Radiative Flow in a Luminous Disk

Jin Fukue
Astronomical Institute, Osaka Kyoiku University, Asahigaoka, Kashiwara, Osaka 582-8582
fukue@cc.osaka-kyoiku.ac.jp

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Abstract

Radiatively-driven flow in a luminous disk is examined in the subrelativistic regime of \((v/c)^1\), taking account of radiation transfer. The flow is assumed to be vertical, and the gravity and gas pressure are ignored. When internal heating is dropped, for a given optical depth and radiation pressure at the flow base (disk “inside”), where the flow speed is zero, the flow is analytically solved under the appropriate boundary condition at the flow top (disk “surface”), where the optical depth is zero. The loaded mass and terminal speed of the flow are both determined by the initial conditions; the mass-loss rate increases as the initial radiation pressure increases, while the flow terminal speed increases as the initial radiation pressure and the loaded mass decrease. In particular, when heating is ignored, the radiative flux \(F\) is constant, and the radiation pressure \(P_0\) at the flow base with optical depth \(\tau_0\) is bound in the range of \(2/3 < cP_0/F < 2/3 + \tau_0\). In this case, in the limit of \(cP_0/F = 2/3 + \tau_0\), the loaded mass diverges and the flow terminal speed becomes zero, while, in the limit of \(cP_0/F = 2/3\), the loaded mass becomes zero and the terminal speed approaches \((3/8)c\), which is the terminal speed above the luminous flat disk under an approximation of the order of \((v/c)^1\). We also examine the case where heating exists, and find that the flow properties are qualitatively similar to the case without heating.

Key words: accretion, accretion disks — astrophysical jets — galaxies: active — radiative transfer — relativity — X-rays: stars

1. Introduction

Accretion disks are tremendous energy sources in the active universe: in young stellar objects, in cataclysmic variables, in galactic X-ray sources and microquasars, and in active galaxies and quasars (see Kato et al. 1998 for a review). In particular, in a supercritical accretion disk, the mass-accretion rate highly exceeds the critical rate, the disk local luminosity exceeds the Eddington one, and the mass loss from the disk surface driven by radiation pressure takes place.

Such a radiatively driven outflow from a luminous accretion disk has been extensively studied in the context of models for astrophysical jets by many researchers (Bisnovatyi-Kogan, Blinnikov 1977; Katz 1980; Icke 1980; Melia, Königl 1989; Misra, Melia 1993; Tajima, Fukue 1996, 1998; Watarai, Fukue 1999; Hirai, Fukue 2001; Fukue et al. 2001; Orihara, Fukue 2003), as on-axis jets (Icke 1989; Sikora et al. 1996; Renaud, Henri 1998; Luo, Protheroe 1999; Fukue 2005a), as outflows confined by a gaseous torus (Lynden-Bell 1978; Davidson, McCray 1980; Sikora, Wilson 1981; Fukue 1982), or as jets confined by the outer flow or corona (Sol et al. 1989; Marcowith et al. 1995; Fukue 1999), and as numerical simulations (Eggleton et al. 1985, 1988). In almost all of these studies, however, the disk radiation field was treated as an external field, and the radiation transfer was not solved.

The radiation transfer in the disk, on the other hand, has been investigated in relation to the structure of a static disk atmosphere and the spectral energy density from the disk surface (e.g., Meyer, Meyer-Hofmeister 1982; Cannizzo, Wheeler 1984; Shaviv, Wehrse 1986; Adam et al. 1988; Hubeny 1990; Ross et al. 1992; Artemova et al. 1996; Hubeny, Hubeny 1997, 1998; Hubeny et al. 2000, 2001; Davis et al. 2005; Hui et al. 2005; see also Mineshige, Wood 1990). In these studies, however, the vertical movement and mass loss were not considered. Moreover, the relativistic effect for radiation transfer was not considered in the current studies.

In this paper we first examine the radiatively driven vertical flow — moving photosphere — in a luminous flat disk within the framework of radiation transfer in the subrelativistic regime of \((v/c)^1\), where we incorporate the effect of radiation drag up to the first order of the flow velocity. At the first stage, we ignore the gravity of the central object as well as the gas pressure, although internal heating is considered in a limited way. We then consider analytical solutions for such a radiative flow, and examine the flow properties.

In the next section we describe basic equations in the vertical direction. In section 3 we examine the radiative flow without internal heating, while the radiative flow with internal heating is discussed in section 4. The final section is devoted to concluding remarks.

2. Basic Equations

Let us suppose a luminous flat disk, deep inside which gravitational or nuclear energy is released via viscous heating or other processes. The radiation energy is trans-
ported in the vertical direction, and the disk gas, itself, also moves in the vertical direction due to the action of radiation pressure (i.e., plane-parallel approximation). For simplicity, in the present paper, the radiation field is considered to be sufficiently intense that both the gravitational field of, e.g., the central object and the gas pressure can be ignored: tenuous cold normal plasmas in the super-Eddington disk, cold pair plasmas in the sub-Eddington disk, or dusty plasmas in the sub-Eddington disk. As for the order of the flow velocity, we consider the sub-relativistic regime, where the terms of the first order of \((v/c)\) are retained, in order to take account of radiation drag. Under these assumptions, the radiation hydrodynamic equations for steady vertical \((z)\) flows are described as follows (Kato et al. 1998).

The continuity equation is
\[
\rho v = J \quad (= \text{const.}),
\]
where \(\rho\) is the gas density, \(v\) the vertical velocity, and \(J\) the mass-loss rate per unit area. The equation of motion is
\[
v \frac{dv}{dz} = \frac{\kappa_{abs} + \kappa_{sca}}{c} [F - (E + P)v],
\]
where \(\kappa_{abs}\) and \(\kappa_{sca}\) are the absorption and scattering opacities (gray), \(E\) the radiation energy density, \(F\) the radiative flux, and \(P\) the radiation pressure. In a gas-pressureless approximation, the energy equation is reduced to
\[
0 = q^+ - \rho \left( j - c \kappa_{abs} E + \kappa_{abs} \frac{2 F v}{c} \right),
\]
where \(q^+\) is the heating and \(j\) is the emissivity. For radiation fields, we have
\[
\frac{dF}{dz} = \rho \left( j - c \kappa_{abs} E + (\kappa_{abs} - \kappa_{sca}) \frac{F v}{c} \right),
\]
\[
\frac{dP}{dz} = \frac{\rho v}{c^2} (j - c \kappa_{abs} E) - \frac{\kappa_{abs} + \kappa_{sca}}{c} [F - (E + P)v],
\]
and the closure relation,
\[
P = \frac{1}{3} E + \frac{4}{3} \frac{F v}{c}.
\]

Eliminating \(j\) and \(E\) with the help of equations (3) and (6), and introducing the optical depth by
\[
d\tau = -(\kappa_{abs} + \kappa_{sca}) \rho dz,
\]
we can rearrange the basic equations up to the order of \((v/c)\) as
\[
\frac{c J}{d\tau} = -(F - 4Pv),
\]
\[
\frac{dF}{d\tau} = -q^+ \frac{1}{(\kappa_{abs} + \kappa_{sca}) \rho} + F \frac{v}{c},
\]
\[
\frac{dP}{d\tau} = F - 4Pv,
\]
\[
(\kappa_{abs} + \kappa_{sca}) \frac{J}{d\tau} = -v.
\]

Finally, integrating the sum of equations (8) and (10) gives the momentum conservation in the present approximation,
\[
Jv + P = K \quad (= \text{const.}),
\]
and similarly from equations (8) and (9) we have the energy conservation,
\[
\frac{1}{2} J v^2 + F = -\int \frac{q^+}{(\kappa_{abs} + \kappa_{sca}) \rho} d\tau,
\]
where the first term on the left-hand side is eventually dropped, although we retain it here to clarify the physical meanings.

We solved equations (8), (12), and (13) for appropriate boundary conditions, and we obtained analytic solutions for several special cases.

As for the boundary conditions, we imposed the following cases. At the flow base (disk “inside”) with an arbitrary optical depth \(\tau_0\), the flow velocity is zero and the radiation pressure is \(P_0\), where subscript 0 denotes the values at the flow base. At the flow top (disk “surface”), where the optical depth is zero, the radiation fields should satisfy the values above the luminous infinite disk. Namely, just above the disk with surface intensity \(I_0\), the radiation energy density \(E_s\), the radiative flux \(F_s\), and the radiation pressure \(P_s\) are \(E_s = (2/c) \pi I_0\), \(F_s = \pi I_0\), and \(P_s = (2/3c) \pi I_0\), respectively, where subscript s denotes the values at the flow top. Rigorously speaking, this latter boundary condition at the flow top is for a static photosphere, and should not be applied to a moving photosphere, although we approximately used this condition in the present subrelativistic regime. The exact boundary conditions are derived and discussed in a separate paper for a fully special relativistic case.

As shown explicitly later, in order for the flow to exist, the radiation pressure \(P_0\) at the flow base should be restricted in some ranges. The lower limit is set due to the negative gradient of the radiation pressure toward the flow top. The upper limit is set when the loaded mass diverges. Beyond the upper limit, the hydrostatic balance without mass loss may be established in the vertical direction.

### 3. Radiative Flow Without Heating

We first examine the case, where there is no viscous or nuclear heating: \(q^+ = 0\). In the luminous accretion disk, it is usually supposed that the heating processes take place deep inside the disk. Hence, the assumption without heating approximately describes the radiative flow in or near the surface envelope of the disk.

In this case and under the present subrelativistic regime, equation (13) means that the radiative flux \(F\) is conserved along the flow,
\[
F = F_s.
\]

Using the boundary conditions at the flow base \((v = 0)\), equation (12) is expressed as
\[
Jv + P = P_0.
\]
Hence, equation (8) becomes
\[ cJ \frac{dv}{dt} = -(F_s - 4P_0v). \] (16)

This equation can be analytically solved to yield the solution for \( v \),
\[ v = \frac{F_s}{4P_0} \left[ 1 - e^{\frac{-4P_0}{cJ}(\tau - \tau_0)} \right]. \] (17)

Thus, the radiative flow from the luminous disk without heating is expressed in terms of the boundary values and the mass-loss rate. In addition, the flow velocity \( v_s \) at the flow top (\( \tau = 0 \)) is
\[ v_s = \frac{F_s}{4P_0} \left( 1 - e^{\frac{-4P_0}{cJ}\tau_0} \right). \] (18)

Using the boundary condition at the flow top, we further impose a condition on the values at boundaries. As already mentioned, at the flow top \( cP = cP_s = (2/3)F_s \). Hence, inserting boundary values into momentum equation (15), we have the following relation:
\[ J \frac{F_s}{4P_0} \left( 1 - e^{\frac{-4P_0}{cJ}\tau_0} \right) + \frac{2}{3c}F_s = P_0. \] (19)

That is, for given \( \tau_0 \) and \( P_0 \) at the flow base, the mass-loss rate \( J \) is determined in units of \( F_s/c^2 \), as an eigen value.

In figure 1 we show the mass-loss rate \( J \) (solid curves) and the terminal velocity \( v_s \) (dashed curves) at the flow top as a function of \( P_0 \) for several values of \( \tau_0 \) at the flow base. The quantities are normalized in units of \( c \) and \( F_s \).

As can be seen in figure 1, the mass-loss rate increases as the initial radiation pressure increases, while the flow terminal speed increases as the initial radiation pressure and the loaded mass decrease.

Moreover, as can be seen in figure 1, and easily shown from equation (19), in order for the flow to exist, the radiation pressure \( P_0 \) at the flow base is restricted in some range,
\[ \frac{2}{3} \leq \frac{cP_0}{F_s} < \frac{2}{3} + \tau_0. \] (20)

At the upper limit of \( cP_0/F_s = 2/3 + \tau_0 \), the loaded mass diverges and the flow terminal speed becomes zero. On the other hand, at the lower limit of \( cP_0/F_s = 2/3 \), the pressure gradient vanishes, the loaded mass becomes zero and the terminal speed approaches \((3/8)c\). This is just the terminal speed of the gas, driven by the radiation pressure under the radiation drag, above the luminous flat infinite disk in the present approximation of the order of \((v/c)^1\) (cf. Icke 1989).

In figure 2 we show the flow velocity \( v \) (solid curves) and the radiation pressure \( P \) (dashed curves) as a function of the optical depth \( \tau \) for several values of \( P_0 \) at the flow base in a few cases of \( \tau_0 \). The quantities are normalized in units of \( c \) and \( F_s \).

When the initial radiation pressure \( P_0 \) at the flow base is large, the pressure gradient between the flow base and the flow top is also large. As a result, the loaded mass \( J \) also becomes large, but the flow final speed \( v_s \) is small due to momentum conservation (12). When the initial radiation pressure \( P_0 \ [> (2/3c)F_s \] is small, on the other hand, the pressure gradient becomes small, and the loaded mass is also small, but the flow final speed becomes large. In the limiting case of \( P_0 = (2/3c)F_s \), the loaded mass \( J \) becomes zero, but the terminal speed \( v_s \) is \((3/8)c\), as already mentioned.

Finally, inserting equation(17) into equation (11), and integrating the resultant equation, we obtain the height \( z \) as a function of the optical depth \( \tau \):
\[ \frac{(\kappa_{abs} + \kappa_{sca}) J}{c} = \frac{F_s}{4cP_0} \left\{ \left( \tau_0 - \tau \right) - \frac{4}{cP_0} \left[ 1 - e^{\frac{-4P_0}{cJ}(\tau - \tau_0)} \right] \right\}, \] (21)
where we use the boundary condition at the flow base: \( z = 0 \) at \( \tau = \tau_0 \). In addition, the height \( z_0 \) of the flow top (disk “surface”) is also expressed as
\[ \frac{(\kappa_{abs} + \kappa_{sca}) J}{c} z_0 = \frac{F_s}{4cP_0} \left( \tau_0 + \frac{2}{3} - \frac{cP_0}{F_s} \right), \] (22)
for the integration to be performed analytically, several functional forms need to be assumed. One simple case, which we assume in the present paper, is a following form:

\[
\frac{q^+}{(κ_{abs} + κ_{sca})} = q_0 (= \text{const.}).
\] (24)

This is quite reasonable since the heating \(q^+\) may be proportional to the density \(ρ\).

In this case, under the present subrelativistic regime and using the boundary conditions at the flow base, equations (8), (13), and (12) become, respectively,

\[
cJ \frac{dv}{dτ} = -(F - 4P_0 v) ,
\] (25)

\[
F = q_0 (τ_0 - τ) = \frac{F_s}{τ_0} (τ_0 - τ) ,
\] (26)

\[
P = P_0 - J v ,
\] (27)

where \(F_s (= q_0 τ_0)\) is again the radiative flux at the flow top (disk “surface”). Here, we impose the boundary condition on \(F\) at the flow base as \(F = 0\) at \(τ = τ_0\).

The above equation (25) can be analytically solved to yield the solution for \(v\):

\[
v = \frac{F_s}{4P_0 τ_0} (τ_0 - τ) - \frac{cJF_s}{(4P_0)^2 τ_0} \left[ 1 - e^{-\frac{4P_0}{cJ}(τ - τ_0)} \right] .
\] (28)

Thus, the radiative flow from the luminous disk with simple heating is also expressed in terms of the boundary values and the mass-loss rate. In addition, the flow velocity \(v_s\) at the flow top \((τ = 0)\) is

\[
v_s = \frac{F_s}{4P_0} - \frac{cJF_s}{(4P_0)^2 τ_0} \left( 1 - e^{-\frac{4P_0}{cJ} τ_0} \right) .
\] (29)

Using the boundary condition at the flow top, we further impose a condition on the values at the boundaries. In this case with simple heating, we have the following relation:

\[
J = \frac{F_s}{4P_0} - \frac{cJF_s}{(4P_0)^2 τ_0} \left( 1 - e^{-\frac{4P_0}{cJ} τ_0} \right) + \frac{2}{3c} F_s = P_0 .
\] (30)

That is, for given \(τ_0\) and \(P_0\) at the flow base, the mass-loss rate \(J\) is determined in units of \(F_s/c^2\).

In figure 3 we show the mass-loss rate \(J\) (solid curves) and the terminal velocity \(v_s\) (dashed curves) at the flow top as a function of \(P_0\) for several values of \(τ_0\) at the flow base in the case with simple heating. The quantities are normalized in units of \(c\) and \(F_s\).

The properties are qualitatively same as those without heating. The mass-loss rate increases as the initial radiation pressure increases, while the flow terminal speed increases as the initial radiation pressure and the loaded mass decreases. Moreover, in order for flow to exist, the radiation pressure \(P_0\) at the flow base is restricted in some range:

\[
\frac{2}{3} < \frac{cP_0}{F_s} < \frac{2}{3} + \frac{1}{2} τ_0 .
\] (31)
Fig. 3. Mass-loss rate $J$ (solid curves) and terminal velocity $v_s$ (dashed curves) in the case with simple heating: (a) $\tau_0 = 1$ and (b) $\tau_0 = 10$. The quantities are normalized in units of $c$ and $F_s$.

At the upper limit, the loaded mass diverges and the flow terminal speed becomes zero. On the other hand, at the lower limit of $cP_0/F_s = 2/3$, the loaded mass becomes zero and the terminal speed approaches $(3/8)c$.

In figure 4 we show the flow velocity $v$ (solid curves) and the radiation pressure $P$ (dashed curves) as a function of the optical depth $\tau$ for several values of $P_0$ at the flow base in a few cases of $\tau_0$ in the case with simple heating. The quantities are normalized in units of $c$ and $F_s$.

In this case with simple heating, the velocity distribution is somewhat different from that in the case without heating. In the case without heating, the flow is rather quickly accelerated to be a terminal speed, since the radiative flux is constant. In the case with heating, on the other hand, the flow is gradually accelerated, since the radiative flux is increasing from the flow base toward the flow top due to heating.

Finally, inserting equation (17) into equation (11), and integrating the resultant equation, we obtain the height $z$ as a function of the optical depth $\tau$:

$$\frac{(\kappa_{\text{abs}} + \kappa_{\text{sca}})J}{c}z = \frac{F_s}{4cP_0} \left\{ \frac{(\tau - \tau_0)^2}{2\tau_0} + \frac{cJ}{4P_0\tau_0}(\tau_0 - \tau) + \left(\frac{cJ}{4P_0}\right)^2 \frac{1}{\tau_0} \left[ 1 - e^{-\frac{4P_0\tau_0}{cJ}(\tau - \tau_0)} \right] \right\},$$

(32)

where we use the boundary condition at the flow base: $z = 0$ at $\tau = \tau_0$. In addition, the height $z_s$ of the flow top (disk “surface”) is also expressed as

$$\frac{(\kappa_{\text{abs}} + \kappa_{\text{sca}})J}{c}z_s = \frac{F_s}{4cP_0} \left\{ \frac{\tau_0}{2} + \frac{2}{3} - \frac{cP_0}{F_s} \right\},$$

(33)

where we use equation (29). Hence, at the upper limit of $cP_0/F_s = 2/3 + \tau_0/2$, where the loaded mass diverges, the height of the flow/disk becomes zero. On the other hand, at the lower limit of $cP_0/F_s = 2/3$, where the loaded mass becomes zero, the height of the flow/disk becomes infinite.

It should be noted that, at the upper limit of $cP_0/F_s = 2/3 + \tau_0/2$ with small terminal speed,$$

T_c^4 = \frac{3}{4} \frac{T_{\text{eff}}^4}{\tau} \left(\frac{2}{3} + \frac{1}{2}\tau_0\right),$$

(34)
where \( T_c \) is the temperature at the flow base (i.e., \( P_0 = \alpha T_c^4 \)), and \( T_{\text{eff}} \) the effective temperature at the flow top (i.e., \( F_\nu = \sigma T_{\text{eff}}^4 \)). This is just the modified Milne approximation.

5. Concluding Remarks

In this paper we have examined the radiative flow from a luminous disk, while taking account of the radiative transfer, in the subrelativistic regime of \((v/c)^4\). The flow velocity, the radiation pressure distribution, and other quantities are analytically solved as a function of the optical depth for the cases without and with heating. At the flow base (disk “inside”), where the flow speed is zero, the initial optical depth is \( \tau_0 \), and the initial radiation pressure is \( P_0 \); in the usual accretion disk these quantities are determined in terms of the central mass, the mass-accretion rate, and the viscous process as a function of the radius. That is, the optical depth is determined by the disk surface density, the radiative flux is given by the central mass and the accretion rate, and the radiation pressure relates to the disk internal structure. At the flow top (disk “surface”), where the optical depth is zero, the radiation fields are assumed to be those above a static flat disk with uniform intensity. In order to match this boundary condition, one relation is imposed on the boundary quantities, and the mass-loss rate \( J \) and terminal speed \( v_\infty \) are both determined by the initial conditions \( \tau_0 \) and \( P_0 \), as eigenvalues. In particular, in order for the flow to exist, the initial radiation pressure is bound in some range: at the upper limit, the loaded mass diverges and the flow terminal speed becomes zero, while, at the lower limit, the loaded mass becomes zero and the terminal speed approaches \((3/8)c\), which is the terminal speed above the luminous disk under the approximation of the order of \((v/c)^4\).

Among various types of astrophysical jets and outflows, some are believed to be accelerated by continuum radiation, and to which the present mechanism can be applied. For example, in supersoft X-ray sources, including accreting white dwarfs, high-velocity mass outflows were reported (e.g., Cowley et al. 1998). The wind velocity measured in these objects is about 3000–5000 \( \text{km s}^{-1} \), which is of the order of the escape velocity of white dwarfs. In addition, the line profile depends on the inclination angle, and the P Cyg profiles are seen. Hence, in supersoft X-ray sources, the outflows should blow off from the innermost accretion disk. However, the gas temperature of the inner disk is so high that the line-driven mechanism cannot operate, and therefore, mass outflows in these objects would be driven by continuum radiation.

In SS 433, a prototype of astrophysical jets, the mass-accretion rate highly exceeds the Eddington one, and twin jets with 0.26 \( c \) would be accelerated by radiation pressure (see, e.g., Cherepashchuk et al. 2005 for a recent observation). The luminosity of microquasar GRS 1915+105 exceeds the Eddington one \((L_{\text{Edd}})\) during its outburst. In quiescent phase, its luminosity never drop below \( \sim 0.3L_{\text{Edd}} \) (Done, Gierliński 2004).

In broad absorption-line quasars (BAL QSOs), strong broad absorption lines in the UV resonance are always blueshifted relative to the emission-line rest frame, which indicates the presence of outflows from the nucleus, with velocities as large as 0.2 \( c \). These moderately high-velocity outflows are supposed to be accelerated by line or continuum radiation (e.g., Murray et al. 1995; Proga et al. 2000; Proga, Kallman 2004). In these objects, the outflows must consist of dense clouds, or be shielded from the central strong source, in order to operate a line-driven mechanism. Otherwise, the material is highly ionized by the central source, and, instead of line-acceleration, continuum acceleration may work (see, e.g., Chelouche, Netzer 2003).

Finally, narrow-line Seyfert 1 galaxies and luminous quasars are also believed to harbor supercritically accreting black holes, and high-velocity outflows are reported as warm absorbers (cf. Blustin et al. 2005). In these objects, radiative acceleration by continuum may play a dominant role.

The radiative flow investigated in the present paper must be quite fundamental problems for accretion disk physics and astrophysical jet formation, although the present paper is only the first step, and there are many simplifications at the present stage.

For example, we have ignored the gravitational field produced by the central object. This means that the flow considered in the present paper would correspond to normal plasmas in the super-Eddington disk, pair plasmas in the sub-Eddington disk, or dusty plasmas in the luminous disk. In other cases, or even for normal plasmas in the super-Eddington disk, the gravitational field would affect the flow properties. In particular, the gravitational field in the vertical direction is somewhat complicated (e.g., Fukue 2002, 2004), and the influence of the gravitational field is important. The effect of gravity will be considered in a separate paper (Fukue 2005b).

We also ignored the gas pressure. This means that the radiation field is sufficiently intense. In general cases, where the gas pressure is considered, there usually appear sonic points (e.g., Fukue 2002, 2004), and the flow is accelerated from subsonic to supersonic. The cross section of the flow also has a similar influence. In this paper we consider a purely vertical flow, and the cross section of the flow is constant. If the cross section of the flow increases along the flow, the flow properties, such as the transonic nature, would be influenced.

In the present paper, we have considered the internal heating processes in a simple form. Since the heating processes should couple with the viscous process or other heating processes, such as a nuclear process, we should treat internal heating more carefully.

In order to take account of the effect of radiation drag, we consider the subrelativistic regime, where the terms of the first order of \((v/c)\) are retained. As a result, the terminal speed \((3/8)c\) for the subrelativistic regime appears in the extreme case. This is quite consistent with the previous knowledge. As is well known, the terminal speed above the luminous flat disk in the full relativity is \((4 - \sqrt{7})/3c \sim 0.45c\) (Icke 1989). To consider highly
relativistic cases, all of the terms up to \( (v/c)^2 \) should be retained (cf. Kato et al. 1998). A related problem is the boundary condition at the flow top. In the present paper, we impose the boundary condition such that the radiation fields at the flow top are those above the luminous flat disk. Rigorously speaking, this boundary condition should be modified, since the disk gas, itself, is now moving. In the highly relativistic case, the deviation would be serious. The exact boundary conditions are derived and discussed in a separate paper for a fully special relativistic case (Fukue 2005c).

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