Radiatively driven electron-positron jets from two component accretion flows.

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Accepted . Received ; in original form

ABSTRACT

Matter accreting onto black holes has long been known to have standing or oscillating shock waves. The post-shock matter puffs up in the form of a torus, which intercepts soft photons from the outer Keplerian disc and inverse Comptonizes to produce hard photons. The post-shock region also produces jets. We study the interaction of both hard photons and soft photons, with on-axis electron-positron jets. We show that the radiation from post-shock torus accelerates the flow to relativistic velocities, while that from the Keplerian disc has marginal effect. We also show that, the velocity at infinity or terminal velocity $\vartheta$, depends on the shock location in the disc.

Key words: Accretion, accretion discs - black hole physics - radiation mechanism: general - radiative transfer - ISM: jets and outflows

1 INTRODUCTION

Jets around quasars and micro-quasars show relativistic terminal speeds. While jets are quite ubiquitous and are associated with a wide range of celestial objects, only some jets around quasars and micro-quasars show highly relativistic terminal speed (\textit{e.g.}, GRS 1915+105, Mirabel & Rodriguez 1994; 3C 273, 3C 345, Zensus \textit{et al.} 1995; M87, Biretta 1993). These relativistic jets are generally associated with compact objects and circumstantial evidences
show that many of these central gravitating objects are black holes. Black holes do not have ‘hard surfaces’ nor do they have atmospheres, hence if observations show that many of these jets come from the vicinity of the black hole, then they must originate from the accretion discs around these black holes.

Inner boundary conditions of matter accreting onto a black hole are (i) supersonic and (ii) sub-Keplerian. Liang and Thompson (1980) showed that sub-Keplerian matter accreting onto a black hole has at least two X-type critical points. In much of the parameter space, it has been shown that the supersonic matter crossing the outer critical point, under goes centrifugal pressure mediated shock (Fukue, 1987; Chakrabarti, 1989), becomes subsonic and enters the black hole through the inner X-type critical point. Entropy is generated at the shock making the post shock region hot. The region in which the flow slows down may be extended if the shock conditions are not satisfied. This hot, slowed down region is puffed up in the form of a torus (hereafter, CENBOL ≡ CENtrifugal pressure supported BOUNDary Layer; see Chakrabarti et al. 1996, hereafter CTKE96). This disc due to its advection term may be called advective accretion disc (Chakrabarti 1989, hereafter C89; Chakrabarti 1990, hereafter C90; Chakrabarti 1996, hereafter C96; CTKE96). Chakrabarti and Titarchuk (1995, hereafter CT95), proposed a disc model which contains both the Keplerian matter and the sub-Keplerian matter. In this model, the Keplerian matter is of higher angular momentum and low specific energy, and settles around the equatorial plane to form the Keplerian disc (see, Shakura & Sunyaev 1973; Novikov & Thorne 1973, hereafter NT73) while the sub-Keplerian matter with high energy and lower viscosity flanks the cooler Keplerian disc from the top and bottom, sandwiching the Keplerian disc known as the sub-Keplerian halo (see, CT95; Chakrabarti 1997, hereafter C97) in the literature. The sub-Keplerian halo may suffer a standing shock at $x_s$, a few tens of Schwarzschild radii, and it may be sustained there if the post-shock thermal pressure is sufficiently high. The shock compresses the pre-shock flow making it denser and at the same time hotter. In the model solution proposed by Chakrabarti & Titarchuk (CT95), the post-shock region is comprised of a mixture of the Keplerian and sub-Keplerian components. Thus, though the sub-Keplerian halo (pre-shock) is optically thin for the radiations from Keplerian disc, the post-shock torus could be optically thin, intermediate, or even thick depending on the Keplerian and sub-Keplerian rates (see CT95, CTKE 96, C97 for details). This is because (1) the Keplerian radiation falls on it at a glancing angle, thereby increasing the path length and (2) the mixed matter in this region has higher density. CT95 showed soft radiation from the cool Keplerian Disc are
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Figure 1. Cross-sectional view of Two Component Accretion Disc Model. Only the top half is shown.

inverse-Comptonized by the CENBOL to produce the hard radiation. If the sub-Keplerian halo rate ($\dot{M}_h$) is higher, then it supplies more hot electrons to the CENBOL than the soft photons from the Keplerian disc and hence the soft photons cannot cool down the CENBOL significantly. Thus CENBOL remains puffed up and hot and can intercept a large number of soft photons and inverse-Comptonize them to produce the hard power-law tail of the accretion disc spectrum — a state called hard state. If, on the other hand, the Keplerian accretion rate ($\dot{M}_K$) is higher, then it supplies more soft photons to cool down the CENBOL region. This results in more power to the soft end of the accretion disc spectrum — a state known as the soft state. Recently, however, Chakrabarti & Mandal (2003) showed that raising the Keplerian rate even higher not necessarily softens the spectrum, for, the Keplerian flow also adds to the number density of electrons in the post-shock region and at some point, the spectrum starts to be hardened once more.

This kind of hybrid disc structure is known as the Two Component Accretion Flow or the TCAF disc (see, CT95, CTKE 96, Ebisawa et al. 1996, C97), and has wide observational support (Smith et al. 2001, 2002). In Fig. (1), a schematic diagram of such a disc is presented. The Figure shows the Keplerian disc, is sandwiched by the sub-Keplerian halo. The shock location ($x_s$) and the compact object are also shown. Jets are shown close to the axis of symmetry. Thus jets are illuminated by the Keplerian disc with soft photons and by the CENBOL with hard photons.

Chakrabarti and his collaborators have also showed that the CENBOL can drive a part of
the infalling matter along the axis of symmetry to form jets (Chakrabarti, 1998; Chakrabarti, 1999; Das & Chakrabarti, 1999; Das et al. 2001). There are wide support that the jets are indeed coming out from a region within $50 - 100$ Schwarzschild radius of the black hole (Junor, Biretta & Livio, 1999). Similarly, it is believed that jets are produced only in hard states (see, Gallo, Fender & Pooley, 2003, and references therein). Thus it is natural to study interaction of hard radiation from the CENBOL and the outflowing jets, with the particular interest of studying, whether momentum deposited to the jet material by these hard photons can accelerate them to ultra-relativistic speeds.

Investigation of interaction of radiation and astrophysical flow is not new. A number of astrophysicists have directed their efforts in this particular field of study while the consideration of the associated accretion disc depended on their personal choice or the popularity of the particular model of accretion at the given time. Icke (1980) studied the effect of radiative acceleration of the gas flow above a Keplerian disc. But the effect of radiation pressure on the gas flow was ignored. Sikora and Wilson (1981) showed that even if the jet is collimated by geometrically thick discs (Lynden-Bell, 1978; Abramowicz and Piran, 1980), radiation drag is important for astrophysical jets. Piran (1982) while calculating the radiative acceleration of outflows about the rotation axis of thick accretion discs, found out that in order to accelerate outflows to $\gamma > 1.5$ (where $\gamma$ is the bulk Lorentz factor), the funnels must be short and steep, but such funnels are found to be unstable. Sol et. al. (1989) proposed a two-flow model for jets, one consists of relativistic particles (electrons and positrons) and of relativistic Lorentz factor, while the other being normal, mildly relativistic plasma. In a very important paper, Icke (1989) considered blobby jets about the axis of symmetry of thin discs and he obtained the ‘magic speed’ of $v_m = 0.451c$ ($c$ — the velocity of light), $v_m$ being the upper limit of terminal speed. Fukue (1996) extended this study for rotating flow above a thin disc and drew similar conclusions, although for rotating flow, away from axis of symmetry the terminal speed was found out to be a little less than the magic speed of Icke. To summarize, earlier works showed that it is not possible to accelerate jets to ultra relativistic terminal velocities, by radiations from the earlier accretion disc models. What is more discouraging is the existence of moderate levels of equilibrium speed ($v_{eq}$ i.e., speeds above which there would be radiative deceleration). Recently Fukue et al. (2001) has done similar investigations on interaction of radiation and pair plasma jets, but they took a disc model which consisted of inner ADAF region (non luminous) and outer slim disc (luminous), which resulted in a relativistic terminal speed.
We are working in a different regime, i.e., less luminous Keplerian disc and more luminous post-shock torus or CENBOL. Since hard radiations are expected to emerge out of the optically thin CENBOL, its intensity (photon counts per unit area per unit time) is low, but it ‘looks’ directly into the jet vertically above and hence eventually deposit its momentum into the later. Radiation from a hot CENBOL is likely a source of pair-production and hence the possibility of radiative momentum deposition is likely to be higher even for radiations from CENBOL hitting the outflow at an angle (see, e.g., Yamasaki, Takahara, and Kusunose, 1999 for the mechanism of pair-production from hot accretion flows.). This is why we believe that the direct deposition of momentum may be important. Chattopadhyay & Chakrabarti (2002a) showed that hard radiations from the post-shock region (CENBOL) do accelerate electron-proton plasma to mildly relativistic terminal speeds. Chattopadhyay & Chakrabarti (2002b) also reported that hard radiations from the CENBOL do not impose any upper limit for terminal speed.

In the present paper, we solve the equations of photo-hydrodynamics of jets for radiations coming out from the TCAF discs, where the radiation fields from both the inner CENBOL and the outer Keplerian thin disc are considered. We show that, while the equilibrium velocity closer to the black hole depends on $x_s$, and the ratio between Keplerian disc and CENBOL luminosities, the terminal speed or the jet velocity at infinity depends on the relative proportions and also on the actual magnitude of various moments of radiation. We also show that, in hard states (in our parlance, CENBOL radiation dominating over Keplerian radiation), optically thin jets can be accelerated to ultra-relativistic speed.

In the next Section, we present the model assumptions and the equations of motion and compute various moments of radiation field. In §3, we present our solutions and finally in §4, we draw our conclusions.

2 ASSUMPTIONS, GOVERNING EQUATIONS AND COMPUTATION OF THE MOMENTS OF RADIATION FIELD

2.1 Assumptions and Governing Equations

In our analysis, the curvature effects due to the presence of the central black hole mass is ignored. The metric is given by, $ds^2 = c^2 dt^2 - dr^2 - r^2 d\phi^2 - dz^2$, where, $r$, $\phi$, and $z$ are the usual coordinates in cylindrical geometry and $ds$ is the metric in four-space. The four-velocities are $u^\mu$. We follow the convention where the Greek indices signify all four components and the
Latin indices represent only the spatial ones. The black hole is assumed to be non-rotating and hence the strong gravity is taken care of by the so-called Paczyński-Wiita potential (e.g., Paczyński & Wiita, 1980).

We also do not consider generation mechanism of jets. As the astrophysical jets are observed to be extremely collimated (Bridle & Perley, 1984), and generally aligned along the normal to the host galaxy, we assume the jet to be along the axis of symmetry. Thus the transverse structure of the jet is ignored, i.e., \( u^r = u^\phi = 0 \) and \( \partial/\partial r = \partial/\partial \phi = 0 \), where \( u^r \) and \( u^\phi \) are radial and azimuthal components of four velocity. We are looking for steady state solutions. Hence, \( \partial/\partial t = 0 \). We also assume the gas pressure is negligible compared to the radiation pressure. This is perhaps the case especially inside the funnel wall close to the axis. The derivation of the equations of motion of radiation hydrodynamics for optically thin plasma, was investigated by a number of workers. A detailed account of such derivation has been presented by Mihalas & Mihalas (1984; hereafter MM84) and Kato et al. (1998; hereafter K98), and are not presented here. Enforcing the above assumptions, the equation of motion presented in MM84 and K98, reduces to:

\[
\frac{u^z}{dz} = -\frac{GM_B}{(z-2)^2} + \frac{\sigma_T}{m} \left[ \gamma \frac{F^z}{c} - \gamma^2 E u^z - u^z P^{zz} + u^z (2\gamma u^z F^z - u^z u^z P^{zz}) \right],
\]

where, \( u^z \) is the \( z \)-component of four velocity, \( G, M_B, \sigma_T, \) and \( m \) are the universal gravitation constant, mass of the black hole, Thomson scattering cross-section and mass of the gas particle, respectively. \( E, F^z \) and \( P^{zz} \) are the radiative energy density, radiative flux and radiative pressure on the axis of symmetry, and \( \gamma(\gamma = u_t) \) is the Lorentz factor. The above equation can be re-written as,

\[
\frac{u^z}{dz} = -\frac{GM_B}{(z-2)^2} + \left[ \gamma F - \gamma^2 E u^z - u^z \mathcal{P} + u^z (2\gamma u^z F - u^z \gamma \mathcal{P}) \right],
\]

For simplicity, we will not compute the shock location \( x_s \) or the the CENBOL luminosity \( (L_C) \) – instead, we will supply them as free parameters. They can be easily computed from accretion parameters (e.g., C89, CT95, Das et al. 2001, Chattopadhyay et al. 2003). We assume that the outflow is made up of purely electron-positron pair plasma.

### 2.2 Computation of radiative moments from TCAF disc

The radiation reaching each point on the jet axis, is coming from two parts of the disc, namely, the CENBOL and Keplerian disc, hence all the radiative moments should have both the contributions.
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Figure 2. Schematic diagram of a two component accretion flow (TCAF). $O$ is the position of the black hole. Centrifugal pressure dominated boundary layer (CENBOL) is the puffed up region between $O$ and $x_s$, the shock location. Thick line between $x_s$ and $x_o$ is the Keplerian disc where $x_o$ is the outer boundary of the Keplerian disc. $DC$ and $D'C'$ are the local normals of the CENBOL and Keplerian disc. $B(0, z)$ is the field point where various moments of radiation fields are computed. $D(r', z')$ is source point on the CENBOL and $D'(r'_K, 0)$ is the source point on the Keplerian disc. $r'$ is radial coordinate of the CENBOL and $r'_K$ is that of the Keplerian disc. The sub-Keplerian halo is not shown.

In Fig. (2), a schematic diagram of the cross-section of the disc structure is presented. The black hole is situated at $O$. The region bounded between $O$ and $x_s$ is the CENBOL, and the thick line between $x_s$ and $x_o$ is the Keplerian disc, where $x_o$ is the outer boundary of the Keplerian disc. Thus $x_s$ is the outer boundary of CENBOL and inner boundary of the Keplerian disc. The inner boundary of Keplerian disc, as seen from the position $B$ is $x_{K_i}$. The shock height is $h_s \sim a_s x_s^{1/2}(x_s - 1)$, where $a_s$ is the equatorial sound speed at shock and depends on $x_s$. In other words $h_s = h_s(x_s)$, but as we are not solving the accretion disc equations simultaneously, we have to make some estimate of $h_s$ which will closely mimic reality. Chakrabarti solutions (C89, C90, CT95, C96) show that if $x_s = 10r_g$ (where $r_g$ is the Schwarzschild radius), then the temperature at the shock is $T_{10} \sim 1.56 \times 10^{11} K$. Assuming that the shock temperature to be $T_s = T_{10}(10/x_s)$, one can estimate the shock height to be $h_s \sim 0.6(x_s - 1)$. The inner surface of the CENBOL is assumed to be conical. The radiations from a point $D(r', z')$ on the surface of CENBOL is primarily along the local normal $DC$. Similarly, the radiation from a point $D'(r'_K, 0)$ on the Keplerian disc, is along its local normal $D'C'$.

The radiative moments at $B$ are:

$$E = \frac{1}{c} \int I d\Omega = \frac{1}{c} \left( \int_C I_C d\Omega_C + \int_K I_K d\Omega_K \right),$$

(2a)
\[
\frac{F_i^i}{c} = \frac{1}{c} \int I^i d\Omega = \frac{1}{c} \left( \int_C I_C^i l_C^i d\Omega_C + \int_K I_K^i l_K^i d\Omega_K \right), \tag{2b}
\]

and
\[
P^{ij} = \frac{1}{c} \int I^i l^j d\Omega = \frac{1}{c} \left( \int_C I_C^i l_C^j d\Omega_C + \int_K I_K^i l_K^j d\Omega_K \right). \tag{2c}
\]

In the Eqs. (2a-2c), \(I, d\Omega, l^i\) are the frequency integrated intensity from the disc, differential solid angle at \(B\), and \(l^i\) are the direction cosines at \(B\), for example \(l_C^z = (z - z')/BD\) and \(l_K^z = z/BD'\). Suffix \(C\) and \(K\) represent quantities linked to CENBOL and the Keplerian disc respectively. The expressions of solid angles subtended at \(B\) from \(D\) and \(D'\) are given respectively by,
\[
d\Omega_C = \frac{r' \csc \theta dr' d\phi}{r'^2 + (z - z')^2} \cos(\angle CD B), \tag{3a}
\]
and
\[
d\Omega_K = \frac{r_K' dr_K d\phi}{r_K'^2 + z^2} \cos(\angle C'D'B), \tag{3b}
\]

where, \(\theta\) is the semi-vertical angle of the inner surface of the CENBOL which, for simplicity, is assumed to be constant. It is to be noted that, from Fig. (2) and Eqs. (3a-3b), in general, \(l_C^z > l_K^z\) and for any unit differential area on the disc, \(d\Omega_C > d\Omega_K\). This implies that the contribution from the CENBOL to the total radiation field moment is much more, compared to that from the Keplerian disc. This will be clear when the comparative study of the moments are presented later [see Fig. (3)].

It is clear from Eq. (2b) that, due to the symmetry about the \(z\)-axis, only non-zero pressure tensor components are \(P^{ii}\) and \(F^r = F^\phi = 0\) on the axis. Only \(F^z \neq 0\). As we consider jets on or about the axis of symmetry and the transverse structure is ignored, \(P^{rr}\) and \(P^{\phi\phi}\) do not enter the equation of motion as these components are only coupled with \(u^r(= 0)\) and \(u^\phi(= 0)\), i.e., we have to only compute \(P^{zz}\), which is exactly what is seen in Eq. (1).

From Fig. (2), it is clear that for the outflowing matter within the funnel, i.e., when \(z < [h_s x_o/(x_o - x_s)]\), radiations from the Keplerian disc do not reach the electrons because these radiations are intercepted by CENBOL. CENBOL will reprocess these intercepted photons, and re-emit them in all directions, especially towards the jets because of especial geometry and directiveness of the CENBOL surface. Due to shadowing effect mentioned above, the radiation from the Keplerian disc reaching \(B\) is from \(x_{Ki}\) to \(x_o\), where \(x_{Ki} = x_s z/(z - h_s)\), with the additional constraints, \(x_s \leq x_{Ki} \leq x_o\) and \(x_{Ki} > 0\).

As far as the CENBOL properties are concerned, we follow CT95, where an effective
temperature was computed for the CENBOL and the radiation intensity was chosen to be uniform. That is: \( I_C = \frac{L_C}{\pi A} = \ell L_{Edd}/\pi A \) = constant, where \( L_C \) and \( A \) are the CENBOL luminosity and the surface area of the CENBOL respectively. \( L_{Edd} \) is the Eddington luminosity and \( \ell \) is the CENBOL luminosity in units of \( L_{Edd} \). The Keplerian disc intensity per unit solid angle is \( I_K = \frac{3GM_B\dot{M}_K}{8\pi r_K^2} \left(1 - \sqrt{\frac{r_g}{r_K}}\right) \) (NT73). Let us now multiply \( \sigma_T/m \) with Eqs. (2a-2c), and then integrate over the whole disc to obtain the following integrated quantities for the moments,

\[
\mathcal{E} = \mathcal{E}_C + \mathcal{E}_K,
\]

where, \( r_{in} \) is the inner boundary of the disc.

\[
\mathcal{F} = \mathcal{F}_C + \mathcal{F}_K
\]

\[
\mathcal{P} = \mathcal{P}_C + \mathcal{P}_K
\]

Constancy of \( I_C \) allows us to have analytical expressions for \( \tilde{E}_C, \tilde{F}_C \) and \( \tilde{P}_C \). These were computed by Chattopadhyay & Chakrabarti (2000, 2002ab) and are not repeated here. We choose the unit of length to be \( 2GM_B/c^2 \) — the Schwarzschild radius \( (r_g) \), unit of time to be \( 2GM_B/c^3 \) and \( M_B \) is the unit of mass. Thus the unit of velocity is \( c \). In such units, constants in Eqs. (4a-4c) are,

\[
\mathcal{E}_C = \mathcal{F}_C = \mathcal{P}_C = \frac{1.3 \times 10^{38} \ell \sigma_T}{2 \pi cmAGM_\odot}
\]

and

\[
\mathcal{E}_K = \mathcal{F}_K = \mathcal{P}_K = \frac{4.32 \times 10^{17} \dot{m}_K \sigma_{TC}}{32 \pi^2 mGM_\odot}.
\]
We have written the Keplerian accretion rate in units of the Eddington accretion rate $\dot{M}_{\text{Edd}}$, i.e., $\dot{m}_K = \dot{M}_K / \dot{M}_{\text{Edd}}$. It is to be noted that, two of the disc parameters i.e., $r_{in} = 1.5r_g$ and $x_o = 1000r_g$ are kept constant throughout the paper. In Fig. (3), the space dependent part of the moments $\tilde{E}$, $\tilde{F}$ and $\tilde{P}$ are compared for two shock locations. In Figs. 3(a-b), the contributions from the CENBOL and the Keplerian disk are plotted for $x_s = 10r_g$ and in Figs. 3(c-d) the contributions from the CENBOL and the Keplerian disc are plotted, where the shock location is $x_s = 20r_g$. We see that as $x_s$ is increased, the CENBOL contributions increases while the Keplerian contributions decrease. We also see another fact that generally, in the hard state, the CENBOL contribution to the total radiative moments dominates over Keplerian counterparts. We see from Fig. (3a) and (3c), that within the funnel (i.e., $z \leq h_s$), $\tilde{E}_C > \tilde{F}_C > \tilde{P}_C$. We also see that, because of the relatively small size of the CENBOL, $\tilde{E}_C \approx \tilde{F}_C \approx \tilde{P}_C$ for smaller values of $z$ [e.g., $z \sim 50.5r_g$ for Fig. (3a) and $z \sim 92.5r_g$ for Fig. (3c)], while we see $\tilde{E}_K \approx \tilde{F}_K \approx \tilde{P}_K$ at much larger distances [e.g., $z \sim 900r_g$ for Fig (3b) and $z \sim 1500r_g$ for Fig. (3d)]. The small size of CENBOL ensures that the direction cosines $l^z \rightarrow 1$ for smaller $z$. In contrast, the larger size of the Keplerian disc ensures $l^z \sim 1$ only for large $z$. From Figs. 3(a-d), it appears that as $\tilde{E}_C \gg \tilde{E}_K$, $\tilde{F}_C \gg \tilde{F}_K$, and $\tilde{P}_C \gg \tilde{P}_K$. So, if CENBOL is more luminous, then CENBOL contributions would dominate the Keplerian ones even at large $z$. While $E_{K0}$ ($= F_{K0} = P_{K0}$) depends only on $\dot{m}_K$, $E_{C0}$ ($= F_{C0} = P_{C0}$), depends on both $\ell$ and the CENBOL surface area. From Figs. (3a) and (3c), we see that with the increase of $x_s$, even if the radiative moments from CENBOL, say for example, $\tilde{E}_C$ increases moderately, but calculations show that $\tilde{E}_{C0}$ decreases appreciably with $x_s$ (e.g., $\tilde{E}_{C0}$ for $x_s = 10r_g$ is about four times larger than $\tilde{E}_{C0}$ for $x_s = 20r_g$). So Keplerian contributions might be comparable to that due to the CENBOL at large $z$, depending on specific combinations of all the parameters $\ell$, $\dot{m}_K$ and $x_s$, especially for large values of $x_s$ and low $\ell$. We also observe that the shadow effect of CENBOL ensures that at $z \leq (h_s x_o)/(x_o - x_s)$, $\tilde{E}_K = \tilde{F}_K = \tilde{P}_K = 0$.

3 RADIATIVE ACCELERATION

In the rest of the paper, we will use geometrical units defined in the last section, but for simplicity we will keep the same symbols representing variables as in Eq. (1). We define a three-velocity $v$ such that $v^2 = -u_i u^i / u_t u^t = -u_z u^z / u_t u^t$. Thus $u^z = \gamma v$ and $u_z = -\gamma v$. 

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Under such considerations Eq. (1) takes the form,

$$
\frac{dv}{dz} = -\frac{1}{2(z-1)^2} + \left[ \gamma F - \gamma^3 v E - \gamma v P + \gamma^3 (2v^2 F - v^3 P) \right] \frac{1}{\gamma^4 v}.
$$

The first term in the right-hand side of Eq. (6) is the gravitational term with dimensionless Paczynski-Wiita potential and the term in the square bracket is the radiative acceleration term. We notice that the radiative acceleration term depends both on $v$ and the radiative moments. The first and fourth term in the square bracket are accelerating terms while the second, third and fifth terms are the decelerating terms, collectively known as the radiation
drag terms. We also see that there is a $\gamma^4$ term in the denominator of the r. h. s of Eq. (6). This term ensures that, apart from radiation drag, outflowing matter will be slowed down as $v\rightarrow 1$.

### 3.1 Equilibrium velocity

Let us now discuss the concept of equilibrium velocity. As a concept, equilibrium velocity is not a new one and has been extensively discussed by a number of astrophysicists (see, Fukue 2003, and references there in for details). We want to study this issue in the context of jets in the radiation field of TCAF discs. It is defined in a manner that at $v = v_{eq}$ the square bracket term in Eq. (6) is zero. Thus, for $v > v_{eq}$ there is deceleration and for $v < v_{eq}$ there is radiative acceleration.

Putting the square bracket term in Eq. (6) to zero we have,

$$ F v_{eq}^2 - (\mathcal{E} + \mathcal{P}) v_{eq} + \mathcal{F} = 0. $$

This is a quadratic equation of $v_{eq}$ whose solution is,

$$ v_{eq}(z) = \frac{(\mathcal{E} + \mathcal{P}) - \sqrt{(\mathcal{E} + \mathcal{P})^2 - 4F^2}}{2F} = \xi - \sqrt{\xi^2 - 1}, $$

where $\xi = (\mathcal{E} + \mathcal{P})/2F$. From Eqs. (4a-5b), we have,

$$ \xi(z) = \frac{(\tilde{E}_C + \tilde{P}_C) + \frac{\dot{m}_K}{\ell} \tilde{E}_K + \tilde{P}_K)}{2(\tilde{E}_C + \frac{\dot{m}_K}{\ell} \tilde{E}_K)}, $$

where $\zeta = 6.6 \times 10^{-13} c^2 A$. It is to be noted that $\xi$ depends on $x_s$ as well as on the ratio $\dot{m}_K/\ell$, but not separately $\dot{m}_K$ and $\ell$. In case $(\dot{m}_K \zeta)/\ell \ll 1$, $v_{eq}$ is completely determined by $\tilde{E}_C$, $\tilde{F}_C$ and $\tilde{P}_C$, and $\ell$ has no effect in determining $v_{eq}$. We also see that $\xi$ does not depend on the mass of the gas particles $m$. Thus $\xi$ (and thus $v_{eq}$) is the same for both electron-proton plasma as well electron-positron plasma, provided $x_s$ and $\dot{m}_K/\ell$ is the same in both the cases. It is clear from Eqs. (8-9), that as $\mathcal{E} \approx \mathcal{F} \approx \mathcal{P}$, $v_{eq} \rightarrow 1$. Therefore if the CENBOL contribution dominates then, $v_{eq} \sim 1$ within few tens of Schwarzschild radii above the disc plane. If Keplerian radiation dominates, then the condition is achieved at much larger distance, as we shall see later. We also see that no outflow is possible i.e., $v_{eq} \leq 0$, i.e., if

$$ \mathcal{F} \leq 0. $$

\[\xi \geq 0 \]
From Fig. (3) we see that very close to the horizon, due to the torus geometry of the CENBOL, \( F < 0 \), hence very close to the black hole, not only enormous gravitational pull but also the radiative force pushes matter inward.

**Icke’s magic speed:**

Let us now recover an important result from Eq. (8). If we consider a thin-Keplerian disc of infinite size, then \( P_{zz}^K = P_{rr}^K = P_{\phi\phi}^K = \frac{1}{3} E_K \) and \( F_z^K = \frac{2}{3} E_K \), and there is no CENBOL in this particular case. Under such conditions \( \mathcal{E} = 2F = 3P \), putting these in Eq. (8) we have,

\[
v_{eq} = v_m = \frac{1}{3} (4 - \sqrt{7}) = 0.451 \equiv \text{magic speed of Icke !}
\]

Thus we see that, if the radiation field, *i.e.*, the radiative properties of the disc, can be prescribed, then \( v_{eq} \) can be found out easily.

**Equilibrium speed from a TCAF disc:**

Let us now concentrate on the radiations from TCAF discs. As discussed in the introduction, jets are produced in the hard state of the accretion disc when the CENBOL is hotter and hard state means more power is on the high energy end of the spectrum. That is, \( L_C > L_K \).

The Keplerian disc luminosity is given by

\[
L_K = r_g^2 \int_{x_o}^{x_s} 2\pi I_K 2\pi r_K dr_K,
\]

which is a function of \( \dot{m}_K, x_s \) and \( x_o \), which when integrated can be expressed as,

\[
L_K = \frac{3}{4} \dot{m}_K \left[ -\frac{1}{r_K} + \frac{2}{3r_K} \sqrt{\frac{3}{r_K}} \right]_{x_o}^{x_s} L_{Edd}.
\]

Thus the Keplerian disk luminosity in units of the Eddington luminosity can be defined as \( \ell_K = L_K/L_{Edd} \). As has been stated before, presently we do not compute \( L_C \), but supply it. Typical values of \( \ell \) that we shall employ should, in general, depend on \( \ell_K \) itself because \( \ell \sim \Lambda \ell_K \), where \( \Lambda \) is the enhancement factor by which incident photon intensity is increased due to Comptonization and has a value of around 20 – 30 in hard states (CT95). Thus, for instance, if \( \ell_K \sim 0.05 \), a typical value of \( \ell \sim 0.1 - 0.15 \). However, while we shall choose \( \ell \) and \( \ell_K \) in these regions, we shall use them as free parameters, since presently we are not interested in computing the terminal speed as a function of the spectral slope \( \alpha \), though, strictly speaking, it would be a function of \( \alpha \).

In our case, we see that, there are two sources of radiation from the disc, (i) the CENBOL and (ii) the Keplerian disc. Let us first concentrate on the CENBOL contribution. From Figs. (3a) and (3c), we see that for \( z \rightarrow \text{few} \times 10 r_g, (\mathcal{E}_C + \mathcal{P}_C) \sim 2F_C \), hence Eq. (8) determines \( v_{eq} \sim 1 \) at such distances.

In general, in the hard state also, \( L_K \) is not negligible. Though at \( z \rightarrow \text{large}, \tilde{E}_K \sim \tilde{F}_K \approx \tilde{P}_K \), but at \( z \sim \text{few} \times 10 r_g, (\mathcal{E}_K + \mathcal{P}_K) > 2F_K \). Hence for higher Keplerian luminosity, \( v_{eq} \sim 1 \).
is achieved at distances around a thousand Schwarzschild radii. In Figs. (4a) and (4b), we have plotted \(v_{eq}\) and \(\xi\) with \(\log(z)\). Various curves correspond to \(\dot{m}_K \sim 0.01\) (solid), \(\dot{m}_K = 1\) (dashed) and \(\dot{m}_K = 6\) (long-dashed). We choose \(\ell = 0.3\) and \(x_s = 20r_g\) for all the plots. We see that close to the black hole, \(v_{eq}\) is independent of \(\ell_K\), because at such distances radiation from the CENBOL dominates. If \(L_C \gg L_K\) \((i.e.,\ solid\ curve)\), then we see that \(v_{eq} \rightarrow 1\) at around \(100r_g\). For \(\dot{m}_K = 1\) \((\equiv \ell_K \sim 0.03, \ i.e.,\ dashed\ curve)\), and for \(\dot{m}_K = 6\) \((i.e.,\ long-dashed\ curve)\), \(v_{eq} \rightarrow 1\) at distances over \(1000r_g\). One should note that at \(z < 1.85r_g\), \(v_{eq} < 0\). The reason for this can be clearly seen from Fig. (3c), which shows that \(\mathcal{F} < 0\)
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at the same distance from the black hole. This means that very close to the black hole not only gravity pushes in matter but also radiation would also push matter inside, thus vindicating Eq. (10). To clarify this point we have plotted $\xi$ with $\log(z)$ in Fig. (4b). Curve styles match those of the corresponding case in Fig. (4a). We see that within the funnel like region of the CENBOL, the radiation is completely dominated by the CENBOL itself. And also due to this particular geometry $\xi > 1$, resulting $v_{eq}$ much lesser than the velocity of light. In case $\dot{m}_K \to$ small (solid), for $z > 100r_g$, $\xi \to 1$, resulting $v_{eq} \to 1$. In case $\dot{m}_K \to$ large (dashed and long-dashed curves), only at around a thousand Schwarzschild radii, $\xi \sim 1$. This means that, if one increases $\dot{m}_K$, then the higher velocities achieved due to the acceleration of jets by CENBOL photons, might be decelerated. It is quite obvious though, that the effect of Keplerian radiation is quite marginal. Of course one should keep in mind that if the radiative moments due to CENBOL and Keplerian radiation are comparable at infinite distances, then there is a possibility of increased terminal speed, with the increase in $\dot{m}_K$.

We have seen from Eq. (9), that $\xi$ not only depends on $\dot{m}_K/\ell$ but also on $x_s$. We now investigate the dependence of $v_{eq}$ on $x_s$. In Fig. (4c), $v_{eq}$ is plotted with $\log(z)$ for $\ell = 0.12$ and $\dot{m}_K = 0.5$. Various curves represent $x_s = 10r_g$ (solid), $x_s = 20r_g$ (dashed) and $x_s = 30r_g$ (long-dashed). With the increase in $x_s$, as the size of CENBOL increases, hence it can only behave like a point source (for which $\xi \to 1$) farther out, and we see that $v_{eq} \to 1$ farther outwards from the black hole. To clarify, we have also plotted $\xi$ with $\log(z)$ for $\ell = 0.12$ and $\dot{m}_K = 0.5$ in Fig. (4d). Various curves represent $x_s = 10r_g$ (solid), $x_s = 20r_g$ (dashed) and $x_s = 30r_g$ (long-dashed). As has been just explained, we see that $\xi \to 1$ for larger value of $z$, as $x_s$ is increased.

3.2 Velocity profile

So far, we have only discussed about the equilibrium velocity and its dependence on $\ell$, $\dot{m}_K$ and $x_s$. This was done to study the upper limit of allowed velocity as a function of $z$. We have seen that, if CENBOL radiation dominates over that from the Keplerian disc, then jets can be accelerated to very high velocities within around $100r_g$, but $v_{eq}$ is only a measure of velocity which signifies only the domain of radiative deceleration or acceleration, and is not the actual velocity. From Eq. (6) it is quite clear that the radiative acceleration and hence $v$ itself, depends on all three quantities $\mathcal{E}$, $\mathcal{F}$, $\mathcal{P}$ and not only just on $\xi$. Equation (6) also shows, radiative acceleration also depends on $v$ in a very complicated way. To compute actual
velocity, one has to integrate Eq. (6). Jets are believed to be produced from the post shock region (Chakrabarti 1999, Das & Chakrabarti 1999, Das et al. 2001). If they are generated from the post-shock region then they should start with very small velocity. In this paper we are not concentrating on the generation of jets. Thus we just put the injection height to be close to the black hole and the injection velocity to be small. We choose injection parameters to be $z_{in} = 5r_g$ and $v_{in} = 10^{-3}$. Before discussing the results, let us define terminal speed. The terminal speed ($\vartheta$) is the constant velocity at infinite distances or $\vartheta = v|_{z \rightarrow \infty}$. In Fig. (5a), three velocity $v$ is plotted with $\log(z)$, where $\dot{m}_K = 0.5$ and $x_s = 20r_g$. Various curves correspond to $\ell = 0.36$ (solid), $\ell = 0.24$ (dashed) and $\ell = 0.12$ (long-dashed). We see that close to the black hole as $v \ll 1$, $F$ dominates resulting in a steep rise in $v$, as $v$ increases the jet starts to feel the drag force, resulting in a somewhat less steep increase in $v$. At $z > 100r_g$, the radiative moments tend to become weak and the jet settles to a constant velocity at large $z$ or the terminal speed $\vartheta$. It is clear that, $v$ increases with $\ell$ and in the three cases depicted, the terminal velocity of jets are $\vartheta \sim 0.93$ (solid), $\vartheta \sim 0.91$ (dashed) and $\vartheta \sim 0.88$ (long-dashed). It may seem curious that even if $\ell$ is increased by equal interval the terminal speed achieved, does not increase by equal interval. The reason is two fold, first of all from

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure5.png}
\caption{(a) Variation of $v$ with $\log(z)$, for $\dot{m}_K = 0.5$. Various curves correspond to $\ell = 0.36$ (solid), $\ell = 0.24$ (dashed) and $\ell = 0.12$ (long-dashed). (b) Variation of $\vartheta$ with $\ell$. Various curves correspond to $\dot{m}_K = 0.4$ (solid), $\dot{m}_K = 1.4$ (dashed) and $\dot{m}_K = 2.4$ (long-dashed). The shock location $x_s = 20r_g$, for both the cases. Injection parameters $z_{in} = 5r_g$, $v_{in} = 10^{-3}$.}
\end{figure}
Eq. (6), one can see that the gradient of $v$ is a non-linear function of the moments. The second reason being the existence of $\gamma^4$ term in Eq. (6), which suppresses the acceleration, as $v$ increases to values close to that of light.

In Fig. (5b), the terminal speed $\vartheta$ is plotted with $\ell$, for $x_s = 20r_g$. Various curves correspond to $\dot{m}_K = 0.4$ (solid), $\dot{m}_K = 1.4$ (dashed) and $\dot{m}_K = 2.4$ (long-dashed). The terminal speed $\vartheta$ increases with $\ell$, though for higher $\ell$, the increase of $\vartheta$ decreases, for the same two reasons discussed for the previous Figure. We also see that the dependence of $\vartheta$ on $\dot{m}_K$ is marginal. In Fig. (3), we have seen that the CENBOL contribution to the various space dependent part of the radiative moments, is a few orders of magnitude higher than that from the Keplerian contribution, hence the marginal dependence of $\vartheta$ with $\dot{m}_K$ is expected. It is to be noted though, that for $\ell > 0.15$, $\vartheta$ decreases with increasing $\dot{m}_K$ and for $\ell < 0.15$, $\vartheta$ increases with increasing $\dot{m}_K$, albeit the dependence is very weak. For the higher values of $\ell$, the jet, powered by the CENBOL radiation, achieves very high velocity within a few tens of Schwarzschild radii. But for the next thousand of Schwarzschild radii or so, $\xi > 1$, as $(E_K + P_K) > 2F_K$ in that region [see, Fig. (3c-3d)]. When $v$ becomes high, the drag force increases, and slows down the jet. At a larger distance, a slight increase in radiative moments, due to the increase in $\dot{m}_K$, is not sufficient to increase $\vartheta$. When $\ell$ is small, the velocity achieved within few tens of Schwarzschild radii is not that high, hence the drag force is less. But at large distances, the radiative moments due to CENBOL and Keplerian disc becomes comparable hence there is a slight increase in $\vartheta$ with $\dot{m}_K$.

In Fig. (6), $\vartheta$ is plotted with $\dot{m}_K$, for $x_s = 20r_g$. Various curves correspond to $\ell = 0.3$ (solid), $\ell = 0.24$ (dashed), $\ell = 0.18$ (long dashed), and $\ell = 0.12$ (dashed-dotted). We see that, as in Fig. (5b), for higher $\ell$, $\vartheta$ decreases with increasing $\dot{m}_K$ (see, solid, dashed and long dashed curves), but for lower values of $\ell$, $\vartheta$ increases with increasing $\dot{m}_K$ (see, dashed-dotted curve). So one can conclude that in the hard state, jets are basically accelerated by radiations from CENBOL and the terminal velocity has a weak dependence on radiation from Keplerian disc.

Let us now investigate the dependence on $x_s$. In Fig. (7a), $v$ is plotted with $\log(z)$, for $\ell = 0.18$ and $\dot{m}_K = 0.5$. Various curves correspond to $x_s = 10r_g$ (solid), $x_s = 20r_g$ (dashed) and $x_s = 30r_g$ (long-dashed). We see that with the increase in $x_s$, the acceleration is reduced. For $x_s = 10r_g$ (solid), we see that the jet is experiencing tremendous acceleration and within a short distance of $z \sim 40r_g$, it achieves a velocity $\sim 0.9$, and gradually settles to a terminal value of $\vartheta = 0.93$. As the shock location is increased to $x_s = 20r_g$ (dashed) and $x_s = 30r_g$
Figure 6. Variation of $\vartheta$ with $\dot{m}_K$, for $x_s = 20r_g$. Various curves correspond to $\ell = 0.3$ (solid), $\ell = 0.24$ (dashed) and $\ell = 0.18$ (long-dashed) and $\ell = 0.12$ (dashed-dotted). Injection parameters are $z_{in} = 5r_g$, $v_{in} = 10^{-3}$.

(long-dashed), acceleration is weaker and the terminal speeds achieved are $\vartheta = 0.9$ and $\vartheta = 0.88$ respectively. Thus, the terminal speed depends strongly on shock location and decreases with increasing shock location.

In Fig. (7b), we have plotted $\vartheta$ with $\ell$. The solid curves correspond to $\dot{m}_K = 1.4$ and the dashed curves correspond to $\dot{m}_K = 2.4$. The shock locations $x_s$, are marked on each pair of solid and dashed curves ($x_s = 10r_g, 20r_g, 30r_g$) in the Figure. We see that in general $\vartheta$ increases with $\ell$. We also see that, for particular values of $x_s$, $\dot{m}_K$ has very limited influence on $\vartheta$ (see each pair of solid and dashed curves marked with values of $x_s = 10, 20, 30$), and also that with increasing $x_s$, $\dot{m}_K$ is less and less effective in determining $\vartheta$. In particular, the solid and the dashed curves marked $x_s = 10$ are distinguishable, but are increasingly less distinguishable for $x_s = 20$ and $x_s = 30$. With the increase in $x_s$, the inner edge of the Keplerian disc is increased. We also know that the magnitude of Keplerian disc intensity is more, closer to the black hole, i.e., if $x_s$ increases then the size of the Keplerian disc as well as Keplerian luminosity decreases. Thus with the increase of $x_s$, Keplerian radiation will be less effective in determining $\vartheta$. Similar to Fig. (5b), we also notice that if $\ell > 0.11$ for the pair of curves (solid and dashed) marked 10 (i.e., $x_s = 10$), then $\vartheta$ increases with decreasing $\dot{m}_K$, but decreases with decreasing $\dot{m}_K$, for $\ell < 0.11$. Though this crossing over value of $\ell$,
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Figure 7. (a) Variation of $v$ with $\log(z)$ for $\ell = 0.18$ and $\dot{m}_K = 0.5$. Various curves correspond to $x_s = 10r_g$ (solid), $x_s = 20r_g$ (dashed) and $x_s = 30r_g$ (long-dashed). (b) Variation of $\vartheta$ with $\ell$. The solid curves correspond to $\dot{m}_K = 1.4$ and dashed curves correspond to $\dot{m}_K = 2.4$. The shock locations $x_s$, are marked on each pair of solid and dashed curves ($x_s = 10r_g, 20r_g, 30r_g$). (c) Variation of $\vartheta$ with $x_s$, for $\dot{m}_K = 0.5$. Various curves represent $\ell = 0.24$ (solid), $\ell = 0.18$ (dashed), $\ell = 0.12$ (long dashed) and $\ell = 0.06$ (dash-dot). (d) Variation of $\vartheta$ with $x_s$, for $\ell = 0.18$. Various curves represent $\dot{m}_K = 1.5$ (solid) $\dot{m}_K = 3$ (dashed) and $\dot{m}_K = 4.5$ (long dashed). Injection parameters $z_{in} = 5r_g$, $v_{in} = 10^{-3}$.

increases with increasing $x_s$, i.e., for $x_s = 20r_g$, the crossing over occurs at $\ell = 0.15$, and for $x_s = 30r_g$ this occurs at $\ell = 0.24$. We also see that, generally, $\vartheta$ decreases with increasing $x_s$. In Fig. (7c), $\vartheta$ is plotted with $x_s$, for $\dot{m}_K = 0.5$. Various curves correspond to $\ell = 0.24$ (solid), $\ell = 0.18$ (dashed), $\ell = 0.12$ (long-dashed) and $\ell = 0.06$ (dashed-dotted). It is clear that $\vartheta$ decreases with $x_s$, though for a fixed value of $x_s$, $\vartheta$ increases with $\ell$. In Fig. (7d), $\vartheta$ is plotted with $x_s$, for constant values of $\ell = 0.18$. Different curves corresponds to $\dot{m}_K = 1.5$ (solid), $\dot{m}_K = 3$ (dashed) and $\dot{m}_K = 4.5$ (long-dashed). The first thing to notice is the weak
dependence of $\vartheta$ on $\dot{m}_K$. We also notice that, for $x_s < 22r_g$, $\vartheta$ decreases with increasing $\dot{m}_K$ and for $x_s > 22r_g$, $\vartheta$ increases with increasing $\dot{m}_K$. The reason for this is that the CENBOL intensity falls with increasing $x_s$. Thus for a larger $x_s$ the radiative moments due to the CENBOL at infinite distances are comparable to that due to the Keplerian disc. Hence with increasing $\dot{m}_K$, $\vartheta$ increases, exactly for the same reason as has been presented while discussing Fig. (5b). In general, one can conclude from Figs. (7a-7d), that jets can be accelerated to, up and around 90% the velocity of light, provided the $x_s \rightarrow 10r_g - 20r_g$, and $\ell \rightarrow 0.1 - 0.2$. As we discussed earlier, such values of $\ell$ are not unreasonable when amplification of photon energy takes place due to Comptonization and $\tau \sim 1$. For smaller $\tau$, amplification factor is higher (so that $\ell \sim 1$ is achievable) but in that case, both the number of very energy photons goes down and also the efficiency to deposite radiative momentum goes down dramatically.

In the preceding paragraphs of this subsection, we have studied the issue of radiative acceleration of jets, and on its dependence on three parameters namely $x_s$, $\ell$ and $\dot{m}_K$. In case of CENBOL radiation, the information we have provided is through its total luminosity, but Keplerian luminosity is governed by two parameters $x_s$ and $\dot{m}_K$ [see, Eq. (11)]. To have a better understanding, now we study how the relative proportions of CENBOL and Keplerian luminosity affect the terminal speed of jet.

In Fig. (8a), $\vartheta$ is plotted with $\ell_K/\ell$ — the ratio of Keplerian and CENBOL luminosities, for $x_s = 20r_g$. Various curves correspond to $\ell = 0.3$ (solid), $\ell = 0.24$ (dashed), $\ell = 0.18$ (long-dashed) and $\ell = 0.12$ (dashed-dotted). The ratio of luminosities is kept less than one to mimic the hard state of the accretion disc. We see that for $\ell = 0.12$, $\vartheta$ increases with the increase of Keplerian luminosity. For higher CENBOL luminosities (see, solid, dashed and long-dashed curves) we observe that $\vartheta$ decreases with increasing Keplerian luminosity. This has been addressed while discussing Fig. (5b), Fig. (6) and Fig. (7b), i.e., for $x_s = 20r_g$, if $\ell < 0.15$, $\vartheta$ increases with the increase of Keplerian luminosity. It is thus clear that if $x_s = 20r_g$ then for $\ell$ higher than 0.15, jets can be accelerated to very high terminal velocities. In order to study the dependence of shock location and the Keplerian luminosity, in Fig. (8b), we have plotted $\vartheta$ with $\ell_K/\ell$ for $\ell = 0.24$. The Keplerian luminosity thus is increased from 3% to 75% of CENBOL luminosity. Various curves correspond to $x_s = 10r_g$ (solid), $x_s = 20r_g$ (dashed) and $x_s = 30r_g$ (long-dashed). For $x_s = 10r_g$, $\vartheta$ is around 0.93 but decreases with the increase of Keplerian luminosity. With the increase in shock location the terminal velocity is less ($\vartheta = 0.91$ for $x_s = 20r_g$), but at the same time, decrement of the terminal velocity due to the increase of Keplerian luminosity is also less. For $x_s = 30r_g$,
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Figure 8. (a) Variation of \( \vartheta \) with \( \ell_K/\ell \) for \( x_s = 20 \) for all the curves. Various curves correspond to \( \ell = 0.3 \) (solid), \( \ell = 0.24 \) (dashed), \( \ell = 0.12 \) (long-dashed) and \( \ell = 0.12 \) (dashed-dotted). (b) Variation of \( \vartheta \) with \( \ell_K/\ell \) for \( \ell = 0.24 \). Various curves correspond to \( x_s = 10r_g \) (solid), \( x_s = 20r_g \) (dashed), \( x_s = 30r_g \) (long-dashed). Injection parameters \( z_{in} = 5r_g, v_{in} = 10^{-3} \).

\( \vartheta \approx 0.895 \) is lesser, but the change due to the increase of Keplerian luminosity is even less and \( \vartheta \) remains almost constant. Thus we conclude that if the shock location is between \( 10r_g - 20r_g \) then jets can be accelerated to terminal speeds above 90% of the velocity of light, for disc luminosities around 20% of the Eddington luminosity.

4 DISCUSSION AND CONCLUDING REMARKS

It is well known that high energy photons can produce particle-antiparticle pairs close to the inner edge of a disk. If the photon energy \( h\nu \gtrsim 2mc^2 \), then an electron-positron pair may be created, where \( h \) is the Planck’s constant, \( \nu \) is photon frequency and \( m \) is the electron (or positron) mass. If, on the other hand, electron and positron collide it will annihilate each other to produce two Gamma-ray photons, a process called pair annihilation. Clearly, to produce electron-positron jets the pair-production process has to dominate over pair annihilation. Workers (e.g. Mishra and Melia, 1993; Yamasaki, Takahara and Kusunose, 1999) discussed the production of electron-positron pairs from inner part of the accretion disc. Observationally, there are reports of pair dominated jets (Sunyaev et al. 1992; Mirabel and Rodriguez 1998; Wardle et al. 1998) from galactic black hole candidates to quasars.
Though there is little doubt on the existence of pair dominated jets, radiative acceleration of such jets on the other hand is a different issue altogether. If the number density of pairs created around the black hole is too high, then radiative acceleration would be ineffective. In this present paper, we have ignored the details of formation of the pair plasma jets and have only concentrated on the radiative acceleration of optically thin pair dominated jets.

In this paper, we supplied the CENBOL intensity \( I_C \), and the shock location \( x_s \) as free, but reasonable parameters. We have separately treated two velocity variables, (i) the equilibrium velocity \( v_{eq} \) and (ii) the actual velocity \( v \). While \( v_{eq} \) decides how much velocity is allowed before deceleration sets in, \( v \) gives us what ‘net’ value of velocity is achieved by actual acceleration. We have shown that if only CENBOL radiation dominates over the radiation from the Keplerian disc, \( v_{eq} \sim 1 \) is achieved within about few tens of Schwarzschild radii, but the same condition is achieved at over a thousand Schwarzschild radii for much higher values of Keplerian luminosity. From Figs. (4a-4b), we have seen as the Keplerian accretion rate is increased, \( v_{eq} \) decreases in the range \( z \rightarrow 15r_g - 1500r_g \).

From Fig. (3d), we have seen that \( (\tilde{E}_K + \tilde{P}_K) > 2\tilde{F}_K \) in this very range, so in this range the radiation drag is higher due to Keplerian radiation. This means that for higher values of \( \ell \) and for fixed \( x_s \) i.e., for higher CENBOL intensity, the flow will tend to achieve very high velocities within few tens of Schwarzschild radii, but at the same time if Keplerian luminosity or \( \ell_K \) is increased, then within the range \( 15r_g - 1500r_g \), Keplerian radiation will tend to reduce the high velocities because of increased radiation drag.

From this study we generally conclude that,

(i) radiative acceleration of electron-positron jets does achieve relativistic terminal speed.
(ii) The space dependant part of the radiative moments from the post-shock region, dominates the corresponding moments from the Keplerian disc.
(iii) In general, the terminal speed of jets increases with increasing post-shock luminosity.
(iv) Post-shock radiative intensity decreases with increasing shock location, so the terminal speed also decreases with increasing shock location.
(v) Keplerian radiation has marginal effect in determining the terminal speed.
(vi) Our calculations show that, if the shock in accretion is located at around \( 10 - 20 \) Schwarzschild radii, and if the post-shock luminosity is about 10% to 20% of the Eddington luminosity, then electron-positron jets can be accelerated to terminal speeds above 90% the speed of light.

This is supported in part by a grant from ESA Prodex project; ESA Contract no.
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14815/00/NL/SFe(IC) [IC] and Department of Science and Technology, Govt. of India, through a grant no. SP/S2/K-15/2001 [SKC and SD].

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