The radial effective temperature distribution of steady-state, mass-losing accretion disks

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ABSTRACT

Mass loss appears to be a common phenomenon among disk-accreting astrophysical systems. An outflow emanating from an accretion disk can act as a sink for mass, angular momentum and energy and can therefore alter the dissipation rates and effective temperatures across the disk. Here, the radial distributions of dissipation rate and effective temperature across a Keplerian, steady-state, mass-losing accretion disk are derived, using a simple, parametric approach that is sufficiently general to be applicable to many types of dynamical disk wind models.

Effective temperature distributions for mass-losing accretion disks in cataclysmic variables are shown explicitly, with parameters chosen to describe both radiation-driven and centrifugally-driven outflows. For realistic wind mass-loss rates of a few percent, only centrifugally-driven outflows – particularly those in which mass loss is concentrated in the inner disk – are likely to alter the disk's effective temperature distribution significantly. Accretion disks that drive such outflows could produce spectra and eclipse light curves that are noticeably different from those produced by standard, conservative disks.

Key words: accretion, accretion disks — binaries: close — stars: mass loss — novae, cataclysmic variables

1 INTRODUCTION

Accretion disks are an important ingredient in our current understanding of many astrophysical systems on all scales. Examples of presumed disk accretors include young stars, compact objects in close binary systems and active galactic nuclei (AGN) and quasars (QSOs). There is also evidence that the process of mass accretion via a disk is often – and perhaps always – associated with mass-loss from the disk in the form of a wind or a jet. For instance, the broad-line region in QSOs, the radio jets in AGN, the jets and bipolar outflows from young stars and the fast winds observed in disk-accreting cataclysmic variables all provide compelling evidence for such a disk-wind connection [see Livio (1999) for a recent review].

Since the outflow provides an additional sink for mass, angular momentum and energy, the spectrum emitted by a mass-losing accretion disk will in general be different from that produced by a "conservative" one. Intuitively, mass loss may be expected to reduce the effective temperatures across the entire disk, but especially in those regions from which matter is preferentially expelled. Based on this expectation, it has been suggested, for example, that mass loss from the inner disk may be responsible for the apparently flatter-than-expected radial brightness temperature distributions inferred for some disk-accreting nova-like CVs (Rutten, vanParadijs & Tinbergen 1992).

Despite this, few attempts have been made to generalize steady-state accretion disk theory to include the family of mass-losing disks. As a result, the standard $T \propto R^{-3/4}$ law that is valid for conservative disks is almost invariably used in practice (i.e. in modeling and in comparisons between theory and observations), even in the context of disk-accreting systems in which mass loss could alter the disk's effective temperature distribution significantly. Where modifications to the standard picture have been attempted, these have usually been ad hoc (e.g. Shlosman, Vitello & Mauche 1996). An exception is the work by

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Piran (1977), which provides analytic expressions (in integral form) for the energy generation rate in a mass-losing accretion disk. However, the bulk of Piran’s paper deals specifically with thermally-driven, evaporative disk winds. The only other type of mass-loss considered explicitly by him is a “delta-function outflow” from a particular radius on the disk.

The purpose of the present work is twofold. First, simple, analytical expressions are derived for the radial distributions of dissipation rate (Section 2) and effective temperature (Section 3) across a steady-state accretion disk that is losing mass and angular momentum to an outflow from its surface. The derivation is designed to closely parallel standard, steady-state accretion disk theory, as described, for example, by Pringle (1981). The parametric model adopted to describe the mass loss is simple, yet sufficiently general to be applicable to many types of dynamical disk wind models, including both radiation-driven outflows (e.g. Murray & Chiang 1996; Proga, Stone & Drew 1998; Feldmaier & Shlosman 1999; Feldmaier, Shlosman & Vitello 1999) and centrifugally-driven, magneto-hydrodynamic (MHD) disk winds (e.g. Blandford & Payne 1982; Cannizzo & Pudritz 1988; Emmering, Blandford & Shlosman 1992; Pelletier & Pudritz 1992).

Second, the formalism is applied to accretion disks in cataclysmic variable stars (CVs; Section 3). The mechanism responsible for the mass loss from these systems is currently still unknown, so both radiation-driven (Section 5.3) and centrifugally-driven outflows (Section 5.2) are considered as possible culprits. It is found that if CV winds are of the latter type, they could leave an observable imprint on the disk’s temperature profile.

2 THE VISCOUS DISSIPATION RATE IN A STEADY-STATE, MASS-LOSING DISK

It is convenient to start by defining the cumulative mass-loss rate from the disk as

$$
\dot{M}_w(R) = \int_{R_*}^{R} \dot{m}_w(R') \, 4\pi R' \, dR'.
$$

Here $R_*$ denotes the radius at the inner edge of the disk ($R_* \approx R_{WD}$ in non-magnetic CVs, where $R_{WD}$ corresponds to the radius of the white dwarf [WD] at the centre of the disk) and $\dot{m}_w(R)$ is the mass-loss rate per unit area from each disk face. With this definition the total mass-loss rate is given by

$$
\dot{M}_{w,\text{total}} = \dot{M}_w(R_{\text{disk}}),
$$

where $R_{\text{disk}}$ is the radius of the accretion disk. The radial rate of change of $\dot{M}_w(R)$ is just

$$
\frac{\partial \dot{M}_w(R)}{\partial R} = 4\pi R \, \dot{m}_w(R).
$$

The equations of mass and angular momentum conservation for a thin disk subject to mass and angular momentum loss can now be derived in close analogy to the usual case by considering a particular annulus on the disk with inner radius $R$ and outer radius $R + \Delta R$ (Pringle 1981; Frank, King & Raine 1992).

The mass of this annulus is $2\pi R \Delta R \Sigma$, where $\Sigma$ is the surface density of the disk – and the time rate of change of this is equal to the rate at which mass flows into the annulus from its radial borders minus the rate at which it loses mass to the wind. In the limit $\Delta R \to 0$, mass conservation then gives

$$
\frac{\partial \Sigma}{\partial t} = -\frac{1}{R} \frac{\partial}{\partial R} (R \Sigma \dot{m}_w) - \frac{1}{2\pi R} \frac{\partial \dot{M}_w}{\partial R},
$$

where $\dot{V}_R$ is the radial inflow velocity of material in the disk ($\dot{V}_R < 0$).

In order to write down the angular momentum equation, it is necessary to specify how, and how much, angular momentum is lost to the wind. Here, it will be assumed that matter ejected at radius $R$ on the disk carries away specific angular momentum $(lR)^2\Omega$, where $\Omega$ is the angular velocity in the disk at $R$. Thus $l = 0$ corresponds to a non-rotating disk wind and $l = 1$ to outflowing material which retains the specific angular momentum it had at the point of ejection. Models in which the lever-arm $l > 1$ effectively describe an outflow with head-on-a-wire rotation out to a radius $R_{\text{out}} = lR$, as envisaged for centrifugally-driven magnetic disk winds. Models such as this can remove a lot of angular momentum from the disk.

With the angular momentum carried by the outflowing material thus fixed, conservation of angular momentum gives

$$
\frac{\partial \Sigma R^2 \Omega}{\partial t} = -\frac{1}{R} \frac{\partial}{\partial R} (R^2 \Sigma \dot{V}_R \Omega) + \frac{1}{R} \frac{\partial}{\partial R} (R^2 \nu \partial \Omega/\partial R) - \frac{(lR)^2 \Omega}{2\pi R} \frac{\partial \dot{M}_w}{\partial R}.
$$

The first two terms on the right hand side of this expression describe the inflow of angular momentum through the boundaries of the annulus and the effects of viscous torques due to shear ($\nu$ here is the effective kinematic viscosity). The third term allows for the angular momentum sink provided by the outflow. Expressions closely analogous to Equations 4 and 5 were first written down by Bath, Edwards & Mantle (1983), who examined the process of stream penetration, i.e. the effects on the disk of a radially varying mass input with constant specific angular momentum.
To keep the discussion simple, the disk will now be assumed to be Keplerian and in a steady-state. The first of these assumptions gives \( \Omega(R) = (GM_\ast/R^3)^{1/2} \), where \( M_\ast \) is the mass of the accretor, whereas the second implies that \( \frac{d\Omega}{dR} = 0 \) in (10) and (11). The conservation equations can then be integrated, which for \( \Omega \) gives

\[
-2\pi R\Sigma R\Omega = \dot{M}_w(R) = \text{constant}.
\]

The first term on the left hand side of this expression is clearly the accretion rate at radius \( R \) on the disk, \( \dot{M}_{\text{acc}}(R) \). The constant of integration in (10) can be found by considering the boundary condition at \( R = R_\ast \), which shows that it is simply the rate of accretion onto the central star, i.e.

\[
\dot{M}_{\text{acc}}(R) = \dot{M}_{\text{acc}}(R_\ast) + \dot{M}_w(R).
\]

The angular momentum integral is somewhat more difficult to compute, since as a result of the factor of \( R^2\Omega \) in the new outflow sink term, the right hand side of Equation (11) is no longer a perfect differential. That term, at least, must therefore be integrated explicitly, which requires \( \dot{M}_w(R) \) or equivalently \( \dot{m}_w(R) \) to be specified. To keep things tractable, a simple power law form will be used here,

\[
\dot{m}_w(R) = KR^\xi,
\]

with the constant \( K \) being fixed by normalizing to \( \dot{M}_{w,\text{total}} \) (which is used as a free parameter). Note that, in principle, the local mass-loss rate may itself depend on the local effective temperature; in this case an iterative procedure would be needed to deal with the resulting non-linearity.

Equation (10) can now be integrated straightforwardly. Using the usual boundary condition, \( \frac{\partial \Omega(R_\ast)}{\partial R} = 0 \) (which really applies at or near the boundary layer, rather than at the stellar surface itself; e.g. Frank, King & Raine 1985), and the identification \(-2\pi R\Sigma R\Omega = \dot{M}_{\text{acc}}(R) \) made above for the mass conservation integral, this yields

\[
3\pi\nu \Sigma = \dot{M}_{\text{acc}}(R) - \dot{M}_{\text{acc}}(R_\ast) \left( \frac{R_\ast}{R} \right)^{1/2} + \dot{M}_{w,\text{total}} \left( \frac{R_\ast}{R} \right)^{1/2} \left[ \frac{\pi^2 (\xi + 2/3)}{(\xi + 5/2)} \right] \left[ \frac{1 - \left( \frac{R_\ast}{R} \right)^{\xi/2}}{\left( \frac{R_\ast}{R} \right)^{\xi/2} - 1} \right] \xi \neq -2, -5/2
\]

\[
D(R) = \frac{1}{2} \nu \Sigma \left( R \frac{\partial \Omega}{\partial R} \right)^2 = 3\pi\nu \Sigma \left( \frac{3GM_\ast}{8\pi R^3} \right)
\]

(e.g. Frank, King & Raine 1985). Together with Equation (11), this yields

\[
D(R) = \left\{ \begin{array}{ll}
\left( \frac{3GM_\ast}{8\pi R^3} \right) \left\{ \dot{M}_{\text{acc}}(R) - \dot{M}_{\text{acc}}(R_\ast) \left( \frac{R_\ast}{R} \right)^{1/2} + \dot{M}_{w,\text{total}} \left( \frac{R_\ast}{R} \right)^{1/2} \left[ \frac{\pi^2 (\xi + 2/3)}{(\xi + 5/2)} \right] \left[ \frac{1 - \left( \frac{R_\ast}{R} \right)^{\xi/2}}{\left( \frac{R_\ast}{R} \right)^{\xi/2} - 1} \right] \xi \neq -2, -5/2 \\
\left( \frac{3GM_\ast}{8\pi R^3} \right) \left\{ \dot{M}_{\text{acc}}(R) - \dot{M}_{\text{acc}}(R_\ast) \left( \frac{R_\ast}{R} \right)^{1/2} + \dot{M}_{w,\text{total}} \left( \frac{R_\ast}{R} \right)^{1/2} \left[ \frac{\pi^2 (\xi + 2/3)}{(\xi + 5/2)} \right] \left[ \frac{1 - \left( \frac{R_\ast}{R} \right)^{\xi/2}}{\left( \frac{R_\ast}{R} \right)^{\xi/2} - 1} \right] \xi \neq -2, -5/2 \\
\left( \frac{3GM_\ast}{8\pi R^3} \right) \left\{ \dot{M}_{\text{acc}}(R) - \dot{M}_{\text{acc}}(R_\ast) \left( \frac{R_\ast}{R} \right)^{1/2} + \dot{M}_{w,\text{total}} \left( \frac{R_\ast}{R} \right)^{1/2} \left[ \frac{\pi^2 (\xi + 2/3)}{(\xi + 5/2)} \right] \left[ \frac{1 - \left( \frac{R_\ast}{R} \right)^{\xi/2}}{\left( \frac{R_\ast}{R} \right)^{\xi/2} - 1} \right] \xi \neq -2, -5/2 \\
\left( \frac{3GM_\ast}{8\pi R^3} \right) \left\{ \dot{M}_{\text{acc}}(R) - \dot{M}_{\text{acc}}(R_\ast) \left( \frac{R_\ast}{R} \right)^{1/2} + \dot{M}_{w,\text{total}} \left( \frac{R_\ast}{R} \right)^{1/2} \left[ \frac{\pi^2 (\xi + 2/3)}{(\xi + 5/2)} \right] \left[ \frac{1 - \left( \frac{R_\ast}{R} \right)^{\xi/2}}{\left( \frac{R_\ast}{R} \right)^{\xi/2} - 1} \right] \xi \neq -2, -5/2
\end{array} \right.
\]

for a Keplerian disk. This expression correctly reduces to the usual result for the case of no mass loss upon setting \( \dot{M}_{w,\text{total}} = 0 \) and noting that \( \dot{M}_{\text{acc}}(R) = \dot{M}_{\text{acc}}(R_{disk}) = \text{const} \) in a conservative disk. The explicit form of Equation (11) for a conservative disk is

\[
D_{\text{no wind}}(R) = \frac{3GM_\ast \dot{M}_{\text{acc}}}{8\pi R^3} \left[ 1 - \left( \frac{R_\ast}{R} \right)^{1/2} \right].
\]

It is convenient to rewrite Equation (11) in such a way that the difference between it and the conservative case, Equation (12), becomes more explicit. After some algebraic manipulation, one finds that \( D(R) \) can be expressed as

\[
D(R) = D_{\text{no wind}}(R) - \frac{3GM_\ast \dot{M}_{w,\text{total}}}{8\pi R^3} X(R),
\]

\[\xi = 000, R = 000, 000-000\]
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where

\[
X(R) = \begin{cases} 
1 - \left( \frac{R}{R_e} \right)^{1/2} & \text{R > } R_W \\
1 - \left( \frac{R}{R_e} \right)^{1/2} & \text{R < } R_W \\
\end{cases}
\]

For reference, it is noted without explicit derivation that if \( \dot{m}(R) \) is a delta-function centered on \( R_W \), then \( D(R) \) can still be written in the form of Equation [3] but with \( X(R) \) now given by (c.f. Equation 2.6 in Piran 1977; note that a power 1/2 is missing from the second term in the \( R > R_W \) part of his equation)

\[
X(R) = \begin{cases} 
\xi - 2, -5/2 & R > R_W \\
\xi = -2 & R > R_W \\
\xi = -5/2 & \text{R < } R_W \\
\end{cases}
\]

3 THE RADIAL EFFECTIVE TEMPERATURE PROFILE OF A STEADY-STATE, MASS-LOSING ACCRETION DISK

Given the rate of viscous dissipation in the steady-state disk, the radial effective temperature profile \( T(R) \) is usually found by assuming that the rate at which accretion energy is released by dissipation is balanced by radiative losses, i.e.

\[
D(R)_{\text{no wind}} = \sigma T_{\text{eff}}^4(R)
\]

if the disk is optically thick. However, if a wind is driven from the surface of the disk, some of the dissipated accretion energy may not be radiated away to infinity, but may go instead into powering the wind. In particular, two energy sinks need to be considered: (i) the binding energy that must be overcome to allow matter to escape at all; (ii) the kinetic energy carried by the flow at infinity.

Strictly speaking, a self-consistent dynamical model of the full disk + wind system is needed to properly describe the process of converting the energy dissipated in the disk into wind energy. However, in the spirit of the current treatment, whose point is to elucidate the effects of mass loss on the disk’s energy balance more generally, it makes sense to simply treat wind energy losses, \( l_w(R) \), as a local cooling term, i.e.

\[
D(R) = \sigma T_{\text{eff}}^4(R) + l_w(R).
\]

Local approximations such as this are crude, since they do not allow for the possibility that the dissipated accretion energy may be redistributed radially before powering the wind. It is also worth noting that, even though Equation [3] can be made energetically correct by construction (see below), it is quite possible for \( l_w(R) \) to have a strong wavelength dependence. For example, in a line-driven wind, energy losses from the disk’s radiation field to the wind may be concentrated in a few strong spectral lines. In this case, estimates of the effective disk temperature derived from the observed continuum would yield values corresponding to \( l_w(R) \sim 0 \) in Equation [3].

Ignoring these subtleties for the moment, some general statements can be made about the conditions under which a wind loss term is required in the energy balance equation at all. First, the rotational energy that is extracted from the disk by the outflow is insufficient to overcome the binding energy of the outflowing matter unless \( \xi = 3/2 \). This can be shown, for example, by integrating \( D(R) \) over \( R \) for the case where \( \dot{m}(R) \) is a delta function (Equations [3] and [15]) and comparing the result to the case of a conservative disk. This yields a “luminosity depletion” of (c.f. Equations 2.7 and 2.8 in Piran 1977)

\[
\Delta L = \frac{GM \dot{M}_w,\text{total}}{2 R_e} - \left( \frac{3}{2} - \xi \right) \frac{GM \dot{M}_w,\text{total}}{R_W}. \tag{18}
\]

* By not explicitly accounting for the thermal energy content of the wind, we are effectively considering only models in which all of the wind’s thermal energy is eventually converted into bulk kinetic energy (i.e. models in which \( T_{\text{wind}} \to 0 \) as \( r \to \infty \)). Close to the disk, the thermal wind energy must in any case be small compared to the binding energy unless the mass loss is in fact thermally driven. This requires temperatures on the order of \( T_{\text{esc}} \approx 3 \times 10^8 (M_*/M_\odot)/(R/10^9 \text{cm})^{-1} \) K near the base of the flow. Such outflows have been considered by Piran (1977) and by Czerny & King (1989ab).
If the wind material had never been accreted at all, the depletion would have been equal to the first term in Equation 18. Thus, unless \( l^2 \geq 3/2 \), the disk still gains energy from the matter that is ultimately lost to the outflow; for \( l = 1 \), this energy is just the binding energy of the wind material. This result should not come as a surprise: for \( l = 1 \), the outflow clearly exerts no net torque on the accretion disk, and thus no energy is extracted from the disk to actually power the wind. The accretion disk therefore dissipates all of gravitational binding energy it extracted from the wind material as it accreted from infinity to \( R_W \) prior to its ejection. For \( l < 1 \), even more energy can be dissipated by the disk, since the outflowing material carries with it less rotational kinetic energy than it had at the point of ejection. Specifically, for \( l = 0 \), the disk gains both the rotational kinetic energy of the wind material \( (GM_m/2R_W) \) and its potential energy \( (GM_m/R_W) \).

Thus unless an external energy source is invoked to help power the wind, \( l_w \) must, in general, contain at least an effective binding energy term for models with \( l^2 < 3/2 \). After all, the outflowing material cannot reach infinity unless at least its binding energy (including any rotational energy it “injected” into the disk) is returned to it. Equation 18 also shows, however, that models with \( l^2 > 3/2 \) do extract sufficient energy from the disk to overcome the binding energy. In fact, if \( l^2 \gg 3/2 \), the energy that is extracted from the disk is sufficient to power the wind to potentially very high terminal velocities of \( \eta = \sqrt{2(l^2 - 3/2)GM_m/R_W} \approx \sqrt{2GM_m/R_W} \). This illustrates how centrifugally-driven accretion disk winds are ultimately powered by accretion energy and shows that explicit binding and kinetic energy terms need not be added to the right-hand side of Equation 17 when modeling disks that drive outflows of this type.

With these general points in mind, the wind loss term, \( l_w \), can be conveniently parameterized as follows. The rate at which energy must be supplied (per unit area) to wind material in order for it to overcome its binding energy and escape from the disk is

\[
l_w(R) = \frac{3}{2} l_m(R) v_{\infty}^2(R),
\]

where \( v_{\infty}(R) \) denotes the Keplerian velocity. Here, no distinction has been made yet between models with different lever arms, which can have very different effective binding energies; this will be accounted for below by means of an efficiency factor.

Similarly, kinetic energy has to be supplied to the wind at a rate of

\[
l_k(R) = \frac{1}{2} \dot{m}_W(R) v_{\infty}^2(R) = \frac{1}{2} \dot{m}_W(R) f^2 v_{\infty}^2(R)
\]

per unit area, where the natural scaling \( v_{\infty}(R) \propto v_{\text{esc}}(R) \propto v_K(R) \) has been assumed, with \( f \) a constant. Here, \( v_{\text{esc}}(R) \) denotes the local escape velocity at a given streamline foot point.

Depending on the length of the lever arm \( l \) and on whether there is an external energy source that helps to power the wind, both \( l_b \) and \( l_k \) may contribute fully or in part to \( l_w \). Allowing for this, \( l_w \) can be written as

\[
l_w(R) = \eta_b l_b(R) + \eta_k l_k(R) = \frac{1}{2} (\eta_b + \eta_k f^2) \dot{m}_W(R) v_K^2(R).
\]

where \( \eta_b \) and \( \eta_k \) are efficiency factors. If there are no external energy sources,

\[
\eta_b = \begin{cases} 3 - 2l^2 & l^2 < 3/2 \\ 0 & l^2 > 3/2 \end{cases}
\]

and

\[
\eta_k = \begin{cases} 1 & l^2 < 3/2 \\ 1 - \frac{(l^2 - 3/2)v_{\infty}^2}{v_{\text{esc}}^2} & l^2 > 3/2. \end{cases}
\]

Equation 17 with \( l_w(R) \) given by Equation 19 is the sought-after expression for the radial effective temperature distribution across a mass-losing accretion disk. On substituting 11 or 13 into 17, it is easy to solve for \( T_{\text{eff}}(R) \) corresponding to any consistent set of parameters. More explicitly, the result is

\[
\sigma T_{\text{eff}}^4(R) = \frac{3GM_m \dot{M}_{\text{acc}}}{8\pi R^3} \left[ 1 - \left( \frac{R_s}{R} \right)^{1/2} \right] - \frac{3GM_m \dot{M}_{\text{w,total}}}{8\pi R^3} X(R) \left( \frac{1}{2} \eta_b + \eta_k f^2 \right) \dot{m}_W(R) v_K^2(R),
\]

where \( X(R) \) is given by Equation 14. The physical meaning of the terms on the right-hand side of this expression is straightforward: the first term is the dissipation rate in the absence of any outflow; the second term accounts for the reduction in dissipation that results from the loss of mass and angular momentum from the disk to the wind; the third term represents the cooling that results if the energy dissipated by the disk must also supply some or all of the wind’s binding and kinetic energies. Within the second term, in the expression for \( X(R) \), terms that involve (do not involve) the lever arm \( l \) arise because the wind acts as a sink for angular momentum (mass).
4 SUMMARY OF THE MODEL PARAMETERS

A total of six parameters are needed to specify how the presence of the accretion disk wind affects the disk’s radial effective temperature distribution:

(1) The length of the rotational lever arm $l$. This variable permits the most meaningful division of the available parameter space and, in principle, allows three types of accretion disk winds to be identified:

(i) $l < 1$: This corresponds to the family of disk winds that carry away less angular momentum than possessed by the wind material before it left the disk surface. Thus each disk annulus loses a larger fraction of its mass than of its angular momentum to the outflow. Even though such disk winds provide a net angular momentum sink (unless $l = 0$), they actually increase the specific angular momentum of the remaining disk material. This family of outflows is unlikely to be physical and will not be considered further below.

(ii) $l = 1$: This corresponds to the family of specific angular momentum-conserving disk winds. This prescription should be most appropriate for radiation-driven outflows (e.g. Murray & Chiang 1996; Proga, Stone & Drew 1998; Feldmaier & Shlosman 1999; Feldmaier, Shlosman & Vitello 1999).

(iii) $l > 1$: This corresponds to the family of accretion disk winds that remove both specific and net angular momentum from the disk. This prescription is appropriate for centrifugally-driven MHD winds (e.g. Blandford & Payne 1982; Cannizzo & Pudritz 1988; Emmering, Blandford & Shlosman 1992; Pelletier & Pudritz 1992). In these, the length of the lever arm corresponds to the Alfvén radius, $R_A$, out to which bead-on-a-wire type rotation on magnetic field is (effectively) enforced. Thus $l = R_A/R$ in this case. Note that, in general, thermally-driven outflows also belong to this class (Piran 1977).

(2) The total wind mass-loss rate, $M_{w,total}$.

(3) The radial mass loss power law index $\xi$.

(4) The velocity parameter $f = v_\infty/v_K$.

(5) The fraction $\eta_k$ of the wind’s kinetic energy that is provided by dissipated accretion energy. Note that for $l < 1$, $\eta_k$ can be greater than unity; see Section 1.

(6) The fraction $\eta_b$ of the wind’s kinetic energy at infinity that is provided by dissipated accretion energy.

5 CONSISTENCY OF PARAMETER COMBINATIONS

The purely parametric approach taken above to obtain the dissipation rate and effective temperature distributions across a steady-state, mass-losing accretion disk does not ensure that a self-consistent solution exists for any given set of specified parameters. This is easily seen, for example, by noting that nothing prevents us from specifying the mass-loss rate to be greater than unity; see Section 3.

One consistency check that can be performed even within the present parametric approach is to identify energetically inconsistent parameter sets. These reveal themselves by giving rise to formally negative dissipation rates and effective temperatures in some or all disk regions, and, strictly speaking, all such solutions should be rejected outright. However, some parameter sets may only just fail to achieve energetic consistency, for instance as a result of the turnover in the (standard) dissipation rate distribution in the innermost disk regions, which the adopted power law mass loss distribution fails to follow. It is therefore useful to define a quantitative measure of the degree of energetic inconsistency exhibited by a given solution and to consider a somewhat less restrictive rejection criterion based on this.

One such measure is the ratio, $R$, of the “negative luminosity” contained in any disk regions in which the dissipation rate and/or effective temperature is formally negative to the positive luminosity that remains in all other disk regions. The absolute value of this ratio is the fraction of the remaining luminosity that would have to be diverted to the negative temperature regions in order to make the solution consistent there. We may then choose to accept solutions for which $|R| < 1$, provided that $|R| << 1$. The plausibility argument behind this more lenient rejection criterion is that only a minor adjustment of the
6 MASS-LOSING ACCRETION DISKS IN NON-MAGNETIC CATACLYSMIC VARIABLES

In this section, the formalism developed above will be used to consider the effect of mass loss in the form of radiation-driven (Section 6.1) and centrifugally-driven winds (Section 6.2) on the effective temperature distribution of accretion disks in cataclysmic variable stars (CVs). These two types of outflows show the most promise for explaining the observed mass loss in these systems.

Kinematic modeling of wind-formed, ultraviolet line profiles in CVs suggests mass-loss rates on the order of a few percent of the accretion rate (Shlosman, Vitello & Mauche 1997; Keggie & Drew 1997; Prinja & Rosen 1995; Hoare & Drew 1993). Guided by this, a value of $M_{\nu,\text{total}}/M_{\text{acc}}(R_{\text{disk}}) = 0.025$ will be adopted for both dynamical pictures below. Note that the current generation of numerical and analytical models for line-driven winds is still struggling to achieve such high efficiencies (Proga, Drew & Stone 1998; Feldmaier & Shlosman 1999; Feldmaier, Shlosman & Vitello 1999). Centrifugally-driven disk winds do not share this problem, but do require the presence of a sufficiently strong, large-scale, ordered magnetic field threading the disk (Section 3.2; Cannizzo & Pudritz 1988; Pelletier & Pudritz 1992).

For the mass loss power law index $\xi$, values of 0,-1,-2,-5/2 and -3 will be considered for all models. Note that not all of these values are equally plausible in the context of a given dynamical picture. This point will be considered in more detail in the relevant sections below. Appropriate choices for the remaining outflow parameters ($f$, $\eta_k$ and $\eta_w$) are also made and explained in these sections.

Finally, the following system parameters are adopted throughout: $\dot{M}_{\text{acc}}(R_{\text{disk}}) = 10^{-8}M_\odot\text{yr}^{-1}$, $M_* = 0.5M_\odot$ and $R_* = 0.0133R_\odot$. These are appropriate for a non-magnetic CV in a high state, e.g. a nova-like variable like UX UMa.

6.1 Radiation-driven disk winds

As already noted above, radiation-driven disk winds are expected to belong to the $l = 1$ family of models. The terminal velocities in these types of outflows are typically on the order of the (local) escape speeds on the disk (e.g. Proga, Stone & Drew 1998; Feldmaier, Shlosman & Vitello 1999), so it is reasonable to take $f = \sqrt{2}$.

Two extreme possibilities are considered for the efficiency parameters $\eta_k$ and $\eta_w$:

(i) Minimal models [$\eta_k = \eta_w = 0$]: For this choice of parameters, the effect of the outflow on the disk is minimized. The minimal model is appropriate if the wind is powered predominantly by an external energy source, such as the boundary layer between the disk and the central star. However, somewhat counter-intuitively, it may also be relevant if the disk does power the wind. This is because a radiation-driven outflow is quite unlikely to drain energy from the disk’s radiation field in a manner that would observationally mimic a cooler (blackbody) disk. Most importantly, in a line-driven wind, the transfer of momentum and energy from the radiation field to the wind often takes place in a few strong lines. In this case, the overall continuous disk spectrum is hardly affected and the minimal model is likely to provide the most relevant benchmark for comparisons with observations.

(ii) Maximal models [$\eta_k = \eta_w = 1$]: For this choice of parameters, the effect of the outflow on the disk is maximized. Energetically, the maximal model corresponds to the assumption that the disk alone is responsible for powering the outflow. However, as just noted above, the minimal model may nevertheless be preferable to describe the observational effects of a purely disk-driven wind. The maximal model is included mainly to set an upper limit on the effect of a radiation-driven wind on the disk’s temperature distribution.

The radial effective temperature distributions corresponding to the minimal and maximal models are shown in Figures 1 and 2. In all of these models, the presence of the outflow results in reduced temperatures over the full range of radii, which simply reflects the removal of mass and angular momentum: the mass-losing disk needs to exert less viscous torque than a conservative to transport material towards the accretor. As a result, the viscous dissipation rates and effective temperatures are also lower.

In most cases, the differences between $T_{\text{eff}}(R)$ with and without mass loss are small everywhere, with $-\Delta T = - (T_{\text{eff}} - T_{\text{eff, no wind}}) \lesssim 10^3$ K. Such minor deviations are not observable with current instruments and techniques.

It is worth noting that, even in a hypothetical “grey” wind (i.e. a radiation-driven outflow with frequency-independent opacity), energy is extracted from the radiation field by red-shifting the entire incident spectrum. This is not, in general, equivalent to (or closely approximated by) replacing the incident spectrum by one that corresponds to a lower temperature. An excellent discussion of momentum and energy transfer in radiation-driven stellar winds is given by Gayley, Owocki & Cranmer (1995).
The only exceptions are the $\eta_b = \eta_k = 1$, $\xi \leq -2$ models. For these parameter combinations, the effective temperature distribution drops precipitously just beyond the stellar surface. The reason for this behaviour is easy to understand. In these models, much or most of the mass lost to the wind is carried away from the inner disk, where the Keplerian and escape velocities are highest. Thus the wind binding and kinetic luminosities are largest in these cases, and, for $\eta_b = \eta_k = 1$, must be supplied entirely by the disk. As a result, most of the accretion energy that is liberated in the innermost disk regions is needed to power the outflow.

How reasonable are these $\eta_b = \eta_k = 1$, $\xi \leq -2$ parameter combinations? Intuitively, it seems inevitable that, in a radiation-driven accretion disk wind, mass loss should occur preferentially from the hot, luminous, inner disk regions. Dynamical models appear to confirm this basic expectation (Proga, Stone & Drew 1998; Feldmaier, Shlosman & Vitello 1999). It is particularly worth noting that, in a conservative disk (and, to an equally good approximation, also in the mass-losing disks considered here), $T_{eff} \propto R^{-3/4}$ for $R \gg R_*$. This implies that $\xi = -3$ corresponds to the plausible situation in which the mass loss rate per unit area is proportional to the rate at which energy is dissipated and radiated locally. In fact, the analytic model of Feldmaier, Shlosman & Vitello (1999) predicts $\dot{m}_w(R) \propto R^{-2.9}$ for $R \gtrsim 5 R_{WD}$ (and $\dot{m}_w(R) \propto R^{-1}$ closer to disk centre). Thus $\xi \leq -2$ is a reasonable choice for a line-driven disk wind.

However, as noted in Section 3 the $\eta_b = \eta_k = 1$ maximal model is likely to overestimate the effects of radiation-driven mass loss on the disk’s continuum spectrum and on the effective temperature distribution that would be inferred for such a disk from observations. Moreover, since deviations from the standard case are so highly localized in the innermost disk regions in the $\xi = -3$ case, it is highly unlikely that even this model would be observationally distinguishable from a conservative disk. Even the most promising technique to infer $T(R)$ from observations, which relies on the inversion of eclipse light curves of high-inclination systems (Horne 1985), cannot currently provide the high spatial resolution that would be required to make such a distinction ($\Delta R \ll 1 R_*$). A final problem is that an optically thick, hot boundary layer between the inner disk and the stellar surface could effectively mask this extremely narrow break in the disk’s temperature distribution close to the WD (Popham & Narayan 1995).

Thus radiation-driven accretion disk winds are unlikely to leave an observable imprint on the disk’s effective temperature distribution.

### 6.2 Centrifugally-driven disk winds

The most important parameter that must be fixed to describe a centrifugally-driven outflow within the present formalism is the length of the wind’s effective lever arm $l$. As noted in Section 4 in this type of outflow $l R$ corresponds to the Alfvén radius, $R_A$.

In the absence of external energy sources, a firm lower bound on $l$ is obtained by noting that $l^2 > 3/2$ is required for such a wind to even overcome its initial binding energy (Section 4). A rough upper bound follows from the fact that a centrifugally-driven wind alone – i.e. without any help from viscous torques – can, in principle, drive mass accretion at a rate of about (e.g. Cannizzo & Pudritz 1988; Pelletier & Pudritz 1992)

$$\frac{\dot{M}_{\text{acc}}}{\dot{M}_{\text{wind}}} \simeq \left( \frac{R_A}{R} \right)^2 = l^2. \tag{25}$$

For $\dot{M}_{\text{w,total}}/\dot{M}_{\text{acc}}(R_{\text{disk}}) = 0.025$, this corresponds to $l = 6.3$. This is an upper limit, because observations of disk radius changes in erupting dwarf nova systems indicate that viscous angular momentum transport must be significant and is probably dominant in CVs (e.g. Livio 1999). One arrives at the same conclusion by noting that a purely wind-driven disk, in which no angular momentum is transported by viscosity and instead all is extracted by an outflow, would actually be dark. This was most recently pointed out by Spruit (1996) and also follows trivially from the fact that there can be no viscous dissipation without viscosity: with $\nu = 0$ in Equation (10), $D(R)$ and $T(R)$ must also vanish. This upper bound on $l$ is only approximate, mainly because boundary effects have been ignored in the derivation of Equation 22. With both of these limits in mind, $l = 5$ is adopted in this section to illustrate the effect of a centrifugally-driven wind on the disk.

Before presenting the results for models of this type, a few preliminary remarks are in order. First, it is of interest to consider the strength of the large-scale, poloidal, magnetic field that is required by centrifugally-driven outflows. Following Cannizzo & Pudritz (1988), one estimate can be obtained by assuming equipartition between the magnetic and gas pressures in the disk photosphere, i.e.

$$\frac{B_{eq}^2}{8\pi} = \rho c_s^2 \tag{26}$$

where $\rho$ is the mass density and $c_s$ is the sound speed. Using the standard $\alpha$-disk solutions (e.g. Frank, King & Raine 1992), one finds equipartition field strengths of $B_{eq} \approx 2000(R/10 R_{WD})^{21/16} G$ for the system parameters adopted here. This estimate is probably close to an upper limit; even though equipartition of gas and magnetic pressures may be a good approximation in the disk atmosphere, it is not necessary for all of the magnetic flux to be in the form of the large-scale field that is needed to launch the wind. An estimate of the minimum strength required for this field (denoted $B_{eq}$) is roughly given by
\[ \frac{B_p^2}{B_{eq}^2} \sim \frac{H}{M_{\text{acc}} R} \]  

where \( H \) is the disk scale-height and \( \alpha \) is the usual Shakura-Sunyaev viscosity parameter (c.f. Pringle 1993). This is larger by a factor \( l \) compared to Pringle’s formula because the wind velocity has been taken to be of order \( l \)-times the Keplerian velocity. Equation 27 shows that the large-scale poloidal field can, in principle, be smaller than \( B_{eq} \) by factors of 10-100. \( \xi \)

Second, it should be stressed that the origin of the large-scale magnetic field, as well as the mechanism for maintaining it (with a topology suitable for driving an outflow), have not been – and need not be – specified in the present formalism. However, it is important to keep in mind that these issues are central to the question of whether centrifugally-driven winds can actually exist in real systems (see, for example, Tout & Pringle 1996; Reyes-Ruiz & Stepinski 1996; Lubow, Papaloizou & Pringle 1994ab; van Ballegooijen 1989).

Third, the ratio \( R_A/R = l \) might, in general, itself be a function of radius \( R \) on the disk (e.g. Pelletier & Praditz 1992). In fact, only \( \xi = -2 \) is consistent with \( l = \text{constant} \) in a purely wind-driven, Keplerian disk. This is because angular momentum has to be removed from such a disk in just the right way to allow it to accrete while still remaining Keplerian. In “mixed” disks, such as those considered here, any combination of \( l \) and \( \xi \), if, in principle, viable, so long as it does not imply the removal of more angular momentum than is available. Partly because of this flexibility of mixed disks, the additional complication involved in considering, say, a power law form for \( l \) did not seem worth pursuing here. If necessary, it would be easy to add this, albeit at the cost of at least one additional free parameter. Note that a mixed disk with \( l = \text{constant} \) and \( \xi = -2 \) should be expected to mimic a conservative disk with a lower accretion rate.

Fourth and finally, the efficiency parameters \( \eta_b \) and \( \eta_c \) are both set to zero for the models discussed in this section. This is because a centrifugally-driven wind already extracts enough energy from the disk to overcome its binding energy and attain terminal velocities well in excess of the local escape velocities (\( v_{\text{esc}} \approx l v_{\text{esc}} \); c.f. Section 3).

With these remarks out of the way, the results for this class of models are shown in Figure 3. In all cases, the effective disk temperatures have been affected more drastically than in their \( l = 1 \) counterparts. This was to be expected, since the outflows here carry away more angular momentum and hence remove more accretion luminosity in the form of rotational kinetic energy. A useful way to compare these two types of models is to note that in both the outflow is accelerated to about \( l \) times the local escape velocity. Thus a centrifugally-driven wind extracts roughly \( l^2 \) times more energy from the disk than even a maximal radiation-driven wind (unless the latter also were to attain \( v_{\text{esc}} \gg v_{\text{esc}} \)).

The effects of a centrifugally-driven outflow on \( T_{\text{eff}}(R) \) depend on \( \xi \) in a straightforward manner. Models in which mass loss does not decline too rapidly with radius (\( \xi > -2 \)) slightly steepen the slope of the temperature distribution in the outer disk region. By contrast, models with \( \xi < -2 \) tend to flatten the temperature distribution further in and eventually again produce a sharp break in the effective temperature distribution near disk center (but note that even for the \( \xi = -3 \) solution \( |R| = 0.027 \ll 1 \)). For \( \xi = -2 \) exactly, temperatures are reduced across the disk by an essentially constant factor.

In dynamical models of centrifugally-driven winds, \( \xi \) is related to the radial distribution of the large-scale magnetic field that threads the disk. It is therefore hard to rule out any particular value for this parameter a priori. The model of Blandford & Payne (1982), for example, corresponds to \( \xi = -2 \). A potential, empirical restriction follows from the fact that the terminal velocity of a centrifugally-driven wind is roughly \( v_{\text{esc}} \approx l v_{\text{esc}} \), where \( v_{\text{esc}} \) is the local escape velocity from a streamline footpoint on the disk. If a significant fraction of the wind escapes from the inner disk, as in the \( \xi \lesssim -2 \) models, one might expect to see maximum outflow velocities of about \( l \)-times the escape velocity from the central object. This is somewhat higher than observed in CVs, where the maximum velocities inferred from wind-formed spectral lines are closer to the escape velocity itself (e.g. Prinja & Rosen 1995). However, it would be premature to conclude that such models are therefore ruled out in CVs. This is because projection effects will tend to reduce the observed maximum velocities. In addition, the outflow might be too dilute to produce much line emission or absorption by the time it reaches its terminal speed. Both effects are seen, for example, in the kinematic disk wind model developed by Knigge & Drew (1997) for the nova-like variable UX UMa.

Figure 3 shows that only models with \( \xi < -2 \) are likely to be observationally distinguishable from a conservative disk. (For \( \xi = -2 \) the net change in the effective temperatures may be large enough, but the temperature distribution closely resembles that of a conservative disk with a lower accretion rate.) A CV disk that drives such a centrifugally-driven wind would have a redder overall spectrum and would produce non-standard eclipse light curves if observed in a high-inclination system. It is interesting to note that such signatures may already have been seen in CVs. Spectroscopically, Long et al. (1994) showed that a “missing” or cooler than expected inner disk could remove the discrepancies between Hopkins Ultraviolet Telescope

\[ \frac{B_p^2}{B_{eq}^2} \sim \frac{H}{M_{\text{acc}} R} \]

Note that even if the strength of the large-scale field were close to \( B_{eq} \), its presence would not imply that such a system would be classified as a “magnetic CV”. The latter term is reserved for systems in which the disk is disrupted, or prevented from forming, by the magnetic field of the WD primary. For reference: a WD dipole field with a surface strength of about \( 5 \times 10^6 \) G would be required to produce a field comparable to \( B_{eq} \) at \( 10 R_{WD} \), (taking the same system parameters as above). A system harbouring such a WD would indeed not be a non-magnetic CV. If a disk could form at all, it would be truncated at a radius of about \( 10 R_{WD} \). Even though this is not the type of model considered here, the interaction region between a stellar magnetic field and a truncated accretion disks can itself be the launch pad for a different type of centrifugally-driven outflow, namely Shu et al.’s (1994) X-celerator wind.

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observations of the nova-like variable IX Vel and model accretion disk spectra. Moreover, IX Vel is not unique in that model disk spectra provide a poor match to its ultraviolet spectrum – standard disk models tend to fail in much the same way for other high-state CVs, both nova-likes and dwarf novae in outburst (Wade 1988; Knigge et al. 1997; Knigge et al. 1998). Similarly, eclipse maps derived from observations of some high-state CVs, most notably the so-called SW Sex stars, suggest that the effective (really: brightness) temperature distribution in the inner disk regions of these systems is considerably flatter than expected (Rutten, van Paradijs & Tinbergen 1992).

Unfortunately, it is far from clear that these observational signatures are really due to the presence of an outflow, and a number of alternative possibilities – e.g. magnetic disk truncation, additional radiation sources in the system or unaccounted-for radiative transfer effects – are discussed in the articles cited above. Nevertheless, given that CVs (and other disk accretors) are known to drive accretion disk winds, attempts to explain discrepancies between standard disk models and observations in terms of the effects of these winds certainly have some appeal.

7 CONCLUSIONS

The radial distributions of dissipation rate and effective temperature across a Keplerian, steady-state, mass-losing accretion disk have been derived and examined. The parametric approach used to achieve this is sufficiently general to be applicable to many types of dynamical disk wind models, including both radiation-driven and centrifugally-driven outflows.

To simplify the calculations, the local wind mass-loss rate was assumed to follow a power law distribution in radius. In addition, wind material moving along a streamline originating at radius $R$ in the disk was assumed to effectively co-rotate with the disk out to a radial distance $lR$, so that non-rotating ($l = 0$), specific angular momentum-conserving ($l = 1$) and angular momentum-extracting ($l > 1$) disk winds could all be described within the same framework. To account for cases where the wind is partially or entirely powered by energy dissipated within disk, fixed (but arbitrary) fractions of the wind’s binding and kinetic luminosities were used as cooling terms in deriving the disks effective temperature distribution. In the absence of external (non-disk) energy sources, the energetics of the disk-wind system constrain these fractions to be simple functions of the length of the lever arm $l$. With these assumptions and simplifications, the mass-losing disk’s effective temperature profile could be obtained analytically.

Effective temperature distributions for CV accretion disks that lose mass to radiation-driven or centrifugally-driven winds have been calculated and discussed in detail. For realistic wind mass-loss rates of a few percent centrifugally-driven outflows – particularly those in which mass loss is concentrated near disk centre – appear to significantly alter the disk’s effective temperature distribution. An accretion disk that is affected by the presence of such an outflow could produce spectra and eclipse light curves that are noticeably different from those produced by a standard, conservative disk. In fact, existing discrepancies between disk models and observations of some systems could probably be reconciled within this framework. However, since other mechanisms might also still be able to account for these discrepancies, the latter cannot yet be regarded as clear signatures of a disk wind. This uniqueness problem will disappear once a self-consistent and reliable dynamical picture of accretion disk winds in CVs and other disk accreting systems has been found. This will allow the free parameters of the parametric disk wind model developed here to be fixed. The formalism outlined above can then be used to obtain a simple, but reliable estimate of the impact of the wind on the disk’s effective temperature distribution.

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Figure 1. Top panel: Radial effective temperature distributions of disks losing mass to a radiation-driven wind. The outflow is assumed to be specific angular momentum-conserving ($l = 1$), with a mass-loss rate of $M_{w,\text{total}}/M_{\text{acc}}(R_{\text{disk}}) = 0.025$. None of the wind’s binding or mechanical luminosity is provided by dissipated accretion energy in these “minimal” models ($\eta_b = \eta_k = 0$). Each line corresponds to a different distribution of the local mass-loss rate per unit area with radius, $\dot{m}_{\text{wind}}(R) \propto R^{-\xi}$, as indicated. Bottom panel: The effective temperature differences $-\Delta T = -\left[T_{\text{eff}} - T_{\text{eff, no wind}}\right]$ between the mass-losing disk models and a conservative disk.
Figure 2. Top panel: Radial effective temperature distributions of disks losing mass to a radiation-driven wind. The outflow is assumed to be specific angular momentum-conserving ($l = 1$), with a mass-loss rate of $\dot{M}_{w, \text{total}}/\dot{M}_{\text{acc}}(R_{\text{disk}}) = 0.025$. All of the wind’s binding or mechanical luminosity is provided by dissipated accretion energy in these “maximal” models ($\eta_b = \eta_k = 1$). Each line corresponds to a different distribution of the local mass-loss rate per unit area with radius, $\dot{m}_{\text{wind}}(R) \propto R^{-\xi}$, as indicated. Bottom panel: The effective temperature differences $-\Delta T = -[T_{\text{eff}} - T_{\text{eff, no wind}}]$ between the mass-losing disk models and a conservative disk.
Figure 3. Top panel: Radial effective temperature distributions of disks losing mass to a centrifugally-driven wind. The outflow extracts angular momentum from the disk ($l = 5$) and has a mass-loss rate of $\dot{M}_{\text{w,total}}/\dot{M}_{\text{acc}}(R_{\text{disk}}) = 0.025$. None of the wind’s binding or mechanical luminosity needs to be provided by dissipated accretion energy in these models ($\eta_b = \eta_k = 0$), since the outflow already extracts sufficient energy from the disk to accelerate the ejected material to velocities in excess of the local escape velocity. Each line corresponds to a different distribution of the local mass-loss rate per unit area with radius, $\dot{m}_{\text{wind}}(R) \propto R^{-\xi}$, as indicated. Bottom panel: The effective temperature differences $-\Delta T = -\left[ T_{\text{eff}} - T_{\text{eff, no wind}} \right]$ between the mass-losing disk models and a conservative disk.
Disks with radiation-driven winds (minimal models)

\[ \dot{M}_{\text{wind}} = 0.025 \dot{M}_{\text{acc}} \]

\[ l = 1 \]

\[ \eta_b = \eta_k = 0 \]
Disks with radiation-driven winds (maximal models)

\[ \dot{M}_{\text{wind}} = 0.025 \dot{M}_{\text{acc}} \]

\[ \eta_b = \eta_k = 1 \]
Disks with centrifugally-driven winds

\[ \dot{M}_{\text{wind}} = 0.025 \dot{M}_{\text{acc}} \]

\[ l = 5 \]

\[ \eta_b = \eta_k = 0 \]