Ratio of kinetic luminosity of the jet to bolometric luminosity of the disk at the “cold” accretion onto a supermassive black hole

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Abstract

Disk accretion onto black holes is accompanied by collimated outflows (jets). In active galactic nuclei (AGN), the kinetic energy flux of the jet may exceed the bolometric luminosity of the disk a few orders of magnitude. This phenomena can be explained in frameworks of so called “cold” disk accretion when the only source of energy of AGN is the energy released at the accretion. The radiation from the disk is suppressed because the wind from the disk carries out almost all the angular momentum and the gravitational energy of the accreted material. In this paper, we calculate an unavoidable radiation from the “cold” disk and the ratio of the kinetic energy luminosity of the outflow to the bolometric luminosity of the accretion disk around a super massive black hole in framework of the paradigm of the optically thick $\alpha$-disk of Shakura & Sunyaev. The ratio of the luminosities is defined predominantly by the ratio of the magnetic field pressure inside the disk to the magnetic field pressure at the base of the wind. The obtained equations applied to the jet of M87 demonstrate good agreement with observations. In the case of Sgr A*, these equations allow us to predict the kinetic energy flux from the disk around Galactic SMBH.

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INTRODUCTION

In the standard theory of disk accretion onto a black hole [1] is assumed that the gravitational energy of the accreted material is entirely carried out by photons radiated from the surface of the optically thick disk. Measurements of the bolometric luminosity of the disk gives total energy released at the accretion. However, the disk accretion can be accompanied by formation of collimated outflows (jets). Moreover, the observations of recent years show that the power of kinetic energy of the jets in some AGN exceeds the bolometric luminosity of the disk.

The famous galaxy M87 is a striking example of an AGN with a very large kinetic-to-bolometric luminosity. According to [2], the bolometric luminosity of the disk does not exceed $10^{42}$ erg/s while the kinetic luminosity of the jets is of the order of $10^{44}$ erg/s [3, 4].

Starting with the paper [5], the energetics and lifetimes of extended double radio sources have been used for calculation of the jet power in radio galaxies and quasars. The ratio of kinetic-to-bolometric luminosity can be estimated from radio and X-ray data. [6–12] argued that the radio and X-ray luminosities are likely to be related to the kinetic and bolometric luminosities, respectively. It is established that in large fraction of AGNs the jet kinetic luminosity exceeds the bolometric luminosity [13–19].

The jet power in 191 quasars detected by the Fermi Large Area Telescope (LAT) in gamma rays, calculated within the frameworks of an one-zone model show that in this sample of blazars $L_{\text{kin}}$ systematically exceeds the bolometric luminosity [20].

An indirect evidence of high kinetic luminosity of an outflow exceeding the bolometric luminosity is provided by observations of the Galactic Center in TeV gamma-rays [21]. To explain the observed diffuse flux of the VHE gamma-rays from the Galactic Center region, the production rate of protons accelerated up to 1 PeV, should be $\sim 10^{38}$ erg/s. Assuming that the accelerator of protons is powered by the kinetic energy of the outflow (a wind or jets) from the super massive Black Hole in the the Galactic Center (Sgr A*), even in the case of 100 % conversion of the bulk kinetic energy to non thermal particles, the kinetic luminosity of the outflow would two orders of magnitude exceed the bolometric luminosity of Sgr A* which is estimated close to $10^{36}$ erg/s [22].

Observations of the very powerful and bright in gamma-rays AGN 3C 454.3 give even more interesting information. During the outbursts of this object, its apparent luminosity in
GeV gamma-rays could exceed $10^{50}$ erg/s \[23\] \[26\]. The mass of the black hole in this AGN is estimated in the region $(0.5 - 4) \cdot 10^9$ M$\odot$. Thus the Eddington luminosity is in the range of $(0.6 - 5) \cdot 10^{47}$ erg/s. Because of the Doppler boosting effect, the intrinsic gamma-ray luminosity of this source is much smaller, by several orders of magnitude, than the apparent luminosity. Yet, the estimates of the jet kinetic luminosity in any realistic scenario give a value exceeding the Eddington luminosity \[27\].

In general, the estimates of the bolometric and kinetic luminosities are model dependent \[28\]. Nevertheless, it is difficult to avoid a conclusion that at least some AGN demonstrate extremely high kinetic luminosities of jets which are not only above the bolometric luminosity, but in some cases can exceed the Eddington luminosity of the central super massive black hole (SMBH). In this case the question about the source of the energy of the jets becomes the central one for understanding of the processes at the accretion onto SMBH \[29\].

Presently, the rotational energy of a black hole is considered as the most likely source of energy which is transformed in to the energy of jets due to the so-called mechanism of Blandford and Znajek \[30\]. In this scenario, SMBH provides an additional (to the accretion) source of energy which, in fact, could be the dominant source of energy of the system. According to numerical simulations \[31\] \[32\] this mechanism can provide the energy flux in the jet $\approx 3\dot{M}c^2$, where $\dot{M}$ is the accretion rate. But the black hole has to rotate with maximal possible angular momentum. In this case the radiation luminosity of the optically thick disk can achieve $\eta\dot{M}c^2$ with $\eta = 0.3$. Thus, the maximal ratio of the kinetic-to-bolometric luminosities is close to 10. This value can be increased if the accretion occurs in radiatively inefficient regime \[33\] \[35\]. In this regime the luminosity of the disk is suppressed because the energy released at the accretion is predominantly advected in to the black hole. But the alternative model is also possible in which the only source of the energy of AGN is the energy released at the accretion.

It was firstly pointed out in \[36\] that the magnetized wind from the disk can carry out a significant fraction of the angular momentum of the accreted matter rather than the viscous stresses. This splendid idea has been explored in a range of papers by Grenoble group \[37\] \[39\] which called this type of flow around black holes as Magnetized Accretion-Ejection structures (MAES). A range of other authors last years explored similar approach \[40\] \[42\] in different physical context. In the last works of the Grenoble group \[43\] \[44\] the radiation from
the Jet Emitting Disks (JED) is discussed. But the discussion is limited by the case when the jet carries out a moderate fraction (20%-80%) of the energy released at the accretion.

Starting with our first work [45] devoted to this problem we stressed that the idea of [36] can be applied to the solution of the problem of high ratio of the kinetic luminosity of the jets to the bolometric luminosity of the disks. Actually [38, 46] specially stressed much before us that at the conventional conditions the jet can carry out almost all the angular momentum and energy from the accretion disk. The radiation of the disk can carry out only a minor fraction of the accretion energy. Therefore, the ratio of the kinetic luminosity of the jets on a few orders of magnitude can exceed the bolometric luminosity of the disks in this regime of accretion. We called this type of accretion as "dissipationless" or "cold" accretion because the dissipative processes like turbulent viscosity play no role in the dynamics of the accretion. Two questions arise in this regard:

- Do solutions describing the regime of "cold" accretion exist?
- If yes then what is the expected ratio of the kinetic-to-bolometric luminosity? Is it really possible to have very high ratio of the kinetic-to bolometric luminosity exceeding 100?

Presently, we have a positive answer to the first question. The selfconsistent solution of the problem of accretion and outflow has been obtained in [38] in the self similar approximation. This allowed to the authors to calculate the structure of the accretion disk assuming that the magnetic field is formed due to advection and diffusion [39, 47].

In the work by [45] a self similar solution of the problem of "cold" accretion has been obtained neglecting the structure of the accretion disk. From one side this simplified the model but from other side this approach allows us to develop a method of numerical solution of the problem without self similarity assumptions and obtain more general numerical solutions [48, 49]. All these works remain no doubts that the "cold" accretion due to magnetized wind is really possible.

In this paper we try to answer to the second question. "Cold" or Jet Emitting Disks [46] can formally exist without any dissipation processes (viscosity and finite electric conductivity) and therefore formally can remain cold. But in reality the heating of the disk is unavoidable because the "cold" disk accretion is possible only in presence of magnetic field at the base of the wind. Therefore, the magnetic field has to be present also inside the disk.
producing magnetic and turbulent viscosity. The objective of the work is to estimate the heating of the disk due to turbulence and determine the expected level of the ratio of the kinetic-to-bolometric luminosity of the "cold" disks.

**PHYSICS OF THE "COLD" ACCRETION**

In the work of the Grenoble group the magnetic field in the disk and outflow has formed under advection and diffusion of the external regular field [50, 51]. The magnetic field lines predominantly vertically crosses the plane of the disk. But this is apparently strongly simplified case. Following to [1], we assume that the magnetic field of the accretion disk is generated by the turbulent motion of plasma due to a dynamo mechanism [52–58]. The formation of regular vertical component of the magnetic field occurs due to emerging, reconnection [59] and opening of magnetic loops by the disk differential rotation [60]. The magnetic field structure inside the disk is shown schematically in fig.1. The wind is formed quite high, in the corona where all the field lines are open, and the magnetic pressure $B^2/8\pi$ exceeds the thermal pressure $p$. The outflow of plasma from the disk occurs due to centrifugal [61] mechanism.

The equation of conservation of the angular momentum can be obtained directly from MHD equations integrated over the central volume shown in the form of rectangular box in fig.1 (for technical details see [45]). This equation in framework of the $\alpha$-disks paradigm [1] has the following form:

$$
\dot{M} \frac{\partial r V_k}{\partial r} + \frac{\partial}{\partial r} 4\pi r^2 t_{r\varphi} h + r^2 < B_\varphi B_z > |_{\text{wind}} = 0,
$$

where $t_{r\varphi}$ is the tangential stress, $V_k$ is the Keplerian velocity of rotation, $\dot{M}$ is the accretion rate, $r$ is the cylindrical radius and $h$ is the semi-height of the disk. $B_\varphi$ and $B_z$ are toroidal and z-components of the magnetic field at the base of the wind. The brackets $<>$ mean averaging over time. The term $\frac{\partial}{\partial r} 4\pi r^2 t_{r\varphi} h$ is responsible for the viscous loss of the angular momentum by the disk. The last term in the left hand side of the equation is responsible for the angular momentum loss due to the magnetized wind. Symbol $|_{\text{wind}}$ means that the value is taken at the base of the wind at the surface of the disk.

The tangential stress in the disk $t_{r\varphi}$ is defined by the turbulent motion and by the
FIG. 1. Schematic structure of the magnetic field inside the disk and in the wind. The wind is formed in the corona where all field lines are open. Inside the disk, the chaotic magnetic field is generated due to turbulent motion of plasma in the disk.

magnetic field. According to [1]

\[ t_{r\varphi} = -\alpha \rho v_s^2. \] (2)

where \( v_s \) is the sound velocity, \( \rho \) is the matter density and \( \alpha \) is the parameter connecting the pressure and tangential stress in the disk.

As it was pointed out by [36], the momentum loss due to the wind will dominate the losses caused by the viscous stresses provided that

\[ 4\pi rt_{r\varphi} h \ll r^2 \left< B_\varphi B_z \right>_\text{wind}. \] (3)

In the opposite case we have the standard [1] version of the disk accretion.

The integration of the energy conservation equation over the control volume gives

\[ \frac{\partial}{\partial r} \left( M \frac{v^2}{2} \right) - \frac{\partial}{\partial r} \left( 4\pi \Omega r^2 t_{r\varphi} h \right) + 4\pi \rho V_z E_{t\text{wind}} + 4\pi r Q = 0, \] (4)

where

\[ E_{\text{tot}} = \frac{v^2}{2} - \frac{B_z B_\varphi \Omega r}{4\pi \rho V_z} - \frac{GM}{r} \] (5)
is the total energy per particle in the wind, $Q$ is the energy radiated from unit square of one side of the disk surface. We neglect by the contribution from the turbulent Ohmic heat production and advection of the energy in the optically thick disk.

Taking into account the mass conservation

$$\frac{\partial \dot{M}}{\partial r} - 4\pi r \rho v_z |_{wind} = 0,$$  \hspace{1cm} (6)

a simple algebra with eqs. (1) and (4) gives like in [1], that the heat production equals

$$Q = t_{r\varphi} h r \frac{\partial \Omega}{\partial r}.$$  \hspace{1cm} (7)

Below we consider the case when the inequality in eq. (3) is fulfilled. Dissipative terms can be neglected in equations (1 and 4). The equation for the angular momentum conservation takes a form

$$\dot{M} \frac{\partial V_k}{\partial r} + r^2 < B_\varphi B_z > |_{wind} = 0.$$  \hspace{1cm} (8)

Energy conservation equation is reduced to

$$\frac{1}{2} \frac{\partial V_k^2 \dot{M}}{\partial r} + 4\pi r \rho v_z E_{tot} |_{wind} = 0.$$  \hspace{1cm} (9)

The luminosity of the disk is calculated by eq. (7).

This system of equations together with the system of MHD equations describing the flow of the wind outside the disk totally defines the dynamics of the disk and the wind in the regime of "cold" accretion. These equations do not contain dissipative terms connected with viscosity or finite conductivity. According to eq. (8),

$$< B_\varphi B_z > |_{wind} = \frac{\dot{M} V_k}{2r^2}.$$  \hspace{1cm} (10)

This allows us to estimate the minimum magnetic field at the base of the wind which can support the regime of the "cold" accretion:

$$B_{min} = \sqrt{\frac{\dot{M} V_k}{2r^2}}.$$  \hspace{1cm} (11)

It is easy to demonstrate that this value is small compared to the field in the Shakura-Sunyaev regime of accretion. Nevertheless, the presence of the magnetic field at the base of the wind means that the magnetic field is present inside the disk and this field generates
turbulence which makes inevitable heating and radiation from the disk. The main question for us is how strong is this radiation?

To connect the magnetic field at the base of the wind with $t_{r\varphi}$ we introduce a parameter

$$\theta = \frac{4\pi t_{r\varphi}}{< B_\varphi B_z > \mid_{\text{wind}}}.$$  \hspace{1cm} (12)

Physical sense of this parameter becomes clear if to pay attention that according to the assumptions of \cite{1} and recent numerical simulations of the magnetic field generation \cite{57}

$$- t_{r\varphi} \approx 0.4 \cdot \frac{B^2}{4\pi}.$$  \hspace{1cm} (13)

where the magnetic field $B$ is taken at the midplane of the disk. Roughly, $\theta$ equals to the ratio of the magnetic pressure inside the disk to the magnetic pressure at the base where the wind starts to flow. The ratio $\theta$ depends on variables like the mass flow rate $\dot{M}$, mass of the central object $M$, radius, and others. This is one of the key difference of our work from other works. We distinguish the magnetic field inside the disk from the magnetic field at the base of the wind. In the works of the Grenoble group $\theta \sim 1$ \cite{38, 39} and other authors \cite{40, 42} $\theta \sim 1$ because the magnetic field vertically crosses the disk. In reality, the magnetic field inside the disk can essentially exceed the field at the base of the wind \cite{57}. Therefore, below we assume that $\theta \geq 1$ but should satisfy to the condition of ”cold“ accretion following from eq.(3)

$$\frac{\theta h}{r} \ll 1.$$  \hspace{1cm} (14)

For the geometrically thin disks with $h/r \ll 1$ the ”cold“ disk accretion can take place even for $\theta > 1$.

**THERMAL RADIATION FROM THE DISK AT THE "COLD" ACCRETION**

In the work \cite{1}, three regimes of the disk accretion have been considered: a) when the radiation pressure exceeds the gas pressure and the Thomson scattering dominates over the free-free absorption; b) when the gas pressure dominates over the radiation pressure but the Thomson scattering dominates the free-free absorption; c) when the gas pressure dominates over the radiation pressure and the opacity of the matter is defined by the free-free absorption. We consider only the case when the gas pressure dominates over the radiation pressure. These are the regimes b) and c). Below we will see that when the radiation dominates, the accretion proceeds in the Shakura-Sunayev regime.
Scattering dominates over free-free absorption

Firstly, we consider the case when the Thomson scattering dominates over the free-free absorption. Hereafter we call this regime as Thomson regime. The pressure of radiation $P_{\text{rad}}$ equals to $\varepsilon/3$, where $\varepsilon = bT^4$. The sound velocity is defined as $v_s^2 = kT/m_p$, where $m_p$ is the proton mass. According to [1] the heat conductivity of the disk is defined by the transport of radiation. Then

$$\varepsilon = 3Q\sigma u_0 \frac{4}{c},$$

where $\sigma = 0.4$ cm$^2$/g is the Thomson opacity and $u_0 = 2\rho h$. Then the rate of heating of the disk equals

$$Q = \frac{3\theta\dot{M}V_k v_s}{16\pi r^2}.$$

We used here that $h = v_s/\Omega$. The solution of this system of equations gives

$$T = \frac{\sqrt{3}}{4\sqrt{\pi}} \left( \frac{\theta^2\dot{M}^2V_k\sigma}{b\alpha c r^3} \right)^{\frac{1}{4}}.$$

The sound velocity equals

$$v_s = \frac{6^{1/4}}{2(2\pi)^{1/4}} \frac{k^{1/2}V_k^{1/8}(\theta\dot{M})^{1/4}\sigma^{1/8}}{m_p^{1/2}b^{1/8}\alpha^{1/8}v_{\text{esc}}^{1/8}r^{-3/8}},$$

and the density flux of radiation from one side of the disk is expressed as

$$Q = \frac{(3\pi)^{5/4}6^{1/4}(\theta\dot{M})^{5/4}V_k^{9/8}k^{1/2}\sigma^{1/8}}{32 \cdot r^{-19/8}m_p^{1/2}b^{1/8}\alpha^{1/8}v_{\text{esc}}^{1/8}}.$$

Let us express $\dot{M} = \dot{m}\dot{M}_{\text{crit}}$, the radius $r$ in $r = (3r_g)x$ and the mass $M$ in the solar masses $M = mM_{\odot}$. In these variables we obtain

$$Q = 0.75 \cdot 10^{23} \frac{(\dot{m})^{5/4}}{m^{9/8}x^{47/16}\alpha^{1/8}} \cdot \text{erg/s/cm}^2.$$

The integration of this expression over the disk gives the bolometric luminosity of the disk

$$L_{\text{bol}} = 0.84 \cdot 10^{36} \frac{(\dot{m})^{5/4}m^{7/8}}{\alpha^{1/8}}, \text{ erg/s}.$$

The kinetic luminosity of the jets equals to the total energy release at accretion. Therefore

$$L_{\text{kin}} = \frac{\dot{M}c^2}{12} = 1.4 \cdot 10^{38} \dot{m}, \text{ erg/s}.$$
Then, the ratio of the kinetic luminosity over the bolometric luminosity equals

\[ \frac{L_{\text{kin}}}{L_{\text{bol}}} = 168 \frac{(m\alpha)^{1/8}}{\dot{m}^{1/4}\theta^{5/4}}. \]  

The bolometric luminosity can be expressed in conventional variables:

\[ L_{\text{bol}} = \frac{4}{5} \theta \dot{M} V_{k0} v_{s0}, \]

where \( V_{k0} \) and \( v_{s0} \) are the Keplerian and sound velocities at the inner edge of the disk. Taking into account that the kinetic luminosity

\[ L_{\text{kin}} = \frac{\dot{M} V_{k0}^2}{2}, \]

the condition \( L_{\text{bol}}/L_{\text{kin}} \ll 1 \) becomes

\[ \frac{8 \theta v_{s0}}{5 V_{k0}} = \frac{8 \theta h}{5 r} \ll 1, \]

which practically coincides with the condition of applicability of the "cold" disk accretion approximation defined by eq. (14). Similar condition has been obtained earlier in [46]. The condition \( L_{\text{kin}} \gg L_{\text{bol}} \) indicates that the accretion occurs in the "cold" regime.

The temperature in the disk

\[ T = 2.5 \cdot 10^7 \frac{\sqrt{\theta \dot{m}}}{\alpha^{1/4} x^{7/8} m^{1/4}} \text{ K}, \]

can be less than the temperature in the [1] disk provided that \( \theta < 100 \).

Let us calculate the ratio of the radiation pressure over the gas pressure in the disk,

\[ \frac{P_{\text{rad}}}{P_{\text{gas}}} = \frac{3}{32\pi} \frac{\theta \dot{M} \sigma}{r c} = 0.85 \frac{\theta \dot{m}}{x}. \]  

This means that all the estimates are valid at \( 0.85\theta \dot{m} < 1 \).

Other parameters are estimated in the disk as follows. Density equals

\[ \rho = \frac{1}{2\sqrt{3\pi}} \frac{\sqrt{\theta \dot{M} V_{k0}^3/4} m_p b^{1/4} c^{1/4}}{r^{5/4} k \alpha^{3/4} \sigma^{1/4}} = 0.6 \frac{\sqrt{\theta \dot{m}}}{m^{3/4} x^{13/8} \alpha^{3/4}} \text{ g/cm}^3. \]

The aspect ratio of the disk is

\[ \frac{h}{r} = 3.7 \cdot 10^{-3} \frac{(\dot{m} \theta)^{1/4} x^{1/16}}{\alpha m^{1/8}}. \]

The true optical depth \( \tau^* = \sqrt{\sigma \cdot \sigma_{ff} \cdot u_0} \) of the disk is expressed as

\[ \tau^* = 51(\theta \dot{m})^{1/8} m^{3/16} x^{5/32} \alpha^{-13/16}, \]
where
\[ \sigma_{ff} = 0.11 \cdot T^{-7/2} n, \; \text{cm}^2/\text{g} \]  
(32)
is the free-free opacity of the disk. The surface temperature of the disk \( T_s \) is defined from the equation \( bcT_s^4/4 = Q \) and has a form
\[ T_s = 6 \cdot 10^6 \frac{(\dot{m}m)^{5/16}}{m^{9/32}x^{47/64}\alpha^{1/32}} \; \text{K}. \]  
(33)

**Free-free absorption dominates over scattering**

At the condition
\[ 4.6 \cdot 10^{-3} \frac{(\alpha m)^{1/10}x^{23/20}}{\dot{m}^{6/17}} > 1 \]  
(34)
the Thomson scattering opacity becomes small compared with the free-free absorption. Hereafter we call this regime as free-free. Similar calculations give the following temperature inside the disk
\[ T = 10^7 \frac{(\dot{m}m)^{6/17}}{x^{12/17}(\alpha \cdot m)^{4/17}} \; \text{K}. \]  
(35)
The bolometric luminosity of the disk is
\[ L_{bol} = 0.6 \cdot 10^{36} (\dot{m}m)^{20/17} m^{15/17} \alpha^{-2/17} \; \text{erg/s}. \]  
(36)
The ratio of the kinetic luminosity over the bolometric luminosity equals
\[ \frac{L_{\text{kin}}}{L_{bol}} = \frac{228(\alpha m)^{2/17}}{\dot{m}^{3/17}x^{20/17}} \]  
(37)
The full optical depth, the density of plasma and the aspect ratio of the disk are
\[ \tau = 93(\dot{m}m)^{4/17} m^{3/17}x^{1/34} \alpha^{-14/17}, \]  
(38)\[ \rho = \frac{1.2(\dot{m}m)^{11/17}}{(\alpha m)^{13/17}x^{61/34}} \; \text{g/cm}^3, \]  
(39)\[ \frac{h}{r} = 2.5 \cdot 10^{-3} x^{5/34} (\dot{m}m)^{3/17} (\alpha m)^{-2/17}. \]  
(40)
Finally, the surface temperature is
\[ T_s = 5.5 \cdot 10^6 \frac{(\dot{m}m)^{5/17}}{m^{19/68}x^{97/136}\alpha^{1/34}} \; \text{K}. \]  
(41)
The fundamental plane encapsulates the relationship between the compact radio luminosity, X-ray luminosity, and the black hole mass and provides a good description of the data over a very large range of black hole mass. There are reasons to believe that the Fundamental Plane (FP) of the black holes reproduces the actual relationship between the kinetic luminosity of jets and bolometric luminosity of the disks. In the work [17], the position of objects of different masses in the coordinates $L_{\text{kin}}/L_{\text{bol}}$ and $L_{\text{bol}}/L_{\text{Edd}}$ has been collected in one FP. If this is true, the FP can be used to extract information about the dependence of $\theta$ on $\dot{m}$ and $m$. All data at the FP can be approximated by a power law function of the form

$$\log \frac{L_{\text{kin}}}{L_{\text{bol}}} = (A - 1) \log \frac{L_{\text{bol}}}{L_{\text{Edd}}} + B$$

with $A$ in the range (0.43 - 0.47) and $B$ in the range from -0.94 to -1.37. For estimates the values $A = 0.457$ and $B = -1.1$ have been chosen which are close to the average. The ratio of $L_{\text{bol}}/L_{\text{Edd}}$ in the Thomson regime is

$$\frac{L_{\text{bol}}}{L_{\text{Edd}}} = 6 \cdot 10^{-3} \frac{(\theta \dot{m})^{5/4}}{(\alpha m)^{1/8}},$$

while in free-free regime this ratio equals to

$$\frac{L_{\text{bol}}}{L_{\text{Edd}}} = 4.4 \cdot 10^{-3} \frac{(\theta \dot{m})^{20/17}}{(\alpha m)^{2/17}},$$

Obviously, at the constant $\theta$ the theoretical predictions are not consistent with observations. The reason is that $\theta$ must depend on $\dot{m}$ and $m$. The most natural option is to assume that $\theta$ depends on $\dot{m}$ as a power law

$$\theta = D \dot{m}^\gamma.$$  

In the Thomson regime

$$X = \frac{L_{\text{bol}}}{L_{\text{Edd}}} = 6 \cdot 10^{-3} \frac{\dot{m}^{5(\gamma + 1)/4} D^{5/4}}{(\alpha m)^{1/8}},$$

and

$$Y = \frac{L_{\text{kin}}}{L_{\text{bol}}} = \frac{168(\alpha m)^{1/8}}{D^{5/4} \dot{m}^{(5\gamma + 1)/4}}$$

After simple algebraic calculations, we obtain that

$$A = \frac{4}{5(\gamma + 1)},$$
For $A = 0.457$ the value $\gamma = 3/4 = 0.75$. Then $D = 5 \cdot 10^3(\alpha m)^{1/10}$. Thus, in the Thomson regime

$$\theta = 5 \cdot 10^3 \dot{m}^{3/4}(\alpha m)^{1/10}$$

(49)

Similar calculations in the free-free regime give

$$\theta = 11.4 \cdot 10^3 \dot{m}^{0.86}(\alpha m)^{1/10}.$$  

(50)

The power at $\dot{m}$ is chosen to provide uniform dependence of $Y$ on $X$ of the form (42) with constant $A$ in both regimes.

The dependencies (49) and (50) seem physically reasonable. They show that the smaller the accretion rate, more uniform is the magnetic field across the disk and, therefore, the $\theta$ is close to 1.

The plot of $\theta/(\alpha m)^{1/10}$ is presented in fig. 2. $\theta$ corresponding to FP agree with the assumption of the "cold" accretion because this curve is located well below the curve separating the regime of the cold accretion from the Shakura-Sunyaev regime. In fig. 2 we present in dashed-dotted the line which separates regions of domination of the gas pressure and radiation pressure defined by equation (28). Thin solid line separates the Thomson regime from the free-free regime.

**COMPARISON WITH THE SPECIFIC SOURCES**

It is interesting to apply the estimated dependencies to the specific sources. Below we consider M87 and the SMBH in galactic center, Sgr A*. We will see later that both sources are in the free-free regime. Therefore eq. (50) has been used at the estimations. For definiteness we accept $\alpha = 0.1$.

**M 87**

For this object $\dot{m}$ and $m$ can be easily estimated. The kinetic luminosity of this object is $L_{kin} = 10^{44}$ erg/s which we assume is equal to the total rate of the gravitational energy released at the accretion. Mass of the central black hole $m = 3.5 \cdot 10^9$ [62]. With the Eddington luminosity equal to $L_{Edd} = 1.4 \cdot 10^{38} m$ erg/s, we find $\dot{m} = L_{kin}/L_{Edd} = 2 \cdot 10^{-4}$. 

13
FIG. 2. Dependence of $\theta/(\alpha m)^{1/10}$ on $\dot{m}$. Shakura-Sunyaev regime of accretion takes place above thick dashed line. Below this line the regime of "cold" accretion takes place. The Thomson scattering dominates above the thin solid line, while below this line the free-free absorption gives the major contribution into the opacity of the medium. Dashed-dotted line (calculated at $m = 10^8$) divides the plane in parts where the radiation pressure (above this line) or the gas pressure (below) dominate.

From eq. (37) we obtain that

$$\frac{L_{kin}}{L_{bol}} = 10^4 \theta^{-20/17}. \quad (51)$$

Eq. (50) gives $\theta = 54$. Then $L_{kin}/L_{bol} \approx 95$ and $L_{bol}$ from eq. (36) equal to $10^{42}$ erg/s in accordance with observations. Optical depth of the disk exceeds $\tau > 10^4$.

**Sgr A***

The kinetic luminosity of the outflow from the disk around SMBH in Sgr A* is not known. The flux of TeV gamma-rays from the Galactic Center can be explained by very high energy accelerated protons with a luminosity close to $10^{38}$ erg/s. The kinetic luminosity of the wind has to be higher. We consider a reversed problem. Assuming that the dependence (50) is
valid for the disk in Sgr A* we can predict what is the kinetic luminosity of wind from the
disk in the Galactic Center. Let us to estimate $\dot{m}$ from the bolometric luminosity of the
disk (see eq. (36)). In this case

$$\dot{m} = \left( \frac{L_{bol}}{3.6 \cdot 10^{40} \text{ (erg/s)}} \right)^{17/20} \frac{1}{m} = 8 \cdot 10^{-6}, \quad (52)$$

$\theta = 1.7$ and from eq. (22) we obtain that

$$L_{kin} = 4.4 \cdot 10^{39}, \text{ erg/s,} \quad (53)$$

at $L_{bol} = 10^{36}$ erg/s. The kinetic luminosity of the wind from the Galactic accretion disk
4.4 $\cdot$ 10$^3$ times exceeds the bolometric luminosity of the disk. Remarkably, this power is
sufficient to explain the flux of the PeV protons in the Galactic Center.

**DISCUSSION**

The main result of this work is the conclusion that the main source of energy in AGN
is the energy released at the accretion even in the case of high ratio of the kinetic-to-
bolometric luminosity. Observation of such objects can point out that the accretion occurs
in the "cold" regime. Outflow in the form of a wind from the accretion disk carry out
the largest fraction of the angular momentum of the accreted material. This results into
suppression of radiation from the disk. Accretion onto SMBH appear very efficient process
which transforms practically all the gravitational energy into the kinetic energy of jets. This
strongly contrasts with also radiationally inefficient ADAF models [63]. In the last case
AGN appears very inefficient machine which transforms almost all the gravitational energy
into mass and rotational energy of the SMBH. This model needs an additional source of
energy which is believed can be the rotational energy of SMBH.

The accretion in the "cold" regime occurs even when the magnetic field inside the accre-
tion disk essentially exceeds the magnetic field at the base of the wind. This is explained by
the geometrical reason. The angular momentum transport due to viscosity is proportional
to the magnetic pressure in the disk times the thickness of the disk $h$ while the flux of the
angular momentum from the disk is proportional to the magnetic pressure at the base of the
wind times the radius. The ratio of the viscous losses to the losses due to wind is $\sim \theta \cdot (\frac{h}{r})$,
where $\theta$ roughly equals to the ratio of the magnetic pressures inside and at the surface of the
disk. Therefore, the Shakura-Sunyaev regime of accretion \[1\] is realized when \( \theta \gg \left( \frac{r}{h} \right) \). The assumption that \( \theta \) is constant contradicts to the observations if to assume that FP correctly reproduces the ratio of the kinetic-to bolometric luminosities. \( \theta \) increases on 2 orders of magnitude with \( \dot{m} \). This value of \( \theta \) agrees with the results of modeling of MRI dynamo \[57\]. This estimate should be considered as only a rough approximation. Nevertheless, it allows us to make certain conclusions about realization of the regime of "cold" disk accretion. It appears that at small accretion rates \( \dot{m} < 10^{-2} \), the estimated value of \( \theta \) is located in the region well below the line were the Shakura-Sunyaev model is valid. The magnetic pressure inside the disk appears less than the magnetic pressure estimated in the model \[1\]. Therefore, it is quite reasonable to assume that at relatively low rates of accretion, \( \dot{m} < 10^{-2} \), the accretion occurs predominantly in the regime of the "cold" accretion. At higher values of \( \dot{m} > 0.1 \) the accretion occurs in the regime of Shakura & Sunyaev. The transition between the regimes takes place at the value of \( \dot{m} \) between 0.01 and 0.1 which well agree with location of the transition from very bright to very dim disks around SMBH with powerful outflow deduced in \[64\]. Remarkably, the rough estimate of the dependence of \( \theta \) on \( \dot{m} \) gives good agreements with observations of two SMBHs, M87 and Sgr A*.

In this paper, we have considered only the case when the gas pressure dominates over the radiation pressure. But according to fig. 2, the radiation pressure dominates basically in the Shakura-Sunyaev accretion regime. The "cold" accretion at the domination of the radiation pressure apparently may be realized in rather narrow range of parameters. This case is planned to be considered separately.

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