrum observed in quiescence and the 6 day delay between the optical rise and the X-ray outburst in GRO J1655–40 imply that the accretion disk does not extend in its "standard" form all the way down to the last stable orbit. The spectral fitting shows that instead, the inner part of the accretion flow is a hot ADAF. It is this ADAF region that is responsible for the X-ray emission detected by ASCA in quiescence. The outer parts of the disk, located at distances larger than about $10^{15}$ cm in quiescence, are cold and subject to the same thermal and viscous instability as dwarf novae. After the instability has been triggered, a heat front propagates inward. When this heat front reaches the transition radius, the divide between the thin disk and the ADAF region, it cannot propagate farther until it rebuilds the inner part of the disk on a viscous timescale (about a week). Because the efficiency of energy release in the ADAF region is very low, the X-ray outburst starts only when the dense parts of the disk can penetrate far enough in (close to the marginally stable orbit) to allow an efficient transformation of gravitational energy into radiation.

We have also shown that both in quiescence and during the initial outburst, the observed properties of GRO J1655–40 are not consistent with a pure locally cooled thin accretion disk without an ADAF component. Invoking such a model would require (1) accepting that the quiescent X-ray emission is not linked to accretion, and (2) tuning the viscosity and the initial surface density of the disk with some care.

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APPENDIX

NUMERICAL PROCEDURE FOR THE TIME-DEPENDENT ACCRETION DISK

The equations that describe the time-dependent accretion disk are standard (see e.g., Cannizzo 1993b):

$$\frac{\partial \Sigma}{\partial t} = - \frac{1}{R} \frac{\partial}{\partial R} (R \Sigma v_r),$$

and

$$\frac{h}{c_p \Sigma} \frac{\partial T}{\partial t} = \frac{\partial}{\partial R} \left[ \frac{1}{R} \frac{\partial}{\partial R} (R \Sigma h v_r) + \frac{\partial}{\partial R} \left( \frac{3}{2} \frac{R^3 \Sigma v}{R} \right) \right],$$

where $v_r$ is the radial velocity in the disc, $h = (GM)^{1/2}$, and $v$ is the viscosity. We do not take into account the torque due to tidal forces, but assume instead that the outer radius of the disk is fixed at a given value. This is equivalent to assuming that the torque is zero for $R < R_{out}$ and infinite beyond. The two boundary conditions associated with the viscous equations are $\Sigma = 0$ at the inner edge of the disk and $2 \pi R \Sigma v_r = -M_0$ at $R = R_{out}$.

The thermal equation is written as

$$\frac{\partial T}{\partial t} = \frac{2(H - C + J)}{C_p \Sigma} - \frac{\partial}{\partial R} 1 \frac{\partial}{\partial R} (R v_r) - v_r \frac{\partial T}{\partial R},$$

where $H = (9/8) \nu \Sigma \Omega^2$ and $C = \sigma T_e^4$ are the heating and cooling rates, $T_e$ is the effective temperature at the surface of the disk, $C_p$ is the specific heat, $\gamma$ is the perfect gas constant, and $\mu$ is the mean molecular weight. The term $J$ accounts for the radial energy flux carried either by viscous processes or radiation. We use here for $J$ a prescription similar to that of Cannizzo (1993b),

$$J = \frac{1}{R} \frac{\partial}{\partial R} (RF_e), \quad F_e = C_p \Sigma v_r \frac{\partial T}{\partial R} l_e = \frac{3}{2} \nu C_p \Sigma \frac{\partial T}{\partial R},$$

where $v_r$ and $l_e$ are the velocity and size, respectively, of the turbulent eddies responsible for the viscosity. Two boundary conditions must be associated with the thermal equation. However, except across the transition front and at the inner edge, the temperature is equal to the equilibrium value given by $H = C$, so that boundary layers are expected and the two boundary conditions are of no physical importance. To minimize numerical difficulties, we have taken $\partial T/\partial R = 0$ at the inner and outer edges of the disk.
In order to solve these equations, we use a method derived from Eggleton (1971), in which the grid points adapt so that rapid variations of the variables are resolved. This is done by considering that 
\[
\frac{dR}{dq} = \Phi G(R, T, \Sigma, \ldots),
\] 
where \(G\) is a function that becomes small whenever the radial derivative of \(T\) or \(\Sigma\) is large, and \(\Phi\) is a constant that satisfies the differential equation \(d\Phi/dq = 0\). The two boundary conditions for this set of equations (for \(R\) and \(\Phi\)) are the values of \(R\) at the inner and outer edges of the thin disk.

The full set of nonlinear equations is then discretized and solved using a generalized Newton method. This method has the advantage of being fully implicit, but convergence may be a problem when the initial guess is not close to the actual solution. We used grids containing 400 or 800 points.

REFERENCES

Abramowicz, M. A., Chen, X., Kato, S., Lasota, J.-P., & Reguev, O. 1995, ApJ, 438, L37
Bailyn, C. D., Orosz, J. A., McClintock, J. E., & Remillard, R. A. 1995, Nature, 378, 157
Cannizzo, J. K. 1993a, in Accretion Disks in Compact Stellar Systems, ed. J. C. Wheeler (Singapore: World Scientific), 6
Cannizzo, J. K., Chen, W., & Livio, M. 1995, ApJ, 454, 880
Cannizzo, J. K., & Kenyon, S. J. 1987, ApJ, 320, 319
Eggleton, P. P. 1971, MNAS, 151, 351
Gontikakis, C., & Hameury, J.-M. 1993, A&A, 271, 118
Hameury, J.-M., King, A. R., & Lasota, J.-P. 1986, A&A, 161, 71
Harmon, B. A., et al. 1995, IAU Circ. No. 6205
Horne, K., & Marsh, T. R. 1986, MNAS, 218, 761
Hjellming, R. M., & Rupen, M. P. 1995, Nature, 357, 464
Kolb, U., King, A. R., Ritter, H., & Frank, J. 1997, ApJ, 485, L33
Kuulkers, E., Hollander, A., van Paradijs, J., & van Paradijs, J. 1991, A&A, 242, 401
Lasota, J.-P. 1996a, in IAU Symp. 165, Compact Stars in Binaries, ed. J. van Paradijs, E. P. J. van den Heuvel, & E. Kuulkers (Dordrecht: Kluwer), 43
Lasota, J.-P., Narayan, R., & Remillard, R. A. 1995, ApJ, 442, 358
Meyers, F., & Meyer-Hofmeister, E. 1994, A&A, 288, 175
Mineshige, S. 1996, PASJ, 48, 93
Mineshige, S., & Wheeler, J. C. 1989, ApJ, 343, 241
Narayan, R. 1996, ApJ, 462, 136
Narayan, R., Garcia, M., & McClintock, J. E. 1997b, ApJ, 478, L79
Narayan, R., McClintock, J. E., & Yi, I. 1996, ApJ, 451, 821 (NMY)
Narayan, R., & Yi, I. 1994, ApJ, 428, L13
Orosz, J. A., Bailyn C. D. 1997, ApJ, 477, 876 (OB)
Orosz, J. A., Remillard, R. A., Bailyn C. D., & McClintock, J. E. 1997, ApJ, 478, L33 (ORBM)
Pringle, J. E., Verbunt, F., & Wade, R. A. 1986, MNRAS, 221, 169
Robinson, C. R., et al. 1997, in preparation
Remillard, R. A., et al. 1996, IAU Circ. No. 6393
Smak, J. 1981, Acta Astron., 31, 395
Tanaka, Y., & Shibazaki, N. 1996, ARA&A, 34, 607
Tingay, S. J., et al. 1995, Nature, 374, 141
Verbunt, F. 1996, in IAU Symp. 165, Compact Stars in Binaries, ed. J. van Paradijs, E. P. J. van den Heuvel, & E. Kuulkers (Dordrecht: Kluwer), 333
Warner, B. 1995, Cataclysmic Variable Stars (Cambridge: Cambridge Univ. Press)
Wheeler, C. J. 1997, in Relativistic Astrophysics, ed. B. Jones & D. Markovic (Cambridge: Cambridge Univ. Press), in press (astro-ph/9606119)
Wilson, C. A., et al. 1995, IAU Circ. No. 6152
Zhang, S. N., et al. 1994, IAU Circ. No. 6046