Energy flows in thick accretion discs and their consequences for black hole feedback

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Accessibility
Black hole feedback from thick accretion discs

Aleksander Sadowski1⋆, Jean-Pierre Lasota2,3⋆, Marek A. Abramowicz3,4⋆ and Ramesh Narayan5⋆

1 MIT Kavli Institute for Astrophysics and Space Research 77 Massachusetts Ave, Cambridge, MA 02139, USA
2 Institut d’Astrophysique de Paris, CNRS et Sorbonne Universités, UPMC Paris 06, UMR 7095, 98bis Bd Arago, 75014 Paris, France
3 Nicolaus Copernicus Astronomical Center, Bartycka 18, 00-716 Warsaw, Poland
4 Physics Department, Gothenburg University, SE-412-96 Göteborg, Sweden
5 Harvard-Smithsonian Center for Astrophysics, 60 Garden St., Cambridge, MA 02134, USA

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ABSTRACT
We study energy flows in geometrically thick accretion discs, both optically thick and thin, using general relativistic, three-dimensional simulations of black hole accretion flows. We find that for non-rotating black holes the efficiency of the total feedback from thick accretion discs is 3% - roughly half of the thin disc efficiency. This amount of energy is ultimately distributed between outflow and radiation, the latter scaling weakly with the accretion rate for supercritical accretion rates, and returned to the interstellar medium. Accretion on to rotating black holes is more efficient because of the additional extraction of rotational energy. However, the outflow and radiation components cover a wide solid angle.

Key words: accretion, accretion discs – black hole physics – relativistic processes – methods: numerical

1 INTRODUCTION
Several aspects of the evolution of galaxies have been puzzling astronomers for decades. Firstly, star formation in galaxies turns out to be efficiently quenched in galactic bulges despite the gas cooling time being much shorter than the age of a given galaxy (Cowie & Binney [1977], Fabian & Nulsen [1977]). Secondly, the galaxy luminosity function features a sharp high-mass cutoff in which the most massive systems are red, dead, and elliptical, inconsistent with the hierarchical growth of structure in the Universe (Croton et al. [2006]). Explaining both phenomena requires additional processes preventing gas from collapsing into stars and limiting the mass of the central galaxies.

Supernova explosions and stellar winds return energy (provide feedback) to the interstellar medium (ISM). Although these processes take place at small scales, they are powerful enough to affect the evolution of the whole galaxy. Without strong stellar feedback, gas inside galaxies would cool efficiently and collapse on a dynamical time resulting in star formation rates inconsistent with observations. As shown recently shown by Hopkins et al. [2014], stellar feedback itself is enough to explain most of the properties of galaxies, e.g., the relation between galaxy stellar mass and halo mass, at stellar masses $M_\star \leq 10^{11} M_\odot$.

Numerical simulations are needed to explain the formation of the most massive galaxies. It is believed that almost every galaxy harbours a supermassive black hole (SMBH) in its nucleus. Being extremely compact such objects can liberate gravitational energy in large amounts. As a black hole (BH) grows to 0.2% of the bulge mass through accreting matter, it releases nearly 100 times the gravitational binding energy of its host galaxy (Fabian et al. [2009]). It is therefore reasonable to expect that, if only the energy returned from accretion (the black hole feedback) is efficiently coupled with the ISM, the central SMBHs can strongly affect the formation and the properties of the host galaxies.

The feedback provided by SMBHs is therefore crucial for studying the evolution of the Universe. It is often accounted for in large scale simulations of galaxy formation, but the adopted models are very simplistic. The large range of scales involved in such simulations does not allow for detailed numerical (and simultaneous) modeling of the BH accretion. Instead, the mass supply rate is estimated (at most) at parsec scales, usually using the Bondi model of spherical accretion, and simple formulae for the feedback efficiency are applied. These are based partly on the standard thin disc models (Shakura & Sunyaev [1973]), but (to be consistent with observed properties of galaxies) involve additional factors arbitrarily rescaling the feedback rate.

These factors reflect our lack of understanding of how accretion on SMBHs works and the efficiency of the feedback it provides. There are two major unknown. Firstly, it is not clear how much matter makes it to the BH, and how much is lost on the way. In other words, what fraction of the gas attracted by the BH near the Bondi radius ultimately crosses the BH horizon and efficiently lib-
erates its binding energy providing the energy source for the feedback (see discussion in Yuan & Narayan 2014). Secondly, it is crucial to understand what fraction of this energy is returned to the ISM.

In this paper we address the second question. The feedback efficiency from an accretion flow is believed to be well established only for geometrically thin discs, corresponding to moderate, sub-Eddington accretion rates, $10^{-3} \lesssim M_{\text{edd}} \lesssim M \lesssim M_{\text{edd}}$ (for the definition of $M_{\text{edd}}$ see Eq. 26). In this case, the accretion flow is radiatively efficient, and all the released binding energy of the gas goes into radiation and is determined by the binding energy of the gas at the disc’s inner edge (e.g., it equals 5.7% of the accreted rest mass energy, $\dot{M}c^2$, for a non-rotating BH).

We address here the question of what amount of energy is extracted if accretion flows are not geometrically thin, i.e., how efficient the BH feedback is if a SMBH accretes either in the radio mode ($M \lesssim 10^{-3} M_{\text{edd}}$), when one expects an optically thin accretion flow and low radiative efficiency, or above the Eddington accretion rate, in an optically thick disc. To this purpose, we analyze a set of state-of-the-art, three-dimensional simulations of the innermost region of BH accretion performed with a general-relativistic, radiative magnetohydrodynamical (MHD) code KORAL (Sadowski et al. 2013).

Our paper has the following structure. In Section 2 we discuss the energy transfer in the standard model of a thin disc. In Section 3 we give the details of the numerical simulations and discuss their properties. In Section 4 we discuss their implications and several caveats. Finally, in Section 5 we summarize our findings.

## 2 ENERGY FLOW IN THIN DISCS

We start by recapitulating the physics of energy transfer in the standard model of a thin accretion disc (e.g. Shakura & Sunyaev 1973; Frank et al. 1992). This will give us a good reference point when discussing energy flows in numerical simulations of accretion flows.

### 2.1 Viscous dissipation

The thin disc model assumes Keplerian azimuthal motion, small vertical thickness of the disc, $h/r \ll 1$, and radiative efficiency. Keplerian angular velocities imply differential rotation and, in the presence of viscosity, non-zero transfer of angular momentum between adjacent rings. The torque exerted by rings on each other is (Frank et al. 1992),

$$T = 2\pi r \Sigma r^2 \frac{d\Omega}{dr} = -3\pi r \Sigma r^2 \Omega,$$

where $\Sigma$ is the surface density at radius $r$, $\Omega = \sqrt{GM/r^3}$ is the Keplerian angular velocity, and $\nu$ is the local kinematic viscosity coefficient corresponding to magnetically induced turbulence.

The torque results in transfer of angular momentum between the rings. Conservation of angular momentum requires,

$$\frac{d}{dr} M r^2 \Omega = -\frac{d T}{dr},$$

where $M > 0$ denotes the accretion rate. Integrating between radius $r$ and the inner edge of the disc at $r_m$, we get,

$$- M \left( \Omega(r) r^2 - \Omega(r_m) r_m^2 \right) = T(r) - T(r_m).$$

Following the standard assumption that the torque at the inner edge of a thin disc vanishes (see Paczynski 2000) we get,

$$T = - \sqrt{GMM} \left( r - \sqrt{r_m} \right).$$

This torque not only transports angular momentum but also dissipates mechanical energy heating up the gas. The dissipation rate (per unit radius) in the whole ring which equals, by assumption of radiative efficiency, minus the radiative cooling rate, is given by,

$$q_{\text{disc}} = -q_{\text{rad}} = -r \frac{d T}{dr} = -\frac{3 G M M}{2 r^2} \left( 1 - \frac{r_m}{r} \right).$$

where the signs have been chosen such that a positive rate corresponds to cooling, and negative to heating. Dividing by the surface area of both sides of the ring we get the well known thin-disc surface radiative flux,

$$Q_{\text{rad}} = \frac{q_{\text{rad}}}{4 \pi r} = \frac{3 G M M}{8 \pi r^3} \left( 1 - \frac{r_m}{r} \right).$$

It is worth reiterating that the viscous dissipation rate (the gas heating rate) does not depend on the particular form of the viscosity (e.g., the $\alpha$-viscosity), but follows from the assumptions of Keplerian motion, the zero-torque boundary condition and angular momentum conservation.

### 2.2 Local energy budget

In the previous paragraph we have shown that viscous dissipation in a differentially rotating flow results in heating of gas and radiative cooling at the rate $q_{\text{rad}}$ (Eq. 5).

The energy required for this radiative emission may come from the gravitational field – gas approaching the BH liberates its own binding energy at the rate,

$$q_{\text{bind}} = M \frac{d q_{\text{bind}}}{dr} = M \frac{d}{dr} \left( \frac{G M}{2r} \right) = \frac{G M M}{2 r^2}.$$  

However, it is clear that,

$$q_{\text{bind}} + q_{\text{rad}} \neq 0. \tag{8}$$

This means that there must be another source or sink component in the local budget of energy.

In the previous section we have seen that viscosity leads to the transport of angular momentum and dissipation of mechanical energy. However, viscosity transports not only angular momentum but also rotational energy. The amount of energy transported in this way is,

$$L_{\text{visc}} = -T \Omega, \tag{9}$$

and the resulting local heating or cooling rate per unit radius is given by,

$$q_{\text{visc}} = \frac{d}{dr} (T \Omega) = \frac{G M M}{2 r^2} \left( 1 - \frac{3}{2} \sqrt{\frac{r_m}{r}} \right). \tag{10}$$

It is straightforward to verify that,

$$q_{\text{bind}} + q_{\text{rad}} + q_{\text{visc}} = 0. \tag{11}$$

The viscous energy transport redistributes energy released in the disc and compensates for the imbalance between the local binding energy release and the rate of radiative cooling.
The maximum is at $r = (49/6) r_g$.

### 2.3 Energy flow

In the previous section we have looked into the local energy balance. Now, let us look into the total amount of energy carried by its various components from one radius to another.

The binding energy is carried by the flow at a rate,

$$L_{\text{bind}} = -M \dot{e}_{\text{bind}} = \frac{G M M}{2r} > 0,$$

where the positive sign reflects the fact that bound gas is falling inward, thus effectively depositing energy at infinity. The luminosity in binding energy may be again decomposed into the gravitational and kinetic components,

$$L_{\text{grav}} = -M \dot{e}_{\text{grav}} = \frac{G M M}{r} > 0,$$

$$L_{\text{kin}} = -M \dot{e}_{\text{kin}} = -\frac{G M M}{2r} < 0.$$

Gravitational energy luminosity is positive, but the kinetic luminosity is negative – kinetic energy of the Keplerian motion is brought inward by the gas.

The radiative cooling rate, $q_{\text{rad}}$, is given by Eq. [5]. Photons are emitted from the disc surface and leave the system. The total radiative luminosity at given radius, $L_{\text{rad}}$, results from the emission inside that radius,

$$L_{\text{rad}} = \int_{r_{\text{in}}}^{r} q_{\text{rad}} \, dr = \frac{3}{2} \frac{G M M}{r} \left( \frac{1}{3} \frac{r}{r_{\text{in}}} + \frac{2}{3} \sqrt{\frac{r_{\text{in}}}{r}} - 1 \right) > 0.$$

This quantity is zero at the inner edge ($r = r_{\text{in}}$) and equals $G M M / 2 r_{\text{in}}$ at infinity. The radiative luminosity of the whole accretion disc is therefore equal to the binding energy of the gas crossing the inner edge.

Finally, the amount of energy carried by viscosity from the inner region outward, $L_{\text{visc}}$, is (Eq. [9]),

$$L_{\text{visc}} = -T \Omega \approx \frac{G M M}{r} \left( 1 - \frac{r_{\text{in}}}{r} \right) > 0.$$

The various integrated energy fluxes introduced above are shown in Fig. [2]. Their magnitudes have been normalized to the amount of accreted rest-mass energy. The blue lines show the luminosities in the binding energy, $L_{\text{bind}}$, and its gravitational and kinetic components, $L_{\text{grav}}$ and $L_{\text{kin}}$. At the inner edge ($r_{\text{in}} = 6 r_g$), the amount of binding energy carried by the gas is

$$L_{\text{bind, in}} = \frac{G M M}{2 r_{\text{in}}} = \frac{1}{12} \frac{M c^2}{r_{\text{in}}}$$

which is, as we will discuss in detail in a moment, the total efficiency of a thin disc in the Newtonian gravitational potential.

The orange line in Fig. [2] shows the radiative luminosity crossing a sphere of radius $r$. As no photons are emitted from inside the inner edge, it starts from zero and gradually grows, reaching finally $G M M / 2 r_{\text{in}}$ at infinity – in a thin disc, the whole energy extracted by the infalling gas ultimately goes into radiation.

It is interesting to note that 50% of the radiation is emitted from outside $r \approx 25 r_g$. At the same time, the gas infalling from infinity down to that radius has extracted only roughly 25% of the available binding energy. The excess in radiative luminosity comes from the extra energy carried by viscosity from the innermost region.

This component of the energy flux is denoted with the pink line at the same plot. The amount of energy carried by viscosity grows rapidly just outside of the inner edge – at these radii viscosity

![Figure 1. Local energy gain in its various forms in the standard thin disc model described in Section 2.](image-url)
is transporting rotational kinetic energy outward. Outside $r \approx 13r_g$ the luminosity of viscous energy transport drops down with radius and the energy taken away from the innermost region is deposited by viscosity into the gas. Summing up all the components of the energy transfer we get the total luminosity,

$$L_{\text{tot}} = L_{\text{bind}} + L_{\text{rad}} + L_{\text{visc}},$$

which is the quantity that is fundamentally conserved in stationary flows, i.e., is independent of radius and no energy accumulates at any location. Indeed, the sum of the three components (red line in Fig. 2) gives a constant value equal to the total efficiency of accretion and the binding energy carried in by the gas through the disc inner edge $(1/12)M_c\dot{M}$. In the Schwarzschild metric this efficiency would be $\sim 0.057 M_c \dot{M}$.

3 ENERGY FLOW IN SIMULATIONS OF ACCRETION FLOWS

Having recapitulated how energy flows in a standard thin disc, we are ready to study the energy redistribution in numerical simulations of accretion flows. In the following Section we describe the numerical method used to perform the simulations. In Section 3.1 we introduce the formalism used to study energy fluxes in numerical solutions. In Sections 3.2 and 3.3 we look in detail into the energy flow in simulations of optically thin and thick discs, respectively.

3.1 Numerical setup

The simulations analyzed in this paper were performed in three dimensions with the general relativistic radiation magnetohydrodynamical (GRRMHD) code RÖRAL [Sadowski et al. 2013] which solves the conservation equations in a fixed, arbitrary spacetime using finite-difference methods. The equations we solve are,

$$\langle pu'' \rangle_\mu = 0,$$  
$$\langle T^\nu_\mu \rangle_\mu = G_\nu,$$  
$$\langle R^\nu_\mu \rangle_\mu = -G_\nu,$$

where $\rho$ is the gas density in the comoving fluid frame, $u''$ are the components of the gas four-velocity, $T^\nu_\mu$ is the MHD stress-energy tensor,

$$T^\nu_\mu = (\rho + u'' + b^2) u''^\nu u''_\mu + (p_g + b^2) \delta^\nu_\mu - b'' b_\nu,$$

$R^\nu_\mu$ is the stress-energy tensor of radiation, and $G_\nu$ is the radiative four-force describing the interaction between gas and radiation (see Sadowski et al. 2014 for a more detailed description). Here, $u_\nu$ and $p_g = (1 - 1)u_\mu$ represent the internal energy and pressure of the gas in the comoving frame and $b''$ is the magnetic field 4-vector [Gammie et al. 2003]. The magnetic pressure is $p_{\text{mag}} = b^2/2$ in geometrical units.

The magnetic field is evolved via the induction equation,

$$\partial_\mu (\sqrt{-g}B^\mu) = -\partial_\nu \left( \sqrt{-g} (b''^\nu - b'' u''^\nu) \right),$$

where $B^\mu$ is the magnetic field three-vector [Komissarov 1999], and $\sqrt{-g}$ is the metric determinant. The divergence-free criterion is enforced using the flux-constrained scheme of Tolstoy (2000).

The radiation field is evolved through its energy density and flux, and the radiation stress-energy tensor is closed by means of the M1 closure scheme [levermore 1984] [Sadowski et al. 2013]. The energy exchange between gas and radiation is by free-free emission/absorption as well as Compton scattering. The latter is treated in the “blackbody” Comptonization approximation as described in Sadowski et al. (2015a).

We use modified Kerr-Shild coordinates with the inner edge of the domain inside the BH horizon. The simulations are run with a moderately high resolution of 252 grid cells spaced logarithmically in radius, 234 grid cells in the polar angle, concentrated towards the equatorial plane, and 128 cells in azimuth.

Three of the four simulations which we analyze in this work are identical to the ones presented in [Sadowski et al. 2015d]. To have a consistent optically thin version of an accretion flow we simulated an additional model with purely magnetohydrodynamical evolution, i.e., without radiation field. This simulation (b091) corresponds to an optically thin, advection dominated accretion flows (ADAF) believed to occur in systems accreting well below the Eddington level [Yuan & Narayan 2014].

Parameters of the models are given in Table 1.

In this work we adopt the following definition for the Eddington mass accretion rate,

$$M_{\text{Edd}} = \frac{L_{\text{Edd}}}{\eta c^2},$$

where $L_{\text{Edd}} = 1.25 \times 10^{38} M/M_c \text{ergs/s}$ is the Eddington luminosity, and $\eta$ is the radiative efficiency of a thin disc around a black hole with a given spin $a_*=a/M$. For zero BH spin, $M_{\text{Edd}} = 2.48 \times 10^{38} M/M_c \text{ergs/s}$. Hereafter, we also use the gravitational radius $r_g = GM/c^2$ as the unit of length, and $r_g/c$ as the unit of time.

In this study we consider simulation output averaged over time. Therefore, whenever we write, e.g., $\langle pu'' \rangle$, we mean the average of the product, i.e., $\langle pu'' \rangle$, where $\langle \rangle$ stands for time averaging.

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Table 1. Model parameters

<table>
<thead>
<tr>
<th>Model</th>
<th>( h_{0.1} )</th>
<th>( r_{0.1} )</th>
<th>( r_{0.03} )</th>
<th>( r_{0.11} )</th>
</tr>
</thead>
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<tr>
<td>MBH</td>
<td>10M⊙</td>
<td>10M⊙</td>
<td>10M⊙</td>
<td>10M⊙</td>
</tr>
<tr>
<td>( M/M_{\text{edd}} )</td>
<td>( \leq 10^{-3} )</td>
<td>10.0</td>
<td>175.8</td>
<td>17.4</td>
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<tr>
<td>( a_\ast )</td>
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<td>0.0</td>
<td>0.0</td>
<td>0.7</td>
</tr>
<tr>
<td>( t_{\text{max}} )</td>
<td>23,000</td>
<td>20,000</td>
<td>19,000</td>
<td>16,100</td>
</tr>
</tbody>
</table>

All models initiated as in Sadowski et al. (2015d).

3.2 Energy fluxes

3.2.1 Fundamental quantities

In quasi-stationary state the accretion rate is constant in radius, i.e., gas does not accumulate anywhere, but rather flows towards the BH in the radial direction. Therefore, we define the rest-mass energy flux from the previous definition, the gravitational energy carried by gas, and by radiation \( (R'_t + \rho u'_t) \). This quantity, however, is not interesting from the point of view of a distant observer. It contains the flux of rest-mass energy which, even if deposited at infinity, will not have observational consequences (since at infinity rest-mass cannot be converted into other forms of energy in a trivial way). Therefore, we define the total luminosity by subtracting the rest-mass energy flux from the previous definition,

\[
L_{\text{tot}} = -\int_0^\infty \int_0^{2\pi} \sqrt{-\mathbf{g}} \rho u'_t d\phi d\theta, \tag{27}
\]

where we integrate the radial flux of energy carried by gas \( (T'_t + R'_t) \) and by radiation \( (R'_t) \). This quantity, however, is not interesting from the point of view of a distant observer. It contains the flux of rest-mass energy which, even if deposited at infinity, will not have observational consequences (since at infinity rest-mass cannot be converted into other forms of energy in a trivial way). Therefore, we define the total luminosity by subtracting the rest-mass energy flux from the previous definition,

\[
L_{\text{tot}} = -\int_0^\infty \int_0^{2\pi} \sqrt{-\mathbf{g}} \rho u'_t d\phi d\theta. \tag{29}
\]

The sign has been chosen in such a way that \( L_{\text{tot}} \) is negative for energy falling in the BH, and positive for energy leaving the system.

In a stationary state the total luminosity is independent of radius (if it was not, energy would accumulate in some regions). It is the luminosity of the whole system, i.e., it is also the luminosity as seen from infinity. Therefore, it determines the rate at which energy is deposited in the interstellar medium or, in other words, \( L_{\text{tot}} \) is the total power of feedback.

3.2.2 Decomposition

The total energy flux consists of multiple components. We decompose it in a way which gives well-known Newtonian limits.

First, we single out the radiative component and define the radiative luminosity,

\[
L_{\text{rad}} = -\int_0^\infty \int_0^{2\pi} \sqrt{-\mathbf{g}} R'_t d\phi d\theta, \tag{30}
\]

which reflects energy carried by photons, either trapped in the gas, or propagating freely. To define other components, let us first write explicitly,

\[
T'_t + \rho u'_t = \rho u'_t (1 + u_t) + (\Gamma u_t + b^2) u_t - b' b_t. \tag{31}
\]

Here we remind the reader that in all the integrals we take averages of products, e.g., \( \rho u'_t (1 + u_t) \) is actually \( \langle \rho u'_t (1 + u_t) \rangle \), where the product is averaged over time. This particular quantity is the average radial flux of binding energy. In detail, it is the sum of advective, \( \langle \rho u'_t (1 + u_t) \rangle \), and Reynolds (turbulent), \( \langle \rho u'_t (1 + u_t) \rangle = \langle \rho u'_t (1 + u_t) \rangle \), components. Similar decomposition applies to the other components of the total energy flux. In this work we will not discriminate between the turbulent and advective fluxes, but instead focus on the net contribution.

It is straightforward to define the luminosity in internal (thermal) energy,

\[
L_{\text{th}} = -\int_0^\infty \int_0^{2\pi} \sqrt{-\mathbf{g}} \Gamma u_t u'_t d\phi d\theta, \tag{32}
\]

which, similarly, contains the advective and convective terms, and the luminosity carried by the magnetic field,

\[
L_{\text{magn}} = -\int_0^\infty \int_0^{2\pi} \sqrt{-\mathbf{g}} (b^2 u'_t u_t - b' b_t) d\phi d\theta. \tag{33}
\]

The gravitational component of the last expression can be singled out by calculating the specific binding energy \( (1 + u_t) \) for a stationary observer. From \( u^t = (u'_t, 0) \) and \( u^r u_r = -1 \) one gets (for a diagonal metric) \( u_t = -\sqrt{1 - u^t u_t} \), and therefore, the luminosity in gravitational energy carried by gas is

\[
L_{\text{grav}} = -\int_0^\infty \int_0^{2\pi} \sqrt{-\mathbf{g}} \rho u'_t (1 + u_t) d\phi d\theta. \tag{35}
\]

The remaining term reflects the luminosity in kinetic energy,

\[
L_{\text{kin}} = L_{\text{th}} - L_{\text{grav}}. \tag{36}
\]

To sum up, we have decomposed the total energy transfer rate into binding, thermal, magnetic, and radiative components,

\[
L_{\text{tot}} = L_{\text{grav}} + L_{\text{th}} + L_{\text{magn}} + L_{\text{rad}}. \tag{37}
\]

3.2.3 Advective and viscous energy fluxes

In a viscous accretion flow energy is transported both by viscosity and by the fluid which advectively carries energy with itself. One may write, \( \dot{L}_{\text{vis}} = \dot{L}_{\text{adv}} + \dot{L}_{\text{visc}} = \dot{M} \dot{B} e - T \Omega, \) where \( \dot{B} e \) is the Bernoulli function of the fluid (which is not constant, because work is done on gas on its way towards the BH), and \( T \Omega \) reflects the viscous rate of energy flow (Eq. [18]). The hydrodynamical quantities defined in the previous section \( (T_{\text{adv}}, L_{\text{vis}}, L_{\text{magn}}) \) are based on time averaged quantities. The turbulence, which provides effective viscosity, is averaged out and contributes to the energy transfer rate. Therefore, as stated in the previous Section, these luminosities include both terms, the advective and the viscous one. It is beyond the
and global simulations of magnetized accretion show that the mag-

This should not be surprising, because it is mostly magnetic field

This quantity is plotted in Fig. 3 with the pink solid line. In the

Figure 3. Total estimated viscous flux of energy (solid line), and the energy
carried by magnetic fields, $L_{\text{magn}}$ (Eq. 33, dashed line).

We estimated the viscous component by calculating

This quantity is plotted in Fig. 3 with the pink solid line. In the

Figure 3 presents the integrated radial fluxes (luminosities) of energy
in various forms for the optically thin simulation h001.  The amount
of binding energy (Eq. 34) carried with the flow is shown with solid
blue line. The closer the gas gets to the BH, the more bound it is, and
the more luminosity it extracts with respect to infinity (once again,
infalling bound gas effectively deposits energy at infinity).  It can be
decomposed into the gravitational (Eq. 25 blue dashed line) and the
kinetic (Eq. 36 blue dotted line) components. Because the flow is only
slightly sub-Keplerian and the radial velocities involved are low, these
two components behave qualitatively in the same way as in the case
of the thin disc discussed in the previous Section.

The magnetic component (Eq. 33), which reflects the energy
carried by effective viscosity, also qualitatively agrees with the thin
disc prediction. It is zero at the inner edge (which is now at the
horizon, not at innermost stable circular orbit (ISCO), because for
thick discs stress is not zero down to the horizon), and becomes
positive, which again reflects the fact that turbulent viscosity takes
energy out of the innermost region (here from $r \lesssim 10$) and carries
it outward. In contrast to the thin disc model, however, there is no
clear decrease in the magnetic luminosity inside the convergence
region of the simulation, i.e., turbulent viscosity does not contribute
to the local heating rate.

Because radiative cooling is not efficient, energy is not trans-
ferred by radiation. The dissipated energy is trapped in the flow,
heats up the gas, and contributes to the thermal energy transport
(Eq. 32). This fact is reflected in the grey line profile in Fig. 4. The
luminosity in thermal energy is no longer negligible, as it was in
the thin disc case. Significant amount of thermal energy is carried
inward with the flow. Because gas becomes hotter when approaching
the BH horizon, the corresponding magnitude increases.

If the studied accretion flow followed the standard model, i.e.,
all the energy released was advected on to the BH, all the com-
ponents contributing to the energy transfer should sum up to zero
total efficiency. This is, however, not the case. The thick red line
in Fig. 4 reflects the total luminosity defined through Eq. 29, i.e.,
composed of binding, magnetic and thermal components. It is flat
to a good accuracy inside $r \approx 25$ proving that the flow has reached
a quasi-stationary state in this region. The efficiency of $\sim 3\% \dot{M}c^2$,
reflects the amount of energy extracted from the accretion flow
and equals roughly 50% of the thin disc efficiency.

In the ideal case of an optically thin disc extending to infinity,
this amount of energy would be deposited at infinity. In practice,
this is the amount by which the BH systems a

3.3 Energy flow in optically thin ADAFs

Let us now look at the energy flow in simulated, multi-dimensional
accretion flows. We start with an optically thin disc (ADAF), model
h001.

According to the standard model [Narayan & Yi] [1994]
[Abramowicz et al.1995] for this mode of accretion, energy locally
dissipated does not have a chance to escape because of low radia-
tive efficiency and is advected with the flow. This fact makes such
disks very hot and geometrically thick. As a result, the expected ef-
iciency of accretion is zero because all the binding energy gained
by gas on its way towards the BH is balanced by thermal energy
advected with it on the BH.

This model, however, does not allow the gas to flow out of
the system. This process, in principle, can provide a path for the
liberated binding energy to escape from the system, and as a result
may increase the efficiency of accretion.

3.3.1 Luminosities

4 If the viscous stress is proportional to shear then it is orthogonal to
the gas velocity, $T_{\text{visc}}^{\mu\nu} u^\nu = 0$. For purely azimuthal motion, $u^\varphi = (0, 0, 0, u^\varphi)$,
one finds that $T_{\varphi\varphi}^{\mu\nu} \Omega = -T^{\mu\nu}_{\varphi}$. Therefore, $T_{\varphi\varphi}^{\mu\nu} \Omega$ indeed gives the radial flux
of energy carried by viscosity. However, in the simulations we performed,
the orthogonality and perfectly circular motion are not enforced and these
conditions are only approximately satisfied.

5 This value gives the total energy of feedback, in contrast to values given
in Sadowski et al. (2013b) who gave only power in jet and wind compo-
nents.
Black hole feedback from thick accretion discs

feedback). Because simulations with inefficient radiative cooling are scale-free, this efficiency is characteristic of an optically thin flow (ADAF) at any accretion rate for which such a solution exists with negligible radiative cooling (i.e., for $\dot{M} \lesssim 10^{-3} \dot{M}$, see Yuan & Narayan 2014). The fate of the energy coming out of the innermost region is discussed below in Section 4.1.

## 3.3.2 Angular distribution

Figure 5 shows the spatial distribution of density (top-most panel) and various components of the energy flux (other panels) in the optically thin simulation $h001$. The streamlines in the top-most panel reflect the velocity of the gas. The second panel shows the corresponding rest-mass energy flux, $\rho u^\mu$. Most of the accretion takes place near the equatorial plane. Within radius $r = 30$, gas at all polar angles falls inward. Only outside this radius (and outside the converged region of the simulation), there is a hint of outflows that may arise from the accretion flow.

The third panel shows the magnitude (colors) and direction on the poloidal plane of the total energy flux, $L_{\text{tot}}$ (Eq. 29). The total energy flows outward, in agreement with the positive total efficiency of 3% (see Fig. 4). Most of the extracted energy flows into the disc, and very little in the polar region.

The fourth panel shows the (negative) binding energy brought inward with the gas. This effectively transports energy outward. This component is more isotropic than the net energy flux, reflecting the fact that gas falling in along the axis is more bound than gas accreting in the equatorial plane.

The fifth panel shows the magnetic component, which, as we argued in the previous section, corresponds roughly to the viscous energy transfer rate. As in the case of a thin disc, viscosity transports energy outward and redistributes it throughout the disc. Most of this energy goes near the equatorial plane, where the density is largest, and viscosity most efficient.

Finally, the bottom-most panel presents the flux of internal energy. Its magnitude is significant (compare Fig. 4) because optically thin flow cannot cool efficiently and becomes very hot. As expected, thermal energy is brought inward with the gas, and the angular distribution is again quasi-spherical – although the accretion rate is highest near the equatorial plane, the gas temperatures there are lower than in the polar region.

## 3.4 Energy flow in optically thick super-critical discs

Accretion flows transferring gas at rates higher than the Eddington limit are optically thick, but they are not as radiatively efficient as thin discs. The vertical optical depth is so large, that the cooling time becomes comparable or larger than the accretion time, and
The distribution of the energy fluxes on the poloidal plane for the super-critical disc accreting at $\sim 10M_{\text{Edd}}$ on a non-rotating BH (model $r001$). The colors have the same meaning as in Fig. 4.

such discs cannot cool efficiently. Instead, significant fraction of radiation may be advected on to a black hole. We now look closely at the simulation $r001$ of a mildly super-critical disc ($10M_{\text{Edd}}$) near a non-rotating BH.

3.4.1 Luminosities

Figure 4 shows the luminosities in various components of energy for the super-critical disc $r001$. It can be directly compared to Fig. 4 corresponding to an optically thin disc. The two figures show qualitatively the same behavior of the binding and magnetic components – bound gas is brought inward and effectively transports energy outward, similarly to the magnetic/viscous energy flux which takes energy liberated in the innermost region and redistributes it in the outer regions.

There is, however, a significant difference in the thermal and radiative components. In the case of an optically thin disc, the liberated energy heats up the gas and results in significant inward flux of thermal energy (grey line in Fig. 4). In the case of a radiative super-critical disc, the flux of thermal energy is negligible. Instead, the radiative component is now significant. It is negative within $r \approx 35r_g$, reflecting the fact that photons are trapped in the optically thick flow and transported with the gas to the BH. Only outside this radius, the amount of radiative energy flowing out exceeds the advected one. The fact that the thermal energy is now subdominant with respect to the radiative energy is consistent with the fact that super-critical discs are radiation-pressure dominated, and therefore, inward motion advectively carries radiative energy, not thermal energy.

Interestingly, the total luminosity of the optically thick disc is again close to $3\% M_c^2$ (thick red line in Fig. 4). This is the amount of energy returned from the system to the ISM in the case of supercritical accretion on to a non-rotating black hole. Noticeably, the magnitude of BH feedback power from a geometrically thick disc near a non-rotating BH is not sensitive to its optical depth.

3.4.2 Angular distribution

The distribution of the energy fluxes on the poloidal plane for the optically thick simulation $r001$ is shown in Fig. 7. The density distribution (top-most panel) shows much larger contrast between the equatorial plane and the polar axis than in the case of an optically thin disc. This fact results from radiative pressure exerted on the gas in the funnel – gas is accelerated vertically and escapes along the axis. This is clearly reflected in the velocity streamlines shown in that panel.

The total energy extracted from the system (third panel) looks different than in the previous case. This time the polar region is not empty of outflowing energy. The optically thin radiation escaping along the polar funnel dominates the energy budget there.

The distributions of binding and magnetic energy fluxes are similar to the optically thin case. Both transfer energy within the bulk of the disc. For the case of the magnetic energy flux (which reflects the effective viscous transport), this fact supports the conjecture that this energy will dissipate at larger radii (it would not dissipate if the magnetic energy has left the disc, e.g., along the axis).

Finally, the bottom-most panel shows the magnitude and direction of the radiative flux. As already mentioned, radiation manages to escape in the polar region. However, it is trapped in the optically thick flow near the equatorial plane.

3.5 Higher accretion rates

Here, we briefly discuss how the picture described in the previous Section changes when the accretion rate increases but the BH spin remains zero. Detailed comparison of models $r001$ (accreting at $10M_{\text{Edd}}$) and $r003$ ($176M_{\text{Edd}}$) was given in Sadowski et al. (2015d).

The most important points are as follows.

The total efficiency in both cases equals approximately 3%. However, because of the larger optical depth in model $r003$, photon trapping is more effective. In particular, the polar region becomes optically thick. Inside $r \approx 30$ the gas is dragged onto to the BH even along the axis. As a result, radiative luminosity of the system goes down. Fig. 8 shows the fractional contribution of the radiative luminosity, $L_{\text{rad}}$, to the total luminosity $L_{\text{tot}}$. Blue and orange lines correspond to simulations $r001$ and $r003$, respectively. It is clear that the latter is less radiatively luminous, and that the effective trapping radius moves outward. This fact, however, turns out not to change the total efficiency.

3.6 Rotating black hole

So far we have been discussing accretion flows around non-rotating BHs. In this Section we briefly discuss what impact non-zero BH spin has on the energy flow properties.

BH spin affects accretion flows in two ways. Firstly, BH rotation modifies the spacetime geometry and for a given BH mass allows for circular orbits getting closer to the horizon with increasing BH spin. This fact results in an increased efficiency of accretion – the closer is the inner edge of the disc, the more binding energy is liberated. Secondly, BH rotational energy can be extracted in the Blandford-Znajek process (Blandford & Znajek 1977). The power of the related jet depends on the value of the BH spin and on the amount of the magnetic flux that has been accumulated at the horizon. The latter is determined by the geometry of the magnetic field in the accreted gas and the efficiency of magnetic field dragging.

It is known not to exceed the value characteristic for the magnetically arrested (MAD) state (igumenshchev et al. 2003; Narayan et al. 2003; Tchekhovskoy et al. 2011).

In this paper we analyze one simulation ($r011$) of super-critical accretion on a mildly-rotating (the non-dimensional spin...
Black hole feedback from thick accretion discs

Figure 7. Similar to Fig. 5 but for a super-critical, optically thick disc (model r001). The bottom-most panel now shows the flux of energy carried by radiation. The flux of thermal energy is negligible.

The average accretion rate in this run ($17\dot{M}_{\text{Edd}}$) is comparable to the fiducial simulation r001, and allows for direct comparison. The amount of the magnetic flux accumulated at the BH (the magnetic flux parameter $\Phi \approx 15$) is far from the MAD limit of $\Phi \approx 50$, but is large enough to study the impact of the extracted rotational energy. Optically thick super-critical accretion flows in the MAD limit were studied in a recent work by McKinney et al. (2015).

The total efficiency of simulation r011 is roughly 8%, significantly higher than that of the comparable simulation on a non-rotating BH (3%). The increase in efficiency comes from both factors mentioned above (modified spacetime geometry and the extraction of the rotational energy). The latter by itself should extract $\sim 6\% \dot{M}c^2$ for the accumulated amount of magnetic flux and spin $a_* = 0.7$, but decomposition of the total energy into these two components is not straightforward.

In Fig. 9 we show the distribution of energy flux components on the poloidal plane. The panels have the same meaning as in the previously discussed Fig. 7. There are a couple of noticeable differences between the two. Most importantly, the amount of total energy extracted into the funnel region is much higher for the rotating BH case. This is expected, because the energy extracted in the Blandford-Znajek is known to go roughly along the axis (Penna et al. 2013c). In the case of optically thin accretion on to a rotating BH (e.g., Tchekhovskoy & McKinney 2012; Sadowski et al. 2013b), the jet power is extracted as magnetic Poynting flux gradually converting (if mass loading is significant) into kinetic energy of gas. In the case of the radiative flow studied here, this extra energy is carried mostly by radiation already for $r \gtrsim 3r_g$. Magnetic flux is significant only in a shell surrounding the funnel region. Fig. 10 shows the magnitude and direction of the radiative flux in the immediate vicinity of the BH. As expected, radiative flux falls on the BH in this innermost region, and it is the magnetic energy which is extracted at the horizon. However, the latter is quickly converted into the radiative energy. This is possible because the magnetic field efficiently pushes hot and optically thick gas along the axis. The gas, in turn, drags the radiation upward.

At the risk of oversimplifying, it is possible to say that the properties of an energy flow in the case of a super-critical accretion on to a rotating BH are a superposition of the disk component (discussed in Section 3.4 for a non-rotating BH) and the jet contribution coming from the Blandford-Znajek process. The power of the latter depends on the BH spin and magnetic flux threading the horizon, and may overwhelm the former in magnitude. At the same time, the jet component is limited only to the polar region. If the confinement provided by the disc is strong enough, it is likely to stay collimated.
4 DISCUSSION

4.1 The fate of the energy flow

We have shown so far that in geometrically thick discs, both optically thin and thick, a significant flux of energy is liberated in the accretion flow and flows out of the system. Although the simulations we performed allowed us to study only the innermost \( r \lesssim 25 \) at the equatorial plane) region of the flow, we were able to infer the total luminosity of the system. This is because in a stationary state this quantity is determined by the energy flux crossing the BH horizon.

However, the adopted method does not allow to study what happens to the extracted energy outside the inflow/outflow equilibrium region (i.e., for \( r \gtrsim 25 \)) – only for gas inside this region the duration of the simulations was longer than the viscous time scale.

As Figs. 4 and 6 show, the total luminosity of the flows around a non-rotating BH (roughly \( 3\% M_c^2 \)) comprises three components at radius \( r = 25 \). The largest in the magnitude is the binding energy flux which effectively deposits energy at infinity. The magnetic component (reflecting the viscous energy transport) is also transporting energy outward in a significant amount. The remainder goes either into the thermal (for optically thin case) or radiative energy (for optically thick) component. In both cases their net effect results in advecting energy inward.

The binding energy consists of the gravitational and kinetic components (plotted with dashed and dotted blue lines, respectively, in Figs. 4 and 6). The former goes to zero with increasing radius, and at infinity no gravitational energy is transported by the gas. The kinetic component is negative inside the computational domain reflecting the fact that gas flows inward and carries kinetic energy of its rotational motion. However, when outflows are efficiently generated it might become ultimately positive outside the computational domain.

As we have shown above, the radial flux of magnetic energy reflects the effective viscous energy transfer. Viscosity not only transports angular momentum and energy, but also leads to dissipation of the latter. Therefore, one may expect that the amount of energy carried by magnetic fields will dissipate sooner or later outside the convergence region of the simulation, adding up to the local heating rate in the same way as for thin discs discussed in Section 2. Thus, no magnetic energy will be ultimately directly deposited in the ISM.

The radiative energy transfer is important only for optically thick accretion flows. Simulations of such flows described in this paper (models \( r001 \) and \( r003 \)) show significant photon trapping in the bulk of the disc which results in negative net flow of radiative energy (see Figs. 6 and 8) in the inner region. However, as radiation gradually diffuses out from the disc, the outflowing component finally overcomes inward advection, and the net radiative luminosity becomes positive. In particular, from the point of view of an observer at infinity, radiation will only carry energy outward. In the following Section we discuss how bright the accretion flow can be.

The thermal energy flux, which contributes significantly to the energy transfer rate in optically thick accretion flows, reflects both the advective and convective contributions. In the inner part of the flow it is the advective component which dominates and results in negative net energy transfer rate – hot gas is accreted and takes its thermal energy with it. However, if outflows are present in the outer region, the net effect may be opposite and the thermal energy carried advectively outward with the outflowing gas may dominate.

Convection can carry energy against the gravity without transporting mass. This component is negligible in the simulations we
performed (thermal energy flows advectively inward), but in principle it may become significant, or even dominant, in the outer region. Several models of accretion flows which are dominated by convection (convectively dominated accretion flows, CDAFs) have been formulated (Quataert & Gruzinov 2000; Narayan et al. 2000; 2002; Abramowicz et al. 2002). Whether convection is important is an open question. Narayan et al. (2012) performed a set of simulations similar to our model h601 and showed that optically thin flows are convectively stable within \( r \leq 100 \). On the other hand, Yuan et al. (2015), who studied non-magnetized viscous flows on large scales, found that the accretion flow is actually convectively unstable. Even if convection is important, it cannot transport energy beyond the outer edge of the disc (or beyond the Bondi radius). One may expect that ultimately all energy transported by convection is released as radiation or generates outflow.

To sum up, geometrically thick accretion flows on to a non-rotating BHs deposit in the ISM roughly 3% of the rest-mass energy crossing the BH horizon. This energy may be transported outward with the outflowing gas, radiation and convection. Which components dominate is currently unclear. Ultimately, however, only radiation and outflow can transport energy beyond the Bondi radius and deposit it in the ISM. A separate jet component may be present in case of a spinning BH which managed to accumulate significant magnetic flux at its horizon. It will result in a collimated, narrow outflow of mostly kinetic energy, unlikely to interact efficiently with the ISM.

4.2 Radiative luminosity
Radiation is one of the ways of extracting energy from accretion flows. The total luminosity (accounting for all forms of energy) for a thick disc near a non-rotating BH seems to be robust – every simulation we have performed indicates that roughly 3% of the accreted rest-mass energy is returned to ISM. The amount of this energy that goes into radiation is not, however, easy to estimate. Only when the photosphere is properly resolved, one can check if radiation reaches the observer at infinity. This is not the case for any of the optically thick simulations discussed here. They were run only for a time which allowed them to reach inflow/outflow equilibrium state at the equatorial plane within \( r \leq 25 \). Because of large optical depths, photospheres are located at large distances (Sadowski et al. 2015), significantly outside the converged region. This fact makes it impossible to directly measure the amount of radiation escaping the system. Only radiation escaping along optically thin funnel, if it exists, is guaranteed to reach a distant observer.

Because of significant photon trapping in the super-Eddington region, the radiative luminosity of the system is not proportional to the accretion rate. What is more, the radiation coming from such an accretion flow must penetrate the optically thick wind region. It cannot be therefore locally significantly super-Eddington, because in such a case it would transfer its energy and momentum to the gas accelerating it. We are inclined to suggest that the result will be similar to the effect of pure photon trapping which results in logarithmic dependence of the luminosity on the accretion rate (already anticipated by Shakura & Sunyaev 1973, see also Begelman 1979).

\[
    L_{\text{rad}} \approx L_{\text{Edd}} \left( 1 + \log \frac{M}{M_{\text{Edd}}} \right).
\]  \hspace{1cm} (40)

A similar logarithmic behavior was found in the old works on super-Eddington Polish Doughnuts and explained by Paczynski and collaborators as being a consequence of the drop in efficiency when, with increasing accretion rate, the inner edge of the accretion disc moves from the ISCO to the innermost bound circular orbit (IBCO), where the efficiency is zero (see Wielgus et al. 2015 for explanation and references).

4.3 Outflow
Our work shows that the existence of outflows is inevitable in outer parts of thick accretion discs. Avoiding them requires all of the extracted energy to be ultimately transported outwards by radiation. However, thick discs are radiatively inefficient (see Eq. 44).

This conclusion is based on disc energetics – a significant fraction of the accreted rest mass energy flows outward through the disc which cannot generate enough radiation to provide the efficient cooling required to get rid of the energy surplus. If convection is not effective, at least in the outermost region, outflow is the only possible way of taking this excess of energy out of the system.

We do not see strong outflows in the simulations we performed (compare topmost panels in Figs. 5, 7, and 9). Only in the funnel region of optically thick simulations r603 and r911 one observes that the radiative luminosity is converted gradually into kinetic energy of the outflowing gas. However, the kinetic luminosity of such gas measured at the outer boundary is still at most ∼ 10% of the total efficiency.

Therefore, one may expect that most of the outflow will be generated at radii larger than covered by the inflow/outflow equilibrium region of the simulations, i.e., at \( r \geq 25r_{\text{g}} \). What will drive these outflows? In principle, there are three acceleration mechanisms likely to act in magnetized accretion flows – magnetocentrifugal (Blandford & Payne 1982), radiative and thermal. Magnetocentrifugal driving is not effective in the simulated inner region of a non-MAD accretion flow (see also Möller & Sadowski 2015), and there is little hope for it to become effective further out, where magnetic field has no reason to be more uniform on large scales. Radiative driving is seen in the funnel region of the simulated super-critical discs, but does not result in significant outflow at larger polar angles – radiation diffusing out of the disc into the optically thick wind region is supporting the disc against gravity, and therefore cannot on average significantly exceed the local Eddington flux.

Thermal wind driving remains the only candidate to balance the energy budget of the discs. It is especially reasonable if we consider the energy flux redistributed through viscosity. When it finally dissipates at larger radii, it will heat up the gas and make it more prone to become unbound and likely to flow out of the system. This is in agreement with the standard ADAF model which predicts positive Bernoulli function for the inflowing gas in the self-similar regime (Narayan & Yi 1994). At the same time, most of the observed outflows in BH accretion flows are believed to be of thermal nature (e.g., Lee et al. 2002; Ponti et al. 2012; Nielsen 2013).

The total luminosity will be ultimately carried by the outflow and radiation. Thus, at infinity, the outflow will carry the amount of energy equal to the difference between the total and radiative luminosities. For a non-spinning BH one will have,

\[
    L_{\text{outflow,∞}} = 0.03Mc^2 - L_{\text{rad,∞}}.
\]  \hspace{1cm} (41)

The latter term is obviously negligible for optically thin discs, for which the whole extracted energy goes into outflow.
4.4 Extent of a thick disc

In the considerations so far we assumed that the accretion flow extends to infinity, and that the gas at infinity has negligible energy, i.e., its Bernoulli number is zero. These conditions do not have to be satisfied in reality. In particular, a thick disc is expected to become thin in the outer region. For example, a super-critical disc becomes radiatively efficient at a sufficiently large radius where its optical depth is no longer large enough to prevent locally generated radiation from diffusing out in an efficient way. Somewhat similarly, optically thin discs (ADAFs) cannot exist above a critical accretion rate, and this limit decreases with radius. Therefore, for a fixed accretion rate at a BH, there is a radius where thick disc must become thin, and radiatively efficient. In such cases, the picture presented here would have to be modified accordingly, e.g., if at the transition radius the outward energy flux inside the disc (i.e., not in the outflow) is still significant, then one could expect that the thin disc, extending from this radius outward, will ultimately release all its energy as (relatively cold) radiation.

4.5 Transition between accretion modes

We did not simulate thin, sub-Eddington accretion flows, for which the standard assumption of being radiatively efficient is satisfied by construction. For such discs, the energy transfer is expected to follow the characteristics described in Section 2, i.e., the total efficiency of BH feedback equals the thin disc efficiency, and ultimately all of it is carried by radiation which is emitted over a wide solid angle. Thin discs are unlikely to drag significant amount of magnetic field on to the BH (e.g., Lubow et al. 1994, Ghosh & Abramowicz 1997, Spruit & Uzdensky 2005, Rothstein & Lovelace 2008, Beckwith et al. 2009), and therefore one does not expect strong Blandford-Znajek jet component.

The transition between optically thick but geometrically thin and thick discs takes place near the Eddington accretion rate. In the past it has been modeled with the so called slim disc model (e.g., Abramowicz et al. 1988, Sadowski 2011) which generalizes the standard thin disc model to higher accretion rates. Recently, numerical simulations similar to the ones that this work is based on, have studied a number of super-Eddington accretion flows (e.g., Sadowski et al. 2015a, Jiang et al. 2014b). Simulations of thin discs, more demanding computationally, have not yet been performed. Below, we will describe the transition between geometrically thin and thick optically thick discs with the help of arbitrary step functions, which make the final formulae agree qualitatively with what we have learned from numerical, multi-dimensional simulations.

The transition between optically thin and thick discs is even less well understood and awaits numerical modeling. It is known than radiatively inefficient optically thin discs flows cannot exist above some critical accretion rate $M_{\text{ADAF}} \approx 10^{-3}M_{\text{Edd}}$ (e.g., Esin et al. 1997). Whether increasing the accretion rate above this threshold results in a dramatic transition to a cold, optically thick disc, or rather the disc takes a form similar to the luminous-hot accretion flow (LHAF, Yuan 2001), has still to be verified.

Below, for simplicity, we assume that whenever accretion rate is below $M_{\text{ADAF}}$, accretion occurs in optically thin disc, and that the transition to optically thick discs (for $M > M_{\text{ADAF}}$) takes place instantaneously.

Having these considerations in mind, one may approximate the total amount of feedback luminosity coming from an accreting system as,

$$ L_{fb} = \frac{1}{2} \eta_{\text{thin}} M c^2 + P_{\text{BZ}}, \quad (42) $$

for $M < M_{\text{ADAF}}$ (opt. thin),

$$ L_{fb} = \eta_{\text{thin}} \left( 1 - \frac{1}{2} f_t \right) M c^2 + f_{\text{BZ}} P_{\text{BZ}}, \quad (43) $$

for $M > M_{\text{ADAF}}$ (opt. thick),

where $\eta_{\text{thin}}$ stands for the efficiency of a standard thin disc with given spin, the 1/2 factor reflects two times smaller efficiency of thick discs, and where we allow for the Blandford-Znajek contribution, $P_{\text{BZ}}$, for thick discs. $M_{\text{ADAF}} \approx 10^{-3}M_{\text{Edd}}$ is the critical accretion rate above which radiatively inefficient optically thin accretion flows do not exist. Functions $f_t$ and $f_{\text{BZ}}$,

$$ f_t = 1 + \left( \frac{3}{M/M_{\text{Edd}}} \right)^{1/3}, \quad (44) $$

$$ f_{\text{BZ}} = 1 + \left( \frac{1}{M/M_{\text{Edd}}} \right)^{1/3}, \quad (45) $$

were chosen to give (arbitrary) smooth transitions between the sub- and super-Eddington regime for the efficiency and the jet power, respectively. The Blandford-Znajek term (given here for saturated magnetic field at the BH, i.e., for the MAD limit, see Tchekhovskoy 2015).

$$ P_{\text{BZ}} = 1.3 \eta^2_{\text{BZ}} M c^2, \quad (46) $$

is strongly damped for thin discs which are not likely to drag the magnetic fields effectively.

Fig. 11 shows the disc and jet components of the total black hole feedback, $L_{fb}$, (Eqs. 42 & 43) as a function of accretion rate for BH spins $a_*=0.0$ and $0.7$. For the latter, we assumed that magnetic field saturated at the BH at half of the MAD limit (in this way the jet power was not overwhelming the power of the disc component).

The solid lines show the efficiency of the feedback coming from the disc. In the thin disc regime ($M_{\text{ADAF}} \leq M \leq M_{\text{Edd}}$), this efficiency equals the standard thin disc efficiency - $\eta = 0.057$ and 0.104 for $a_*=0.0$ and 0.7, respectively. For accretion rates significantly exceeding the Eddington limit, this efficiency drops down to roughly half of the thin disc efficiency, i.e., to $\eta = 0.03$ for a non-rotating BH. The proposed formulæ make the transition between the two regimes smooth. In the limit of low accretion rates $M < M_{\text{ADAF}}$, one expect accretion flows to be optically thin with similar efficiency of $\eta = 0.03$. The transition to the thin disc limit is probably more violent, and we did not apply any smoothing function there.

The dashed lines reflect the power of the jet feedback component. It is non-zero only for the case of a rotating BH. For geometrically thick discs jet production is efficient and given for magnetic flux at the BH saturated at half of the MAD limit) by $1/4 P_{\text{BZ}}$ (see Eq. 46). Thin discs are unlikely to provide strong jet components and therefore we damped the jet power in this regime. One has to keep in mind that the jet component will be highly collimated and may not interact efficiently with the ISM.

5 SUMMARY

In this paper we have studied the flow of energy in geometrically thick discs, both optically thin and thick. We based our study on
a set of state-of-the-art, three-dimensional simulations of accretion flows performed in the framework of general relativity. Our results are as follows:

(i) **Total feedback**: Thick accretion flows on a non-rotating BH show the same total efficiency of $3\% M c^2$ (roughly 50% of the thin disc efficiency) independent of the accretion rate. Both optically thin, ADAF-like flows, and super-Eddington optically thick flows liberate energy at the same rate. This energy is ultimately distributed between the energy carried by outflow and radiation. The efficiency of accretion flows onto a rotating BH is increased by the modified spacetime geometry and the rate at which BH rotational energy is extracted through the Blandford-Znajek process.

(ii) **Approximated formulae**: One may approximate the total amount of feedback coming from an accreting system using Eqs. \[ (42) \] and \[ (43) \]. These formulae assume that the total efficiency of the feedback disc component is proportional to the amount of energy extracted from the disc itself, and not from the jet, equals half of the thin disc efficiency, i.e., $\epsilon_{\text{disc}} \approx \epsilon_{\text{jet}} = 0.5 \epsilon_{\text{thin}}$.

(iii) **Energy in the outflow**: The energy outflowing from the system can be ultimately carried away only by radiation or outflowing gas. If a disc cannot cool efficiently, i.e., if it is not luminous (in radiation), then most of the liberated energy must be carried by the outflow. This is true both for optically thin discs and for optically thick discs which at sufficiently high accretion rates efficiently trap radiation in the gas. Therefore, we can infer the existence of outflows even if they are not emerging strongly within the computational domain.

The amount of energy, either kinetic or thermal, carried by the outflow equals,

$$L_{\text{outflow}} = L_{\text{th}} = L_{\text{rad}}.$$  

where the radiative luminosity is zero for optically thin discs and may be approximated as,

$$L_{\text{rad}} \approx L_{\text{Edd}} \left(1 + \log \frac{M}{M_{\text{Edd}}} \right).$$  

for super-critical discs.

(iv) **Outflowing mass**: Our study is based on simulations covering only the innermost region of BH accretion. We find that significant amount of energy flows out from that region and likely results in outflowing mass from larger radii. However, despite the fact that we know how energetic the outflow can be, we are not able to say in what amount gas is blown away. The relation between the two depends on the Bernoulli function of the outflowing gas, e.g., marginally bound gas will carry virtually zero energy per unit mass. Because of similar reasons we cannot determine the amount of momentum carried with the outflow. As [Begelman (2012)](2012) points out, accretion through thick, advective discs leads to either winds or breezes. To find how much gas is lost on the way towards the BH, one has to solve the problem consistently on larger scales than covered by the simulations presented here. Recently, a significant progress in this direction has been made by [Yuan et al. (2015)](2015) and [Bu et al. (2015)](2015), who studied optically thin accretion flows and found that gas is likely lost between $r = 40$ and the Bondi radius, and that the mass loss rate in the wind increases proportionally to radius according to $\dot{M}_{\text{out}} = \dot{M}_{\text{BH}} (r/40)$.

(v) **Angular distribution of feedback**: In the case of optically thin accretion, the liberated energy can flow out in two channels. The jet component related to the extraction of BH rotational energy is collimated along the axis and ultimately results in a narrow, relativistic magnetized jet. The accretion component flows outward in the bulk of the disc and is responsible for driving the outflows at large radii or ultimately leaves the system in convective eddies. Such energy flows will have a quasi-spherical distribution in space and will likely interact efficiently with ISM.

To some extent similar properties characterize outflows in case of super-critical, optically thick discs. The jet component is likely to be collimated along the axis, while the outflow component covers wide range of angles. The radiation coming out of the system may have initially a mildly collimated component in the funnel region (the radiative jet, see [Sadowski et al. (2015b)](2015b)). However, it either converts into kinetic jet (if there is enough coupling between radiation and gas in the funnel), or ultimately diffuses when the funnel opens because of only mild collimation of the photon beams ([Jiang et al. (2014b)](2014b) [Narayan et al. (2015)](2015)). Therefore, the radiation component should be expected to cover large solid angle from the point of view of a distant observer. Thin accretion discs, which we did not study here, are expected to produce largely isotropic radiative feedback.

(vi) **Models of thick discs**: Our study shows that outflows in some form or convection is inevitable for thick discs. This is not surprising because advection dominated accretion involve fluid which is only weakly bound to the BH ([Narayan & Yi (1994)](1994) [Blandford & Begelman (1999)](1999)). Existence of outflows or convection in principle rules out well known and celebrated models of thick accretion flows which assume that gas is not lost on the way towards the BH, and which do not allow for convection, i.e., optically thick slim discs ([Abramowicz et al. (1988)](1988) and optically thin ADAFs [Narayan & Yi (1994) [Abramowicz et al. (1995)](1995)]. However, the outflow and convective regions do not extend all the way down to the BH. Therefore, the innermost region can be described with the use of these models.

In the outer region the situation is less clear because all the proposed semi-analytic models for convection and winds suffer from some problems, or they have been developed, as the recent inflow-outflow solution by [Begelman (2012)](2012), in application for the inner part of the flow. In particular, the models of convectively dominated discs ([Quataert & Gruzinov (2000)](2000) [Narayan et al. (2000) [2002] [Abramowicz et al. (2002)](2002) in optically thin flows are self-similar. The [Dotan & Shaviv (2011)](2011) model for slim discs with winds uses sophisticated descriptions of the disc and the wind separately, but assumes an ad hoc wind launching mechanism. Finally, simple and widely used ADIOS model ([Blandford & Begelman (1999)](1999) takes strong assumptions which have been criticized by [Abramowicz et al. (2000)](2000)).

![Figure 11. Total efficiency of the feedback.](image)
6 ACKNOWLEDGEMENTS

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