ON THE PROPERTIES OF THERMAL DISK WINDS IN X-RAY TRANSIENT SOURCES: 
A CASE STUDY OF GRO J1655−40

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ABSTRACT

We present the results of hydrodynamical simulations of the disk photosphere irradiated by strong X-rays produced in the innermost part of the disk of an accreting black hole. As expected, the irradiation heats the photosphere and drives a thermal wind. To apply our results to the well-studied X-ray transient source GRO J1655−40, we adopted the observed mass of its black hole and the observed properties of its X-ray radiation. To compare the results with the observations, we also computed transmitted X-ray spectra based on the wind solution. Our main finding is that the density of the fast-moving part of the wind is more than 1 order of magnitude lower than that inferred from the observations. Consequently, the model fails to predict spectra with line absorption as strong and as blueshifted as those observed. However, despite the thermal wind being weak and Compton thin, the ratio between the mass-loss rate and the mass-accretion rate is about seven. This high ratio is insensitive to the accretion luminosity, in the limit of lower luminosities. Most of the mass is lost from the disk between 0.07 and 0.2 of the Compton radius. We discovered that beyond this range the wind solution is self-similar. In particular, soon after it leaves the disk, the wind flows at a constant angle with respect to the disk. Overall, the thermal winds generated in our comprehensive simulations do not match the wind spectra observed in GRO J1655−40. This supports the conclusion of Miller et al. and Kallman et al. that the wind in GRO J1655−40, and possibly in other X-ray transients, may be driven by magnetic processes. This in turn implies that the disk wind carries even more material than our simulations predict and as such has a very significant impact on the accretion disk structure and dynamics.

Key words: accretion, accretion disks – hydrodynamics – methods: numerical

Online-only material: color figures

1. INTRODUCTION

Most X-ray sources are powered by disk accretion onto compact objects. Therefore, a main challenge for X-ray astronomy is to understand the mechanisms that enable this process. High quality and high spectral-resolution observations obtained with Chandra, XMM-Newton, and Suzaku allow us to study disk accretion and related outflows better than ever before. These observations are especially revealing if taken of relatively bright, well-studied objects such as X-ray transient sources, with GRO J1655−40 being a very fine example.

Many properties of GRO J1655−40 are well constrained. For example, GRO J1655−40 is a black hole binary at a distance of 3.2 kpc containing a black hole with a mass of 7.0 $M_{\odot}$ that accretes from an F3 IV−F6 IV star with a mass of 2.3 $M_{\odot}$ in a 2.6 day orbit. The inner disk is viewed at an inclination of 67°−85° (nearly edge-on; Orosz & Bailyn 1997). When GRO J1655−40 is in an X-ray-bright phase, it is possible to obtain with Chandra high signal-to-noise spectra as shown by Miller et al. (2006a, hereafter M06). In fact, the quality of the spectra obtained by M06 is good enough to reveal many absorption lines significant at the 5σ level of confidence or higher. Over 70 of these lines can be identified as resonance lines from over 32 charge states. This is in sharp contrast to high-resolution spectra of other systems (see below) and of GRO J1655−40 in other spectral states, when only the absorption lines of Fe XXV and XXVI are detected (e.g., Miller et al. 2004, 2006b; Kubota et al. 2007; Neilsen & Lee 2009).

It is hoped that such high quality spectroscopy can provide new and surprising results. Indeed, absorption lines discovered by M06 are blueshifted and likely produced in a disk wind, making GRO J1655−40 the best case for an X-ray binary with an X-ray absorbing disk wind. But there are also other cases. For example, Chandra and RTXE spectroscopy of the microquasar H 1743−322 reveals blueshifted absorption lines that are likely formed in a disk wind (Miller et al. 2006b). Another example is the X-ray transient 4U 1630−472. Outburst spectra of this source obtained with Suzaku show iron absorption lines indicative of a disk wind (Kubota et al. 2007). More recently, Neilsen & Lee (2009) found an X-ray absorbing disk wind in the microquasar GRS 1915+105. In addition, there is an outflowing X-ray absorber in Circinus X-1, which could either be a disk wind or a wind from a massive companion star (Schulz & Brandt 2002).

The disk wind discovered in GRO J1655−40 gives us a good testbed for constraining wind properties using X-ray observations. For example, by fitting to observations one-dimensional models, M06 and Miller et al. (2008) put strict limits on the ionization balance in the wind. This was facilitated by the fact that some of the detected lines provide a density diagnostic. The main results from the fitting yielded the following parameters: number density $n$ between $5 \times 10^{13} \text{cm}^{-3}$ and $2 \times 10^{14} \text{cm}^{-3}$, column density $N_{H} = 7.4 \times 10^{24} \text{cm}^{-2}$, and velocity $v_r = 500 \text{km s}^{-1}$. The lines show blueshifts in the 300−1600 km s$^{-1}$ range.

Using these wind properties, together with the observed wind speed and the system luminosity, M06 concluded that the inferred wind location is well within the Compton radius (the radius where the gravitational and thermal pressures are equal),...
defined by

\[ R_{IC} = \frac{GM_{BH} m_p \mu}{kT_{IC}}, \]  

where \( M_{BH} \) is the mass of the black hole, \( m_p \) is the proton mass, and \( T_{IC} \) is the Compton temperature (for \( M_{BH} = 7 M_\odot \), \( \mu = 0.6 \), and \( T_{IC} = 1.4 \times 10^7 \) K, the Compton radius is \( R_{IC} = 4.8 \times 10^{11} \) cm). This wind location is then inconsistent with an outflow being driven by thermal expansion, even though weak flows are possible at \( \sim 0.1 \) of the Compton radius, as found by Woods et al. (1996; see also Proga & Kallman 2002, hereafter PK02).

On the other hand, the wind cannot be driven by radiation pressure in this system because the luminosity, \( L_\star \), relative to the Eddington limit, \( L/L_{Edd} = 0.03 \), is too low (e.g., PK02). Therefore, by a process of elimination, M06 concluded that the only plausible mechanisms that could drive the wind are magnetic processes. M06 also concluded that disk accretion itself is driven by magnetic fields.

Although magnetic forces can drive a disk wind (e.g., Blandford & Payne 1982) and magnetic fields are a very important ingredient of accretion disks (Balbus & Hawley 1998), M06’s arguments for magnetic driving are indirect (see also Proga 2006). Ideally, one should demonstrate that magnetic forces can drive a disk wind capable of reproducing the observed spectra. However, that is a relatively challenging task, and instead one could try to verify first whether thermal driving is indeed unsatisfactory. This is a very relevant question especially in light of the work done by Netzer (2006, hereafter N06), who argued that it is possible to produce a simple thermally driven wind that is consistent with the observed properties of GRO J1655−40.

One can argue for or against thermal driving by comparing the escape velocity and the isothermal sound speed at the Compton temperature (equivalently, one can compare the wind launching radius \( R_l \) with \( R_{IC} \) as M06 did). To use this basic physical argument, one needs to estimate the wind temperature and \( R_l \). The former can be constrained relatively well for the Compton heated gas, whereas the latter is much more difficult to constrain. In fact, the controversy about the role of thermal driving has its origin in two different estimates for \( R_l \): M06 estimated \( R_l = 5 \times 10^8 \) cm, whereas N06 estimated \( R_l = 5 \times 10^{10} \) to \( 5 \times 10^{11} \) cm. These two estimates differ by 2–3 orders of magnitude! For the lower estimate, the wind cannot be thermally driven, whereas for the higher estimate it can. We note that the two groups also found different values for \( R_{IC} \) (i.e., M06’s estimate is \( 7 \times 10^{12} \) cm, whereas N06’s estimate is \( 5 \times 10^{11} \) cm). Additionally, the two groups disagree about the gross properties of the thermal wind. In particular, contrary to N06, M06 claimed that both theory and simulations (Begelman et al. 1983; Woods et al. 1996) predict mass-loss rates of thermal winds that are too much low to account for the observed density and blueshift.

The reason for the above disagreement can be traced down to the differences of the absorption lines of Fe xxii and Fe xxiii. N06 assumed saturated lines and the absolute covering factor of the X-ray source to be 0.75, whereas M06 took the lines to be unsaturated and used a covering factor of 1.

Miller et al. (2008) confirmed their previous results and stated that if N06’s claim of the Fe xxii line being saturated were correct then the ratios of some other lines should be different than what is observed. Miller et al. also reiterated their point that Fe xxii lines are optically thin and the covering factor along the line of the sight is near unity. They also stated that N06 used too low a density (by factor of 5–10) as a consequence of his line saturation assumption.

The conclusion of M06 and Miller et al. (2008) was confirmed by Kallman et al. (2009) who used spectral fitting to show that the ionization conditions in a one-dimensional model are not consistent with the wind being driven by thermal expansion.

Generally, M06, N06, and Kallman et al. (2009) agree that thermal driving operates in this system but disagree about the properties of the thermal wind. Overall, it appears that thermal driving is unfavorable. However, this conclusion is based on simplified one-dimensional wind models and needs to be confirmed by a detailed physical model that takes into account the intrinsic multi-dimensional geometry of disk winds.

In this paper, we present 2.5 dimensional axisymmetric, time-dependent hydrodynamical numerical simulations of thermally driven disk winds. We focus on the results obtained after the simulations reached the steady state and compare the results with observations. Our simulations do not include any magnetic processes and therefore do not address the possibility of the wind being magnetically driven, but they instead focus on the thermal contribution to the wind. The main goal of our simulations is to assess within the assumptions whether or not the thermally driven wind can account for the observations. We also consider possible implications of the fact that the mass-loss rate of the wind is several times higher than the total mass-accretion rate.

The paper is organized in the following way. Section 2 describes the methods used; Section 3 lays out the initial conditions of the simulation and discusses the properties of the fiducial run as well as its differences from other runs; Section 4 compares the results to the observations; and Section 5 summarizes and discusses our results.

2. METHOD

2.1. Hydrodynamics

To compute the structure and evolution of a disk irradiated by the central source, we numerically solve the equations of hydrodynamics

\[ \frac{D\rho}{Dt} + \rho \nabla \cdot \mathbf{v} = 0, \]

\[ \rho \frac{D\mathbf{v}}{Dt} = -\nabla P + \rho \mathbf{g}, \]

\[ \rho \frac{D}{Dt} \left( \frac{e}{\rho} \right) = -P \nabla \cdot \mathbf{v} + \rho \mathbf{L}, \]

where \( \rho \) is the mass density, \( P \) is the gas pressure, \( \mathbf{v} \) is the velocity, \( e \) is the internal energy density, \( \mathbf{L} \) is the net cooling rate, and \( \mathbf{g} \) is the gravitational acceleration of the central object. We adopt an adiabatic equation of state, \( P = (\gamma - 1) e \), and consider models with the adiabatic index \( \gamma = 5/3 \).

We use the ZEUS-2D code (Stone & Norman 1992) extended by Proga et al. (2000) to solve Equations (2)–(4). We perform our calculations in spherical polar coordinates \((r, \theta, \phi)\) assuming axial symmetry about the rotational axis of the accretion disk \((\theta = 0^\circ)\).

Our computational domain is defined to occupy the angular range \(0^\circ \leq \theta \leq 90^\circ\) and the radial range \(r_1 \leq r \leq r_\alpha\). The numerical resolution consists of 200 cells in the \( r \) direction and 100 cells in the \( \theta \) direction. In the angular direction, we used the following ratios: \( d\theta_{k+1}/d\theta_k = 0.95, 0.97, 0.99, \) and 1.00 (i.e., the zone spacing increases toward the pole). Gridding in this manner ensures good spatial resolution close to the disk. In the
r direction, we used $dr_{k+1}/dr_k = 1.04$, which enables a good resolution at smaller radii.

We chose the boundary condition at the pole (i.e., $\theta = 0^\circ$) as an axis-of-symmetry boundary condition. At the disk (i.e., $\theta = 90^\circ$), we applied the reflecting boundary condition. For the inner and outer radial boundaries, we used an outflow boundary condition (i.e., to extrapolate the flow beyond the boundary, we set values of variables in the ghost zones equal to the values in the corresponding active zones; see Stone & Norman 1992 for more details).

2.2. Radiation Field

We deal with the radiation field and radiation heating and cooling in the same manner as described by Proga & Kallman (2002; see also Proga et al. 2000).

The net cooling rate is a function of the photoionization parameter, which is defined as

$$\xi = \frac{4\pi F_X}{n},$$

(5)

where $F_X$ is the local X-ray flux and $n = \rho/(\rho_p \mu)$ is the number density of the gas. We consider a fully ionized gas with $\mu = 0.6$. The local X-ray flux is corrected for the optical depth effects,

$$F_X = F_\star \exp(-\tau_X),$$

(6)

where $\tau_X$ is the X-ray optical depth and

$$F_\star = \frac{L_\star}{4\pi r^2},$$

(7)

with $L_\star$ being the luminosity of the central source.

We estimate $\tau_X$ between the central source and a point in the flow from

$$\tau_X = \int_0^r \kappa_X \rho(r, \theta) dr,$$

(8)

where $\kappa_X$ is the absorption coefficient and $r$ is the distance from the central source. We assume $\kappa_X = 0.4$ g cm$^{-2}$ which is numerically the value of electron scattering. We found that in most cases, the wind is optically thin and the wind column density,

$$N_\theta(\theta) = \int_{r_i}^{r} \frac{\rho(r, \theta)}{\mu m_p} dr,$$

(9)

is less than $10^{23}$ cm$^{-2}$.

3. RESULTS

To complete the specification of the simulations, we need to assign the properties of the disk atmosphere (the lower boundary condition for the wind solution). For the density, we use a simple distribution at $\theta = 90^\circ$ of the form $\rho = \rho_0 (r/r_{IC})^{-\alpha}$, where $\alpha$ and $\rho_0$ are constants. We run the model for $\alpha = 0, 1, 2, 3$ and for $\rho_0 = 10^{-12}, 10^{-11}$, and $10^{-10}$ g cm$^{-3}$ (see Table 1). The temperature is set to $10^4$ K, while the radial velocity is set to zero. In addition, we enforce Keplerian rotation at $\theta = 90^\circ$. The size of the computational domain is defined in the following way: $r_i = 0.05 R_{IC}$, while $r_o = 20 R_{IC}$. We follow M06 and N06 and adopt the following properties of the system: the mass of the central black hole $M_{BH} = 7 M_\odot$, the luminosity $L_\star = 0.03$ in units of the Eddington limit, $L_{Edd}$, and the Compton temperature, $T_{IC} = 1.4 \times 10^7$ K. The adopted $L_\star$ corresponds to the mass-accretion rate $M_\star = 4.4 \times 10^{17}$ g s$^{-1}$.

3.1. Properties of the Fiducial Run

Our fiducial run (C8) has the following parameters: $\alpha = 2$, $\rho_0 = 10^{-11}$ g cm$^{-3}$, $d\theta_{k+1}/d\theta_k = 0.97$, and $dr_{k+1}/dr_k = 1.04$. This run settles to a steady state after about three sound-crossing times.

We show the geometry and structure of the fiducial run in Figure 1. The top left panel shows a density distribution with a significant departure from spherical symmetry. The temperature (the top right panel) is close to 0.1 $T_{IC}$ in the outer part of the wind. In a very narrow region around the rotational axis, the temperature is much less than $T_{IC}$ at $r > 16 R_{IC}$, whereas at small radii the temperature is comparable to $T_{IC}$. The X-ray irradiated gas accelerates rapidly and becomes supersonic relatively close to the disk. As shown in the bottom left panel, the contour for the Mach number of one corresponding to the wind

| No. of Run | $\alpha$ | $\rho_0$ (g cm$^{-3}$) | $d\theta_{k+1}/d\theta_k$ | $\rho(\theta_o)$ (g cm$^{-3}$) | $M_{wind}$ (10$^{18}$ g s$^{-1}$) | $\theta_{max}$ (deg) | $v_\theta$ (km s$^{-1}$) | $\gamma$ (2.2 $\times$ 10$^5$ s$^{-1}$) |
|-----------|---------|-----------------------|--------------------------|-------------------------------|-----------------------------|---------------------|---------------------|---------------------|
| A         | 0       | $10^{-10}$            | 0.97                     | $10^{-9}$                      | 0.33                        | 59                  | 587                 | 10                  |
| B         | 1       | $10^{-10}$            | 0.97                     | $10^{-9}$                      | 0.33                        | 67                  | 599                 | 10                  |
| C1        | 2       | $10^{-12}$            | 1.00                     | $10^{-10}$                     | 6.6                         | 90                  | 746                 | 10                  |
| C2        | 2       | $10^{-12}$            | 0.99                     | $10^{-10}$                     | 6.8                         | 90                  | 708                 | 10                  |
| C3        | 2       | $10^{-12}$            | 0.98                     | $10^{-10}$                     | 6.9                         | 90                  | 675                 | 10                  |
| C4        | 2       | $10^{-12}$            | 0.97                     | $10^{-10}$                     | 7.0                         | 90                  | 642                 | 10                  |
| C5        | 2       | $10^{-12}$            | 0.96                     | $10^{-10}$                     | 7.1                         | 90                  | 619                 | 10                  |
| C6        | 2       | $10^{-12}$            | 0.95                     | $10^{-10}$                     | 7.1                         | 90                  | 604                 | 30                  |
| C7        | 2       | $10^{-11}$            | 0.99                     | $10^{-9}$                      | 0.32                        | 90                  | 703                 | 10                  |
| C8        | 2       | $10^{-11}$            | 0.97                     | $10^{-9}$                      | 0.33                        | 90                  | 627                 | 10                  |
| C9        | 2       | $10^{-11}$            | 0.95                     | $10^{-9}$                      | 0.33                        | 90                  | 587                 | 10                  |
| C10       | 2       | $10^{-10}$            | 0.99                     | $10^{-8}$                      | 0.33                        | 63                  | 692                 | 30                  |
| C11       | 2       | $10^{-10}$            | 0.97                     | $10^{-8}$                      | 0.33                        | 62                  | 603                 | 30                  |
| C12       | 2       | $10^{-10}$            | 0.95                     | $10^{-8}$                      | 0.33                        | 63                  | 553                 | 10                  |
| D1        | 3       | $10^{-11}$            | 0.97                     | $10^{-7}$                      | 0.11                        | 90                  | 670                 | 10                  |
| D2        | 3       | $10^{-10}$            | 0.97                     | $10^{-7}$                      | 0.32                        | 90                  | 621                 | 30                  |

Note. * $2.2 \times 10^5$ is the sound-crossing time $(r_o/c_s)$, calculated for $T = T_{IC}$. 

Table 1

Summary of the Results

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launched at relatively large radii is almost a straight line making an angle of about 15° with the disk midplane. Streamlines (the bottom right panel) show that most of the disk wind follows almost perfect straight lines inclined at the angle of about 45° with respect to the disk midplane. The streamlines are curved only in zones near the rotational axis and very close to the disk. It appears that the outer wind is self-similar (we will return to this point at the end of this subsection). The streamlines also show that the flow expands in the angular direction, reminiscent of a spherical outflow, only for streamlines originating at \( r \ll 3R_{\text{IC}} \).

An accretion disk is geometrically extended and as such it intercepts the central radiation over a large range of radii. However, the thermal expansion drives a significant wind only within a relatively narrow radial range. One can show this wind property by plotting the product of the density and the velocity normal to the disk as a function of the radius along the disk midplane (see Figure 2). Most of the outflow in the fiducial run comes from the narrow ring on the disk in the zone between 0.07 \( R_{\text{IC}} \) and 0.2 \( R_{\text{IC}} \), having its maximum at \( \sim 0.1 R_{\text{IC}} \). For radii \( > 0.2 R_{\text{IC}} \), the mass flux density scales like \( r^{-q} \), with \( q = 1.76 \).

The wind is launched in a non-spherical way and it remains non-spherical as it expands. To illustrate this point, Figure 3 shows various flow properties at \( r_o \) as functions of \( \theta \). In particular, the top panel shows the product of density and radial velocity (solid line) and the accumulated mass-loss rate (dotted line). We evaluate the accumulative mass-loss rate throughout \( r_o \) using the following formula:

\[
\dot{m}(\theta) = 4\pi r_o^2 \int_0^\theta \rho v_r \sin \theta d\theta.
\]

One finds that most of the mass flows out of the computational domain through the outer spherical boundary over a broad range of angles between 20° and 75°. The outflow for large \( \theta \) contributes insignificantly. The total mass-loss rate for the fiducial run is \( \dot{M}_w = \dot{m}(90^\circ) = 3.3 \times 10^{18} \text{ g s}^{-1} \). Note that this is larger than the assumed accretion rate \( \dot{M}_a \) by a factor of 7.5. We will return to this point in Section 5.

The middle panel of Figure 3 shows that the photoionization parameter (solid line) decreases with increasing \( \theta \), which is

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**Figure 1.** Fiducial run C8. Top left panel: color density map. Top right panel: color temperature map. Bottom left panel: the Mach number \( M \) contours (the poloidal component only). Contours are for \( M = 1, 2, \) and 3 (bottom to top). Bottom right panel: the flow streamlines. Note the self-similarity of the streamlines especially those arising from the outer disk. In all panels, the rotation axis of the disk is along the left-hand vertical frame, while the midplane of the disk is along the lower horizontal frame. Lengths are expressed in units of the Compton radius, \( R_{\text{IC}} \).

(A color version of this figure is available in the online journal.)

**Figure 2.** Mass flux density, \( \rho v_r \), as the function of radius along the disk midplane, \( \theta = 90^\circ \) for the fiducial run.
another indication that the wind is not spherically symmetric. In particular, $\xi$ changes from $10^2$ to $10^3$, whereas the column density (dotted line) increases from below $10^{21}$ cm$^{-3}$ to $10^{23}$ cm$^{-3}$.

The bottom panel of Figure 3 presents the radial velocity (solid line) and the number density (dotted line) versus $\theta$. It is evident from this panel why the zones close to the disk and the rotational axis contribute insignificantly to the total mass-loss rate: the zone near the disk is very dense but also very slow, while the zone near the axis of rotation has extremely low density, so that the product of the density and radial velocity is very small.

To specify the departure of the disk wind from a purely radial wind, Figure 4 shows $\xi$ (top panel) and $n$ (bottom panel) as a function of radius for various $\theta$. In the case of an optically thin radial wind with a constant velocity, $\xi$ is constant and $n \propto r^{-2}$. In our solution, $\xi$ decreases with increasing radius especially for small and intermediate $\theta$, whereas $n$ decreases slower than $r^{-2}$ and can even increase with radius for $\theta < 62^\circ$.

How does our solution compare to a $\beta$-velocity law, in terms of approaching a constant value at large $r$? As expected, our solution is in agreement with a $\beta$-velocity law for large radii. Figure 5 shows an example of the radial velocity as a function of radius for $\theta = 76^\circ$ (this corresponds to the dot-dashed lines in Figure 4). The radial velocity increases up to the Compton radius and then remains roughly constant. Also, there is a visible dip at $0.4R_{\text{IC}}$, which is a consequence of the radial line crossing different streamlines. The function has a nearly flat shape for radii larger than $1.0R_{\text{IC}}$.

We conclude that a thermal disk wind cannot be approximated by a radial outflow. Font et al. (2004) found that a disk wind expands quasi-spherically if a relatively steep decline of the mass flux density with radius is assumed ($\rho v_\theta \propto r^{-q}$, with $q \gtrsim 2.5$). As we mentioned above, in our self-consistent simulations, $q \approx 1.76$ in the self-similar part of the wind. It is unfortunate that the wind expansion is non-radial because a radial outflow is a very simple case that can be (and has been) straightforwardly implemented into photoionization/spectra calculations. Is it then possible to approximate the thermal wind with some other simple model? It is beyond the scope of this paper to answer this question.

### 3.2. Comparison of the Fiducial Run with Other Runs

To check the robustness of the solution presented above, we explored effects of the numerical resolution and the density along the lower boundary. We present the most significant properties of those runs in Table 1. The columns are organized in the following way: (1) the name of the run; (2) the $\alpha$ parameter used in the density profile; (3) $R_{\text{IC}}$, the normalized density at the lower boundary of the computational domain; (4) $d\theta_{\text{max}}/d\theta$, the angular resolution; (5) $\rho(r_{\text{in}})$, the disk density at the inner radius; (6) $M_w$, the total wind mass-loss rate; (7) the maximal angle $\theta_{\text{max}}$ for which the integration of the total wind mass-loss is computed ($M_w = d\dot{m}(\theta_{\text{max}})$); (8) the maximal radial velocity on the outer shell; and (9) the total time of the run.
As mentioned in Section 3, the density along the lower boundary ($\theta = 90^\circ$) is specified by two parameters, $\alpha$ and $\rho_0$. We have run several models for various $\alpha$ and $\rho_0$ to check if our solution depends on the density along the disk. As for radiation driven disk winds (see, e.g., Proga et al. 1998), we find that as long as the density along the disk is high enough, the gross properties of the thermal disk wind do not depend on the assumed disk density.

For a relatively small density along the disk, the X-rays heat the gas to the Compton temperature even at $\theta = 90^\circ$. This means that the computational domain does not capture a cold disk and its part that is in hydrostatic equilibrium. Consequently, the simulations do not follow the transition between a cold and a hot flow. Therefore, the mass flux density from the lower boundary is set not by the physics of a cold disk being heated by X-rays but by a choice of the density. Comparing $M_w$ for runs C4, C8, C11, and D2, we find that the mass-loss rate depends on the density if the density is too low. Only when the density along the disk is high enough that the central X-rays do not heat the gas along the lower boundary do the simulations capture the transition from the cold to hot phase. The mass-loss rate then becomes independent of a particular choice of the density in the cold part of the disk.

However, we also find that for very high densities along the disk there is a technical problem with the proper measurement of $M_w$. Our goal is to compute the wind from a geometrically thin disk. But if we choose too high a density along the lower boundary, the disk will flare and its thickness at large radii becomes substantial. During the simulations, the dense disk remains cold and nearly in hydrostatic equilibrium but it will subsonically fluctuate. Therefore, when computing the wind mass-loss rate, one needs to keep in mind these fluctuations of the dense disk and exclude the region very close to the disk midplane. Otherwise the subsonic oscillations of the very dense disk material would be counted as the disk outflow and this will yield an erroneous estimate of the wind mass-loss rate. In practice, when calculating the total outflow rate for some runs, we stop the integration in Equation (10) at $\theta_{\text{max}} < 90^\circ$, where the transition from the wind to the disk occurs. This transition is not difficult to identify as it occurs where the density sharply increases with $\theta$ (see the bottom panel in Figure 3). The seventh column in Table 1 lists the maximum angles we used.

For the fiducial run as well as some other runs, the region near the disk midplane does not affect estimations of $M_w$. This is one of the reasons we chose the run C8 with $\rho_0 = 10^{-11}$ g cm$^{-3}$ as the most suitable run in terms of having not too high a density, as only a small portion of the dense disk will enter the domain. On the other hand, the density cannot be so small as to make the transition from the disk to the wind outside the computational domain. All runs with the density comparable to or larger than that in the fiducial run have a total mass-loss rate of about $3.3 \times 10^{17}$ g s$^{-1}$. Other wind properties are also similar (see Table 1).

As for the resolution study, we find that changes in the $d\theta_{\text{res}}/d\theta_k$ factor produced small differences in the final steady state and such effects just slightly changed calculated physical quantities and properties of the solution.

4. COMPARISON WITH OBSERVATIONS

We start our comparison with the observations by considering the wind properties inferred from the observation and then we present an example of the synthetic spectrum directly compared with the observed X-ray spectrum. As we described in Section 1, the number density of the wind is inferred to be quite high, i.e., $5 \times 10^{15}$ cm$^{-3}$.

Figure 6 shows the scatter plot of the photoionization parameter versus $v_r$ for the fiducial run. Dots correspond to $n < 10^{12}$ cm$^{-3}$, whereas diamonds correspond $n \geq 10^{12}$ cm$^{-3}$. The lack of points corresponding to $v_r > 100$ km s$^{-1}$ means that the thermal wind cannot account for a high density fast outflow as observed in GRO J1655—40 during an X-ray-bright state.

As another test of our model, we computed synthetic spectra using the photoionization code XSTAR (Kallman & Bautista 2001; Bautista & Kallman 2001). Our method of performing this calculation is very similar to that used and described in Dorodntysyn et al. (2008). We have computed spectra of the fiducial model for several inclination angles. We have also calculated the emissivity and opacity at each point in the hydrodynamic flow at the end of the simulations and then integrated the formal solution of the transfer equation along radial rays to get the spectrum. In doing this, we ignore emission. Since emission will tend to fill in absorption lines, thus reducing their strength, these models represent an overestimate to the amount of absorption coming from such a flow. For most inclinations, the spectra show no absorption features or only very weak ones. Figure 7 compares the count spectrum observed by the HETG during an X-ray-bright phase of the GRO J1655—40 2005 outburst (see M06 for more details) with an example of our synthetic spectrum.

For this figure, we have chosen an almost $90^\circ$ inclination because it yields the best fit to the data and the right ionization for the most lines. Nonetheless, even for this very high inclination, the range of ionization is too narrow, and the ionization is probably lower than that observed. It is possible to obtain higher ionization for lower inclinations, but then the column density is lower, making the fit is even worse. To see if we can improve the fit for a given density structure, we have to put in the 600 km s$^{-1}$ blueshift by hand while computing this spectrum (note that our wind model does not predict such a high velocity of the dense wind at this inclination angle). However, neither this nor any other simple model fits the Fe Kα line profile showing the high velocity ($\sim 1600$ km s$^{-1}$) component. We also found that the curve of growth of the lines does not fit: there are no saturated...
Figure 7. Comparison of the X-ray spectrum of GRO J1655−40 with the model spectrum computed based on our thermal disk wind (see Section 4). The points with the error bars are the data and the model is the red curve. The wavelength is in angstroms.

(A color version of this figure is available in the online journal.)

lines in the model. Still, close to half of the observed lines are predicted in the model, including the Fe XXII lines.

In addition to an inadequate mass-loss rate, the model shows a geometric structure different from that required by the spectra of GRO J1655−40. From the lack of emission lines, M06 and N06 concluded that the global covering factor of the wind is small, while Figure 3 shows that the density is reasonably constant over a broad range of angles.

5. SUMMARY AND DISCUSSION

In this paper, we presented axisymmetric, time-dependent hydrodynamical numerical simulations of thermally driven disk winds. To apply our results to GRO J1655−40, we adopted the observed mass of its black hole and the properties of its X-ray radiation. We performed the simulations using the same code and in a similar fashion to PK02 who computed wind for low-mass X-ray binaries. The main difference is that we turned off radiation driving here because, as found by PK02, it is negligible for the stellar black hole accretors. Turning off radiation driving makes the simulations run much faster. The main goal of our simulations is to check whether, in the absence of magnetic fields and within the given simplifications and assumptions, the thermally driven wind can account for the observation.

To compare the results with the observations, we also computed transmitted X-ray spectra using the photoionization code XSTAR and our wind solution. Our main findings are as follows.

1. The density of the fast-moving part of the wind is more than 1 order of magnitude lower than that inferred from the observations. Consequently, contrary to the claim made by N06, the thermal model fails to predict synthetic spectra with line absorption as strong and as blueshifted as those observed. Overall, our results support the conclusion reached by Miller et al. (2008) and Kallman et al. (2009) that GRO J1655−40 and likely other X-ray transient sources have thermal winds insufficient to explain the observed spectra.

2. Despite the thermal wind being weak and Compton thin, the ratio between the mass-loss rate and the mass-accretion rate is about seven.

3. We discovered that the outer wind is self-similar.

One should ask if our simulations are conclusive. Namely, can we rule out thermal driving as a mechanism responsible for the observed wind even though we did not consider some effects such as thermal conduction? Conduction could increase
the degree of disk heating and could in principle increase the gas density above the disk photosphere. However, it is unclear if this increase would lead to an increase in the density at the wind base and consequently in an increase of the density of the fast wind, which is what is needed to account for the observations. The density of the wind base could also be sensitive to the details of radiative transfer and spatial resolution. Hydrostatic models of the ionization layer on the surface of the disk by Jimenez-Garate et al. (2002, 2005) showed that there is much more opacity than the electron scattering opacity in the very thin layer that produces the UV and most of the X-ray emission on top of the optically thick disk. Therefore, our pure electron scattering assumption might underestimate the opacity. However, if this were true then that would mean that thermal driving is even less efficient because the irradiation would penetrate the disk even less. We expect that the mass-loss rate is set at the sonic point, which would occur in the million Kelvin temperature plateau seen in the hydrostatic models. In that region, the models indicate that photoabsorption opacity due to Fe XXV and Fe XXVI is comparable to the electron scattering opacity. Radiative heating can also be affected by disk flaring or the shape of the disk, in general. For example, at a given radius a strongly flared disk might intercept more central radiation than a flat disk. Consequently, the resulting outflow might be more collimated and denser. We intend to extend our model to explore these effects in the near future. However, before that we plan to check if the inclusion of magnetic driving will produce a dense fast wind.

It is important to remember that the outburst spectra of GRO J1655–40 and GRS 1915+05 are exceptional. The spectra of these objects in their low states, as well as the spectra of other black hole binaries, show absorption lines of only Fe XXV and Fe XXVI. That suggests lower densities and mass-loss rates, and it raises the question of whether thermally driven winds might account for those observations. The microquasar H1743–322 (Miller et al. 2006) is a good example, with Fe XXV and Fe XXVI equivalent widths of up to 4.3 mÅ and 6.8 mÅ, respectively. The densities are below 10¹³ cm⁻³ and the ionization parameters are above 10⁵ in four separate observations. We have examined our models at sightlines farther from the disk, where the velocities match those reported by Miller et al. (2006). The Fe XXV and Fe XXVI column densities are too small to account for the observed equivalent widths by a factor of a few, and we conclude that thermally driven winds cannot account for the typical low state winds.

In the models presented here, we found that \( M_\infty / M_a \approx 7 \). This is a fairly large ratio, which could mean that the thermally driven wind can significantly change the mass flow in the disk. Therefore, we decided to check if our result is consistent with the results of others who modeled thermally driven winds. Of particular interest is the work of Woods et al. (1996), who studied thermally driven disk winds in great detail and performed many axisymmetric wind simulations for various luminosities. They summarized their results in an analytic formula which fitted the mass flux density distribution obtained from simulations for various \( L_\ast \) (their Equation (5.2)).

To find the \( M_\infty / M_a \) predicted by Woods et al.’s simulations, we integrated their formula over the disk surface to obtain \( M_\text{wind} \) and then expressed \( L_\ast \) through the following equation of the accretion luminosity:

\[
L_\ast = \frac{G M_\text{BH} M_a}{2 r_s},
\]

where \( r_s \) is the Schwarzschild radius, which yields a relation between \( M_\infty \) and \( M_a \). We found that the Woods et al. simulations predict \( M_\infty / M_a \) in a range from 2.0 to 6.0 depending on the luminosity. The \( L_\ast \) dependence is weak for \( L_\ast / L_{\text{Edd}} < 0.1 \). Thus, \( M_\infty / M_a \) is higher than, but comparable to, Woods et al.’s results for \( L_\ast \) on the lower side of their luminosity range. Nevertheless, both sets of simulations predict \( M_\infty / M_a > 1 \).

Our conclusion from this analysis is that \( M_\infty \) is greater than the rate at which the central engine is fueled. This in turn can make the disk variable. One could even expect this relatively strong wind to cause a recurrent disappearance of the inner disk even if the disk is fed at a constant rate at large radii. Neilsen & Lee (2009) suggested such a process while interpreting jet/wind/radiation variability in the microquasar GRS 1915+105. It is beyond the scope of this work to model the effect of the wind on the disk. However, this problem was studied by Shields et al. (1986), who showed that due to viscous processes no oscillations appear for \( M_\infty / M_a < 15 \). We conclude that the thermal wind is too weak to cause the oscillation. But the thermal wind is also too weak to account for the observed wind. Therefore, it is possible that in GRO J1655–40 and other sources, e.g., GRS 1915+105, the wind responsible for the observed X-ray absorption will be so strong that \( M_\infty / M_a > 15 \) and as such cause disk oscillations and contribute to the observed disk variability.

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