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Low-luminosity AGNs

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ABSTRACT

Context. We propose that low-luminosity AGNs (LLAGNs), or some of them, are sources extracting their energy from the black hole rotation by the Blandford-Znajek mechanism.

Aims. It is shown that almost all energy of the black hole rotation is converted to relativistic protons in a jet. Owing to the high magnetic-field magnitude near the black hole, required for the Blandford-Znajek mechanism, electrons are not strongly accelerated because of their high synchrotron losses. Conversely, protons gain energies on the order of $(10^4–10^5)m_pc^2$ when crossing the light cylinder surface. Protons are also accelerated in a disk by 2D turbulent motion of the disk matter.

Methods. We calculate the luminosity of the synchrotron radiation by fast protons in the disk, the frequencies of this radiation being in the infrared band, and the luminosities corresponding to LLAGNs. We measure the very high energy (VHE) radiation luminosities from the disk and the jet, finding that VHE radiation is produced by collisions of accelerated protons with surrounding matter.

Results. We predict a correlation between the infrared luminosity $L_{\rm IR}$ and the VHE luminosity $L_{\rm VHE}$ of the disk, $L_{\rm VHE} \propto L_{\rm IR}^{1/2} M^{3/2}$, where M is the mass of a black hole. Two low-luminosity sources Sgr A* and M 87, for which luminosities $L_{\rm IR}$ and $L_{\rm VHE}$ are known, appear to follow this scheme.

Conclusions. The discovery of new bright VHE sources from LLAGNs could confirm our hypotheses that they are energy sources powered by the Blandford-Znajek mechanism.

Key words. galaxies: active – acceleration of particles – accretion, accretion disks – gamma rays: galaxies

1. Introduction

The common point of view about active galactic nuclei (AGN) is that they are massive black holes placed in the galactic center surrounded by an accretion disk. Energy-release mechanisms that have been proposed up to now are the transformation of the gravitation energy of accreting matter into heat and radiation, and the extraction of energy from the black hole rotation. The latter is known as the Blandford-Znajek mechanism (1977). The former mechanism gives an AGN luminosity $L = \eta_r \dot{M} c^2$, which is proportional to the mass rate accretion $\dot{M}$, η_r being the radiative efficiency (the standard value of which is $\eta_r \simeq 0.1$ for bright AGNs). For the Eddington luminosity $L_{\rm Edd} = 4\pi G M m_p c^2/\sigma_T = 1.3 \times 10^{38}(M/M_\odot)$ erg/s, the accretion rate is $\dot{M}_{\rm Edd} = 1.4 \times 10^{17}\eta_r^{-1}(M/M_\odot)$ g/s $= 2.3\times10^{-9}\eta_r^{-1}(M/M_\odot)\ M_\odot\ {\rm yr}^{-1}$, where G is the gravitational constant, m_p is the proton mass, c is the speed of light, and σ_T is the Thomson cross-section. The value M is the black hole mass, and $M_\odot$ is the solar mass. Accretion of the interstellar gas into a black hole is defined by its gravitational radius $R_G = 2MG/c_s^2$, at which the gravitational energy of a particle in the field of a black hole equals the particle thermal energy. The value of c_s is the gas sound velocity. In this case, the accretion rate is equal to $\dot{M} \simeq \pi R_G^2 c_s \rho$, where ρ is the gas density. The gravitational radius R_G is much larger than the radius of the horizon of a Schwarzchild black hole $r_H = 2MG/c^2$, $R_G = r_H(c/c_s)^2$. Nevertheless, the accretion rate of interstellar gas into a black hole is small because of the low gas density, $\dot{M} \simeq \pi r_H^2 c^2 c_s^{-1}\rho = 4.2\times10^2(M/M_\odot)^2(T/100\ {\rm K})^{-1/2}(n/0.1\ {\rm cm}^{-3})$g s^{-1}, where T is the gas temperature and n is the gas concentration $n = \rho/m_p$. If the accretion is fed by a dense molecular cloud or a disrupted star, then the accretion flow has the form of an accretion disk, which possesses a large angular momentum. The accretion rate $\dot{M}$ for disks can be as high as the Eddington rate $\dot{M}_{\rm Edd}$ and depends on the disk structure. The bolometric luminosity of AGN L_b is often less than the Eddington luminosity $L_{\rm Edd}$ by a factor $10^{-3}–10^{-1}$. A noticeable fraction of nearby galaxies do not contain bright AGNs, but instead contain low-luminosity AGNs (LLAGNs) for which $L_b \approx (10^{-5}–10^{-3})L_{\rm Edd}$. The center of our Galaxy Sgr A* is dimmer, its bolometric luminosity being only $L_b \approx 10^{36}$ erg/s $\approx 2\times10^{-9}\ L_{\rm Edd}(M \approx 3.6\times10^6\ M_\odot)$. The reason for this low luminosity is the low value of the radiative efficiency η_r or low value of the accretion mass rate $\dot{M}$. Low radiative efficiency is inherent to the advection-dominated accretion flow (ADAF) (Narayan 1994). At low accretion mass rates of $\dot{M} \ll \dot{M}_{\rm Edd}$, the radiative cooling of the disk becomes inefficient. These accretion flows are generally called radiatively inefficient accretion flows (RIAFs) (Narayan 2002). This class of disk flow also includes ADAFs. The accretion mechanism of the energy release is ineffective for LLAGNs.

We discuss the situation when the energy release due to accretion into the black hole of an AGN is small, and the energy activity is determined generally by the Blandford-Znajek mechanism. As we see below, the main energy output from a rotating black hole is in a relativistic jet. Thus, we mean LLAGNs with a jet. They are low-luminosity Seyfert galaxies and LINERs, the central source of which is a low luminosity AGN. A LINER is the low-ionization nuclear emission-line region at the center of a bright galaxy. LINERs are characterized by collisionally excited lines of neutral and singly ionized gas (Maoz 2007).

We first describe in Sect. 2 how the Blandford- Znajek mechanism operates. We describe configurations of the magnetic field and electric currents in both a disk and the black hole magnetosphere above a disk. We also show how the energy and angular momentum are transmitted from the black hole rotation to a jet. In Sect. 3, we calculate the synchrotron radiation of fast protons in a disk. The production of high energy γ-rays by relativistic protons is presented in Sect. 4. The observational consequences of the suggested scheme are described in the last section.

2. Blandford-Znajek mechanism

The Blandford-Znajek mechanism assumes that energy and angular momentum are extracted from the black hole rotation. The rotation energy $E_{\rm rot}$ stored in the black hole rotation is large, for slow rotation being $E_{\rm rot} = Mr_{\rm H}^2\Omega_{\rm H}^2/2 = a^2Mc^2/8 = 2.25\times10^{53}a^2(M/M_\odot)$ erg. Here we introduce the dimensionless parameter of rotation, $a = Jc/M^2G$, where J is the angular momentum of the black hole. For a black hole, $a < 1$. The value of $\Omega_{\rm H}$ is the angular velocity of black hole rotation, $\Omega_{\rm H} = ac/2r_{\rm H}$. The extraction of the rotation energy is possible if there exists a poloidal magnetic field B near the black hole horizon. The black hole in this case works as a dynamo machine, creating the voltage U, $U = \Omega_{\rm H} f_{\rm H}/2\pi c$ (Landau & Lifshitz 1984; Thorne et al. 1986), where $f_{\rm H}$ is the flux of the magnetic field reaching the horizon, $f_{\rm H} \approx \pi B r_{\rm H}^2$. The voltage generates the electric current $I = U/(R + R_{\rm H})$, which on one side is closed at the black-hole horizon surface of resistivity $R_{\rm H} = 4\pi/c \approx 377$ ohms (Thorne et al. 1986). The resistivity of the outer part of the system is R. Thus, the extracted power is $L = U^2R/(R + R_{\rm H})^2 = a^2B^2r_{\rm H}^2R/16(R + R_{\rm H})^2$. The power L reaches its maximum value $L_{\rm m}$ when $R = R_{\rm H}$, $L_{\rm m} = a^2B^2r_{\rm H}^2c/256\pi$. This value is proportional to the square of the black hole mass, $L_{\rm m} \propto M^2$, and for large enough magnetic fields can exceed the Eddington luminosity, $B > B_{\rm Edd} = 5.5\times10^9a^{-1}(M/M_\odot)^{-1/2}$ Gauss. For high masses of AGN black holes, the value of the magnetic field $B_{\rm Edd}$ is quite moderate, $B_{\rm Edd} \simeq 10^5$ Gauss ($a \simeq 1$).

We note that to ensure in general that the gravitation energy release is extremely efficient, there must be a high mass accretion rate $\dot{M}$ on to the massive black hole and a high radiative efficiency $\eta_{\rm r}$, whereas the Blandford-Znajek mechanism provides such efficiencies when there is black hole rotation, $\Omega_{\rm H} \neq 0$, and a strong enough magnetic field B near the black hole horizon. Formally, the Blandford-Znajek mechanism does not need accretion. In addition, the accumulation of a strong magnetic field requires some accretion process, but accretion is not a source of energy.

2.1. Magnetic field and electric current configurations

For disk accretion, the magnetic field inside the disk and nearby must have no component that is perpendicular to the disk (Istomin & Sol 2009). In the axisymmetric stationary electromagnetic field, a charged particle conserves the generalized angular momentum, $\rho p_\phi + qf/c = $ const., where ρ is the cylindrical distance from the center, f is the flux of the poloidal magnetic field, and q is the charge of a particle. However, two terms in this relation are practically not commensurable in the case of disk accretion. The first one is proportional to the frequency of rotation of a particle in the disk v_ϕ/ρ, the second one is proportional to the cyclotron frequency, $\omega_{\rm c} = qB_z/mc$, of the particle rotation in the perpendicular magnetic field B_z. A charged particle then can move in the radial direction if it is not magnetized, $\omega_{\rm c} \simeq v_\phi/\rho$, i.e. in a practically zero perpendicular magnetic field. The radial B_ρ and the azimuthal B_ϕ components can be arbitrary. We see that the accretion of matter provides only the radial component of the poloidal magnetic field towards the black hole vicinity. And the magnetic field in the expression for the Blandford-Znajek luminosity is the radial magnetic field $B = B_\rho$ produced by the conducting matter of the accretion disk. For there to be a zero component of the magnetic field perpendicular to the disk, an electric current I_ρ must flow through the disk, which is just the current $I = -I_\rho$ generated by the voltage U. The accretion disk is not only the origin of the flux of the matter on to the black hole, but also conducts the electric current. This current then flows onto the black hole horizon and closes through a jet in outer space. Owing to the spiral motion of charged particles in the disk, there exist not only a radial current j_ρ but also an azimuthal current j_ϕ. Only the azimuthal current creates the radial magnetic field. The ratio of the radial current to the azimuthal one $\alpha_j = j_\rho/j_\phi$ can be found by considering that from one side the current I is produced by the black hole rotation, $I = U/(R + R_{\rm H}) = (a/16\pi)B_\rho r_{\rm H}c[R_{\rm H}/(R + R_{\rm H})]$, and from another side the Maxwell equation determines the radial magnetic field through the azimuthal current, $I = \alpha_j r_{\rm H} c B_\rho$. We obtain $\alpha_j = (a/16\pi)[R_{\rm H}/(R + R_{\rm H})] \ll 1$, which shows that the toroidal magnetic field B_ϕ is much less than the radial magnetic field B_ρ, $B_\phi = \alpha_j B_\rho$.

2.2. Jet power

The radial magnetic field at the black hole horizon allows the transfer of energy and angular momentum from the black hole rotation to particles of the black hole magnetosphere and to particles leaving the accretion disk. The lines of the radial magnetic field begin to rotate with the angular velocity $\Omega_{\rm F} < \Omega_{\rm H}$, $\Omega_{\rm F} = \Omega_{\rm H}R/(R + R_{\rm H})$ (Thorne et al. 1986). For the optimal condition $R = R_{\rm H}$, $\Omega_{\rm F} = \Omega_{\rm H}/2$ (Blandford & Znajek 1977). Forced by the centrifugal acceleration towards the light cylinder surface ($r_{\rm L} = c/\Omega_{\rm F} = 2a^{-1}r_{\rm H}R/(R + R_{\rm H})$), particles achieve azimuthal velocities, which are close to the speed of light, and significant energies. The main energy is in protons because of their low synchrotron losses in the strong magnetic field. The Lorentz factor γ of protons on the light cylinder surface is $\gamma = (\omega_{\rm cL}/\Omega_{\rm F})^{1/2}$ (Istomin & Sol 2009), where $\Omega_{\rm cL} = eB_{\rm L}/m_{\rm p}c$ is the non – relativistic proton cyclotron frequency in the magnetic field $B_{\rm L}$ on the light cylinder surface. Taking into account that the radial magnetic field falls as ρ^{-1}, we obtain $\gamma = (\omega_{\rm cH}r_{\rm H}/c)^{1/2}$, where $\omega_{\rm cH}$ is the proton cyclotron frequency near a black hole. Almost all particle energy is in the azimuthal motion, $p_\phi \simeq m_{\rm p}\gamma c$, and only a small part is in the radial one, $p_\rho \simeq m_{\rm p}\gamma^{1/2}c$ (Istomin & Sol 2009). Thus, rotating energetic protons slowly flow outside the light cylinder surface, $v_\rho \simeq c\gamma^{-1/2}$, forming the relativistic jet. The energy density of particles on the light cylinder is equal to the density of the electromagnetic energy $B_{\rm L}^2/8\pi + E_{\rm L}^2/8\pi = B_{\rm L}^2/4\pi\,(E_{\rm L} = B_{\rm L})$ (Istomin 2010). We can then calculate the jet luminosity $L_{\rm J} = cr_{\rm L}^2B_{\rm L}^2\gamma^{-1/2}/2 = B^2r_{\rm H}^2c(\omega_{cH}r_{\rm H}/c)^{-1/4}/2$. The total Blandford-Znajek luminosity L must, of course, be higher than the jet luminosity $L_{\rm J}$. This implies the definite condition for the magnetic field stress near the black hole

$$\frac{\omega_{\rm cH}r_{\rm H}}{c} \geq (128\pi)^4a^{-8}\left[\frac{(R + R_{\rm H})^2}{4RR_{\rm H}}\right]^4, \tag{1}$$

or for $R = R_{\rm H}$

$$B \geq 2.7 \times 10^{11} a^{-8} \frac{M_\odot}{M} \text{ Gauss.} \tag{2}$$

For AGN black-hole masses of $M \simeq 10^8\ M_\odot$, this condition ($B > 3 \times 10^3$ Gauss) is not onerous and AGN with such moderate magnetic fields can produce a relativistic jet. Less massive black holes must have higher magnetic fields. At the center of our Galaxy, for example, there must be $B \approx 10^5$ Gauss. However, a magnetic field of the order of 10^{11} Gauss for micro-quasars seems problematic when producing relativistic jets. However, Karitskaya et al. (2009) measured the disk magnetic field to be 600 Gauss at a distance $2 \times 10^5 r_{\rm H}$ in Cygnus X-1. We see in Eqs. (1) and (2) a strong dependence of the magnitude of the magnetic field near the black hole B on the black hole rotation, $B \propto a^{-8}$. This increases the value of magnetic fields for slowly rotating black holes. We note, however, that the rotation cannot be too slow: the light cylinder surface must be inside the jet radius for producing a relativistic jet, $r_{\rm J} > r_{\rm L}$, $a > 4r_{\rm H}/r_{\rm J}$.

2.3. The resistivity R

The effective work done by the Blandford-Znajek mechanism depends on the value of resistivity R. The maximum output is for $R = R_{\rm H}$. However, the real resistivity of the system can differ from this value. We now estimate the resistivity $R_{\rm c}$ of the current loop created by the unipolar inductor voltage U. The current resistivity, of course, is determined by the electron motion. We assume the conducting system to be a box with the cross-section S and the length l along the direction of the electric current. The resistivity of this system is then $R_{\rm c} = l/S\sigma$, where σ is the electron conductivity, $\sigma = ne^2\tau_{\rm e}/m_{\rm e}$, n is the electron density, $m_{\rm e}$ and e are its mass and charge, and $\tau_{\rm e}$ is the relaxation time of electrons. Coulomb collisions of electrons dominate even in a low ionized plasma, and we can write the electron conductivity in the form

$$\sigma = \eta_\sigma \begin{cases} \frac{\varepsilon^{3/2}}{\pi m_{\rm e}^{1/2} e^2 \Lambda}, & \varepsilon < m_{\rm e}c^2 \\ \frac{\varepsilon^2}{\pi m_{\rm e} c e^2 \Lambda}, & \varepsilon > m_{\rm e}c^2. \end{cases} \tag{3}$$

We introduce the coefficient $\eta_\sigma < 1$, which takes into account the possible abnormal electron conductivity due to turbulent or another processes decreasing the electron conductivity. The other parameters are ε the mean electron energy, and Λ the Coulomb logarithm, $\Lambda \approx 15{-}20$. Thus, the electron resistivity is

$$R_{\rm c} = \frac{\pi l \Lambda \eta_\sigma^{-1} e^2}{S} \begin{cases} \frac{m_{\rm e}^{1/2}}{\varepsilon^{3/2}}, & \varepsilon < m_{\rm e}c^2 \\ \frac{m_{\rm e}c}{\varepsilon^2}, & \varepsilon > m_{\rm e}c^2. \end{cases} \tag{4}$$

The electron resistivity strongly depends on the electron energy, decreasing as the energy increases. The energy ε can be estimated from the bolometric luminosity $L_{\rm b}$ in a continuous spectrum, $L_{\rm b} = S_1 \sigma_S \varepsilon^4$, where S_1 is the surface of the current system, $S_1 \approx 2S + 4S^{1/2} l$. The constant $\sigma_{\rm S}$ is the Stefan-Boltzmann constant in energetic units, $\sigma_{\rm S} = \pi^2/60\hbar^3c^2 = 1.6 \times 10^{59}$ erg^{-3} cm^{-2} s^{-1}. Substituting the expression $\varepsilon = (L_{\rm b}/S_1\sigma_{\rm S})^{1/4}$ into Eq. (4), we get

$$R_{\rm c} = R_{\rm H} \left(\frac{L_1}{L_{\rm b}}\right)^{3/8}, \ \varepsilon < m_{\rm e}c^2, \tag{5}$$

$$R_{\rm c} = R_{\rm H} \left(\frac{L_2}{L_{\rm b}}\right)^{1/2}, \ \varepsilon > m_{\rm e}c^2. \tag{6}$$

Characteristic luminosities L_1 and L_2 are defined as

$$\begin{aligned} L_1 &= S_1 \left(\frac{l}{S}\right)^{8/3} \left(\frac{\Lambda \eta_\sigma^{-1} m_{\rm e}^{1/2} e^2 c}{4}\right)^{8/3} \sigma_{\rm S} \\ &= 3.8 \times 10^{15} \left(\frac{S_1}{1\ {\rm pc}^2}\right) \left(\frac{l \times 1\ {\rm pc}}{S}\right)^{8/3} \left(\frac{\Lambda}{20}\right)^{8/3} \eta_\sigma^{-8/3} {\rm erg/s}, \end{aligned} \tag{7}$$

$$\begin{aligned} L_2 &= \left(\frac{l^2 S_1}{S^2}\right) \left(\frac{\Lambda \eta_\sigma^{-1} m_{\rm e} e^2 c^2}{4}\right)^2 \sigma_{\rm S} \\ &= 1.4 \times 10^{11} \frac{l^2 S_1}{S^2} \left(\frac{\Lambda}{20}\right)^2 \eta_\sigma^{-2}\ {\rm erg/s}. \end{aligned} \tag{8}$$

Equations (7) and (8) show that for the real bolometric luminosities of AGN and LLAGN the resistivity $R_{\rm c}$ is very small, $R_{\rm c} \ll R_{\rm H}$, and the current system is far from the optimal condition $R_{\rm c} = R_{\rm H}$. We note that for relativistic electrons the characteristic luminosity L_2 does not depend on the size of the system, and the factor $l^2 S_1/S^2$ is only the geometric factor which is of the order of unity. In contrast, the luminosity L_1 for non-relativistic electrons is inversely proportional to the system size l, $L_1 \propto l^{-2/3}$. According to this estimation, we can conclude that the Blandford-Znajek mechanism for extracting energy from rotating black holes is ineffective for Ohmic heating of outer space. The only possible way to extract power by means of the Blandford-Znajek mechanism is to transform it into relativistic jet luminosity $L_{\rm J}$. We can attribute to the jet some value of the resistivity $R_{\rm J}$ using the relation $L_{\rm J} = R_{\rm J} I^2$. Neglecting the electron resistivity, we have $L = L_{\rm J}$. This implies that the equation for determining the jet resistivity is

$$\frac{(R_{\rm J} + R_{\rm H})^2}{4 R_{\rm J} R_{\rm H}} = \frac{a^2}{128\pi} \left(\frac{\omega_{\rm cH} r_{\rm H}}{c}\right)^{1/4}. \tag{9}$$

Denoting the quantity $\kappa = a^2(\omega_{\rm cH} r_{\rm H}/c)^{1/4}/128\pi$, ($\kappa \geq 1$), we find that

$$R_{\rm J} = R_{\rm H}[2\kappa - 1 \pm 2(\kappa^2 - \kappa)^{1/2}]. \tag{10}$$

Only for $\kappa = 1$ do we achieve the optimal efficiency of the central machine, $R = R_{\rm H}$. For $\kappa > 1$, we can assign to the jet two values of resistivity, one greater than $R_{\rm H}$ (the sign + in Eq. (10)), another less than $R_{\rm H}$ (the sign – in Eq. (10)). For both values, the jet luminosity is the same. However, the solution in Eq. (10) with the negative sign is unstable because a fluctuation of the magnetic field δB in this case results in a negative feedback with the power of the central machine, $\delta L/\delta B < 0$, while the value $\delta L_{\rm J}/\delta B$ is always positive since $\delta L_{\rm J}/\delta B = 7L_{\rm J}/4B$.

3. Acceleration of protons in accretion disk: proton synchrotron radiation

Producing a relativistic jet, a rotating black hole transmits an electric current I of high magnitude through an accretion disk. This current creates the magnetic field not only outside the disk, but also inside. Internal magnetic fields B_ρ and B_ϕ, such that $B_\phi \ll B_\rho$, are of the same order as that above the disk, except that they are equal to zero at the disk equator. Fields are frozen to the disk plasma motion. Therefore, a turbulent motion in the disk induces a turbulent electric field $\tilde{\boldsymbol{E}} = -[\tilde{\boldsymbol{u}}\boldsymbol{B}]/c$, where $\tilde{\boldsymbol{u}}$ is the turbulent plasma velocity. This stochastic electric field accelerates disk particles. This scenario of particle acceleration by large-scale 2D turbulence in a disk was discussed by Istomin & Sol (2009). They found that protons are indeed accelerated by this mechanism up to high energies. The maximum value of the proton Lorentz factor $\gamma_{\rm m} \gg 1$. Electrons are almost not accelerated at all because of large synchrotron losses (Istomin & Sol 2009). Thus, for the pure Blandford-Znajek mechanism almost no non-thermal high-energy radiation is produced by disk electrons because of the strong magnetic field and, connected with this, large synchrotron losses in any acceleration process.

In contrast, fast disk protons can radiate the synchrotron emission in strong magnetic fields. The distribution function of fast protons follows a power-law function $f_{\rm p} = b\gamma^{-\beta}$. The index β is the ratio of the loss energy rate to the rate of acceleration by stochastic electric field (Istomin & Sol 2009)

$$\beta = \sigma_{\rm E} c n \tau_{\rm c} \frac{c^4}{\langle \tilde{u}_\phi^2 \rangle u_\phi^2}, \tag{11}$$

where n is the proton density, $\sigma_{\rm E}$ is the cross-section of the proton-proton collisions in the disk $\sigma_{\rm E} \approx 10^{-26}$ cm^2, and the time $\tau_{\rm c}$ is the correlation time of the turbulence. We consider that relativistic protons have energies in the range $10 < \gamma < 10^8$ eV$/T_{\rm d}$, where the pricipal means of energy loss for fast protons is in their collisions with disk protons and $T_{\rm d}$ is the disk temperature in eV units. The correlation time $\tau_{\rm c}$ is of the order of the time of the azimuthal gyration of the matter in the disk, $\tau_{\rm c} \simeq \rho/u_\phi \propto \rho^{3/2}$. We consider the velocity u_ϕ to be $u_\phi \propto \rho^{-1/2}$, as in the Keplerian disk. The same law applies to the turbulent motion, $\tilde{u}_\phi \propto u_\phi$, and the disk density is $n \propto \rho^{-2}$. We see that the power-law index β increases with the radius ρ as $\rho^{3/2}$, $\beta = \beta_{\rm H}(\rho/r_{\rm H})^{3/2}$, where $\beta_{\rm H}$ is the index value near the black hole. The density of energetic protons is determined by the condition that their energy density is of the order of the energy density of the stochastic electric field $\tilde{E}^2/8\pi = B_\rho^2 \langle \tilde{u}_\phi^2 \rangle / 8\pi c^2$. Therefore, the distribution function of fast protons is

$$f_{\rm p} = (\beta - 2) \frac{B_\rho^2 \langle \tilde{u}_\phi^2 \rangle}{8\pi m_{\rm p} c^4} \gamma^{-\beta}, \; \beta > 2. \tag{12}$$

We see that the fast proton density decreases as $\propto \rho^{-3}$. Substituting this function into the well known expression for synchrotron radiation of one particle $P_1 = 2(e^2/m_{\rm p}c^2)^2 c B^2 \gamma^2/3$ and integrating over γ up to $\gamma_{\rm m}$, we find that the density of the synchrotron power W is

$$W = \frac{1}{12\pi} \frac{(\beta-2)}{(3-\beta)} \left(\gamma_{\rm m}^{3-\beta} - 1\right) \frac{\langle \tilde{u}_\phi^2 \rangle}{c^2} \left(\frac{\omega_{\rm c}}{c}\right)^3 m_{\rm p} c^2 \omega_{\rm c}. \tag{13}$$

The total synchrotron luminosity $L_{\rm s}$ is the integral of Eq. (13) over the disk volume

$$L_{\rm s} = \frac{1}{12} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \left(\frac{\omega_{\rm cH} r_{\rm H}}{c}\right)^3 m_{\rm p} c^2 \omega_{\rm cH} \times \int_1^\infty x^{-3} \frac{(\beta-2)}{(\beta-3)} (\gamma_{\rm m}^{3-\beta} - 1) {\rm d}x, \; x = \rho/r_{\rm H}. \tag{14}$$

Here we introduce the dimensionless disk width near the black hole, $h = H/r_{\rm H}$. Because the index β increases with distance from the black hole, i.e., $\beta = \beta_{\rm H} x^{3/2}$, the main contribution to the total synchrotron luminosity comes from the inner part of the disk, $x < x_1 = (3/\beta_{\rm H})^{2/3}$ if $\beta_{\rm H} < 3$. For $\beta_{\rm H} > 3$, the luminosity is lower. For $\ln\gamma_{\rm m} \gg 1$, the result of the integration in Eq. (14) with logarithmic accuracy is

$$L_{\rm s} = \frac{1}{18} \frac{(\beta_{\rm H}-2)}{\beta_{\rm H}(3-\beta_{\rm H})} \frac{\gamma_{\rm m}^{3-\beta_{\rm H}}}{\ln\gamma_{\rm m}} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \times \left(\frac{\omega_{\rm cH} r_{\rm H}}{c}\right)^3 m_{\rm p} c^2 \omega_{\rm cH}, \text{ if } \beta_{\rm H} < 3, \tag{15}$$

$$L_{\rm s} = \frac{1}{24} \frac{(\beta_{\rm H}-2)}{(\beta_{\rm H}-3)} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \left(\frac{\omega_{\rm cH} r_{\rm H}}{c}\right)^3 m_{\rm p} c^2 \omega_{\rm cH}, \text{ if } \beta_{\rm H} > 3. \tag{16}$$

To generate a relativistic proton jet, an AGN must have a strong magnetic field near the central black hole. This condition is given by Eq. (1), $\omega_{\rm cH} r_{\rm H}/c \geq (128\pi)^4 a^{-8}$. Using that, we obtain

$$L_{\rm s} \geq \frac{(128\pi)^{16}}{18} a^{-32} \frac{(\beta_{\rm H}-2)}{\beta_{\rm H}(3-\beta_{\rm H})} \frac{\gamma_{\rm m}^{3-\beta_{\rm H}}}{\ln\gamma_{\rm m}} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \frac{m_{\rm p} c^3}{r_{\rm H}} = 2\times 10^{41} a^{-32} \frac{(\beta_{\rm H}-2)}{\beta_{\rm H}(3-\beta_{\rm H})} \frac{\gamma_{\rm m}^{3-\beta_{\rm H}}}{\ln\gamma_{\rm m}} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \frac{M_\odot}{M} \text{ erg/s, if } \beta_{\rm H} < 3, \tag{17}$$

$$L_{\rm s} \geq \frac{(128\pi)^{16}}{24} a^{-32} \frac{(\beta_{\rm H}-2)}{(\beta_{\rm H}-3)} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \frac{m_{\rm p} c^3}{r_{\rm H}} = 1.6\times 10^{41} a^{-32} \frac{(\beta_{\rm H}-2)}{(\beta_{\rm H}-3)} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \frac{M_\odot}{M} \text{ erg/s, if } \beta_{\rm H} > 3. \tag{18}$$

For our Galaxy, the estimated synchrotron luminosity of the disk is $L_{\rm s} \simeq 10^{35}$ erg/s, which is close to its bolometric luminosity $L_{\rm b} \simeq 10^{36}$ erg/s. In any case, the synchrotron luminosity from a turbulent disk given by Eqs. ((15), (16)) is always much less than the total power extracted from a AGN rotating black hole.

The frequencies of radiation are in the range $\nu < 0.07\omega_{\rm cH}\gamma_{\rm m}^2$. According to Eq. (1) $\omega_{\rm cH} \geq (128\pi)^4 a^{-8} c/r_{\rm H}$, we estimate that $\nu \simeq 2\times 10^{14} a^{-8} (M_\odot/M) \gamma_{\rm m}^2\, s^{-1}$. For $\gamma_{\rm m} \simeq 10^3$–$10^4$, frequencies are in the infrared band for AGNs ($M \simeq 10^8\, M_\odot$). Observed LLAGNs with radiatively inefficient accretion flow indeed show a peak in infrared emission (Maoz 2007). The center also radiates infrared light, and according to observations (Genzel et al. 2003; Nishiyama et al. 2009), this emission comes from the rotating accretion disk. We suggest that this infrared emission is due to the proton synchrotron radiation from the disk. The emitted spectrum of radiation is locally a power law $F(\nu) \propto \nu^{-(\beta-1)/2}$ because the fast proton distribution function is a power law with index β. However, β changes in the disk since $\beta = \beta_{\rm H}(\rho/r_{\rm H})^{3/2}$. The integration over the disk gives the following dependence $F(\nu) \propto \nu^{-(\beta_{\rm H}-1)/2}/\ln(\nu)$, which is almost a power law, but corrected by the logarithmic function. We note that the observed power-law index of the infrared radiation from the Galactic center is -0.6 (Meyer et al. 2009), implying that index of the proton distribution is $\beta_{\rm H} \simeq 2.2$.

4. Very high energy radiation

In the Blandford-Znajek mechanism, almost all energy is transformed into protons, a jet, or a disk. Thus, it appears to be a barionic scenario for the very high energy (VHE) photon production. VHE photons are measured by Cherenkov telescopes and have energy in the TeV band. Energetic protons collide with the ambient matter and produce pions and then gamma quanta. Sources of the VHE radiation can be in the disk and the jet. We first calculate the VHE radiation from a disk. The spectrum of photons reproduces the spectrum of fast protons in Eq. (12), $\gamma \gg 1$, and is equal to

$$F(\gamma_{\rm ph}) = \frac{\sigma_{\rm E}}{8\pi m_{\rm p} c^3} \int 2^{2-\beta} (\beta-2) n B_\rho^2 \langle \tilde{u}_\phi^2 \rangle \gamma_{\rm ph}^{-\beta} {\rm d}V, \tag{19}$$

where n is the disk proton density, $\gamma_{\rm ph}$ is the photon energy $E_{\rm ph}$ in units of the proton rest energy, $\gamma_{\rm ph} = E_{\rm ph}/m_{\rm p}c^2$. The integration in Eq. (19) is over the disk volume. The quantities n, B_ρ, $\tilde{u}_\phi$, and β depend on the radial distance ρ as we have discussed. The result of the integration with logarithmic accuracy, $\ln\gamma_{\rm ph} > 1$, is

$$F(\gamma_{\rm ph}) = \frac{2^{2-\beta_{\rm H}}(\beta_{\rm H}-2)}{6\beta_{\rm H}} \sigma_{\rm E} c n_{\rm H} \frac{B^2 r_{\rm H}^3}{m_{\rm p} c^2} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \frac{\gamma_{\rm ph}^{-\beta_{\rm H}}}{\ln\gamma_{\rm ph}}\cdot \tag{20}$$

We obtain the total luminosity of VHE radiation from the disk by integrating the spectrum given by Eq. (19) over photon energies $m_p c^2 \gamma_{ph}$

$$L^{d}_{\rm VHE} = \frac{0.15}{8}\sigma_{\rm E} c n_{\rm H} B^2 r_{\rm H}^3 h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \,. \tag{21}$$

We see that the VHE luminosity is proportional to the energy of the magnetic field near the black hole, $L_{\rm VHE} \propto B^2 M^3$, and increases with the black hole mass. Substituting the condition in Eq. (2) in to the expression (21), we get

$$L^{\rm d}_{\rm VHE} \geq 1.1\times 10^{29} a^{-16} \frac{n_{\rm H}}{10^7 {\rm cm}^{-3}} \frac{M}{M_\odot} h \frac{\langle \tilde{u}_\phi^2 \rangle_{\rm H}}{c^2} \; {\rm erg/s}, \tag{22}$$

which implies that $L^{\rm d}_{\rm VHE} \simeq 10^{38}$ erg/s for $M \simeq 10^8\, M_\odot$. We note that the luminosity of VHE photons from the disk increases with the black hole mass in Eq. (22), while the bolometric luminosity of the proton synchrotron radiation in the disk represented by Eqs. (17) and (18) decreases with mass because of the different dependences of luminosities on the magnetic field B, $L_{\rm s} \propto B^4$, $L_{\rm VHE} \propto B^2$. The magnetic field must be stronger for low-mass black holes (see Eq. (2)).

Fast protons of the jet can also produce VHE photons. Their energy density on the light cylinder surface, $\rho = r_{\rm L} = c/\Omega_{\rm F} \simeq 4a^{-1} r_{\rm H}$, is equal to the energy density of the electromagnetic field near this surface $B_{\rm L}^2/4\pi = a^2 B^2/64\pi$ (Istomin 2010). Thus, the jet proton density on the light cylinder surface is $n_{\rm L} = a^2 B^2/64\pi m_p c^2 \gamma$. Moving further away the light surface, protons diminish in density $n_{\rm J}$ in line with the jet poloidal magnetic field to which they are frozen, $n_{\rm J} = n_{\rm L}(R_{\rm L}/\rho)^2$. The total luminosity of the jet in γ-rays is

$$L^{\rm J}_{\rm VHE} = 0.15 \sigma_{\rm E} c n m_p c^2 \gamma \int n_{\rm J} {\rm d}V,$$

where n is the density of the interstellar gas inside the jet and V is the jet volume. The integration provides a simple formula analogous to Eq. (21)

$$L^{\rm J}_{\rm VHE} = \frac{0.15}{2}\sigma_{\rm E} c n B^2 r_{\rm H}^3 \frac{l}{r_{\rm H}} \ln\left(\frac{r_{\rm J}}{r_{\rm L}}\right), \tag{23}$$

where l is the jet length and $r_{\rm J}$ is the outer radius of the jet at its base. Comparing Eqs. (21) and (23), we conclude that they are similar and that both contain the column density of the matter, which for the disk is $n_{\rm H} H$ and the jet is nl. We assume an angular resolution $\Delta\phi$. In the field of the central source, there is a contribution of the jet emission along its length $l = D\Delta\phi$, where D is the distance to the source. For a resolution $\Delta\phi < (r_{\rm H}/D)(n_{\rm H}/n)$, this means that the VHE luminosity of the disk dominates over the observed flux of the central source. In contrast, for a resolution $\Delta\phi > (r_{\rm H}/D)(n_{\rm H}/n)$ we will observe only the radiation of the jet. Using the condition given by Eq. (2), the expression in Eq. (23) becomes

$$L^{\rm J}_{\rm VHE} \geq 0.45\times 10^{22} a^{-16} \frac{n}{0.1\ {\rm cm}^{-3}} \frac{M}{M_\odot} \frac{l}{r_{\rm H}} \ln\left(\frac{r_{\rm J}}{r_{\rm L}}\right) \; {\rm erg/s}. \tag{24}$$

If the jet length l does not depend on the black-hole scale length $r_{\rm H}$, as it occurs over the accretion disk width H, then the VHE luminosity of the jet does not depend on the black hole mass and is defined only by the column density nl, $L^{\rm J}_{\rm VHE} \simeq 3\times 10^{35} nl/(1\ {\rm pc}\cdot 0.1\ {\rm cm}^{-3})$ erg/s.

5. Discussion

It seems that LLAGNs are good candidates in which to observe the Blandford-Znajek mechanism. Accreting matter onto a black hole in LLAGNs is a weak source of energy because of either the low accretion mass rate or low radiative efficiency. It may then be possible to observe a black hole operating like a dynamo machine. The accretion disk would play the role of a conductor through which the electric current would flow and the electric current would then follow a jet. For jet creation, a strong magnetic field near the black hole is needed. For AGNs of mass $10^8\, M_\odot$, this field should not be too high, $B \geq 3\times 10^3$ Gauss. The magnetic field could be accumulated during previous epochs of high accretion rate. The rotating black hole loses its rotation energy and angular momentum, which are both transmitted to the jet. Rotating with the black hole, the radial magnetic field transfers its rotation to the surrounding matter, which leeds to relativistic energies being attained on the light surface. The energy is mainly in protons, which form the relativistic jet. The accretion disk around the black hole can be observed in millimetre and infrared bands. To accrete matter, the disk must be turbulent (abnormal transport coefficients). The turbulent motion in the strong magnetic field generates a turbulent electric field, which accelerates disk ions. Electrons are not accelerated to relativistic energies because of their large synchrotron losses. Disk fast protons radiate synchrotron emission in the infrared range. The high proton energies should correspond to a high disk luminosity in the very high energy (VHE) photon range, and that a correlation exists between the infrared luminosity $L_{\rm s}$ of the disk and its VHE luminosity $L^{\rm d}_{\rm VHE}$. Using Eqs. (15) and (21), and excluding the unknown value of the magnetic field, we find the correlation

$$\log L^{\rm d}_{\rm VHE} = \frac{1}{2}\log L_{\rm s} + \frac{3}{2}\log\left(\frac{M}{M_\odot}\right) + \log\left(\frac{n_{\rm H}}{10^7\ {\rm cm}^{-3}}\right) + {\rm const.} \tag{25}$$

We can check the validity of this relation for low-luminosity AGNs with known luminosities $L_{\rm VHE}$ and $L_{\rm s}$, namely the Galactic center (Sgr A*), M 87, and Centaurus A. The nucleus of Centaurus A has a high bolometric luminosity $L_{\rm b} \simeq 1.3\times 10^{41}$ erg/s (Meisenheimer et al. 2009). At this luminosity, a high density of infrared photons of energy $\varepsilon_{\rm ph} \simeq 0.1$ eV in the central source prevents the free escape of VHE photons due to photon-photon collisions and the production of electron-positron pairs. The simple estimate $L_{\rm b} < \varepsilon_{\rm ph} c d/\sigma_{\rm T}$, where $\sigma_{\rm T}$ is the Thomson cross-section and d is the length scale of the central engine $d \simeq 10^2 r_{\rm H} = 1.5\times 10^{15}$cm for Centaurus A for which $M = 5\times 10^7\, M_\odot$, gives the condition $L_{\rm b} < 1.1\times 10^{37}$ erg/s at which the photon-photon annihilation is ineffective. The absorption of VHE quanta and the generation of e^+e^- pairs result in the re-radiation of VHE emission, as described by Stawarz et al. (2006), for the interaction of VHE radiation with the starlight radiation from stars of the host galaxy. Thus, for Centaurus A the observed luminosity $L_{\rm VHE}$ does not reflect the direct VHE radiation from the black hole vicinity.

For Sgr A*, we have $L_{\rm VHE} = 3\times 10^{34}$ erg/s (Aharonian et al. 2009a) and $L_{\rm s} = L_{\rm IR} = 10^{36}$ erg/s (Yuan et al. 2003), $M = 3.6\times 10^6\, M_\odot$, and for M 87 $L_{\rm VHE} = 3\times 10^{40}$ erg/s (Aharonian et al. 2006) and $L_{\rm s} = L_{\rm IR} = 10^{39}$ erg/s (Perlman et al. 2007), $M = 3\times 10^9\, M_\odot$. Substituting these values in to Eq. (25), we find

$$\log\left(\frac{n_{\rm H}({\rm SgrA}*)}{10^7\ {\rm cm}^{-3}}\right) + {\rm const.} = 6.64$$

for Sgr A*, and

$$\log\left(\frac{n_{\rm H}({\rm M\,87})}{10^7\ {\rm cm}^{-3}}\right) + {\rm const.} = 6.76$$

for M 87. These data do not contradict our model. To reach this conclusion, we must be sure that the observed VHE luminosity $L_{\rm VHE}$ comes from the disk, and not the jet. For the angular resolution of current VHE instruments, this requires that $n_{\rm H}/n > 10^{-3}(D/r_{\rm H})$. For M 87, the argument in favor of a disk origin of the VHE radiation is its short time variability, which excludes the large scale jet of 2 kpc length (Aharonian et al. 2006). On the basis of our estimate of the VHE luminosity of the total jet given by Eq. (21), $L^{\rm J}_{\rm VHE} \simeq 3\times10^{35}(l/1\ {\rm pc})$ erg/s, the jet is less luminous than the luminosity observed, while $L^{\rm d}_{\rm VHE} \geq 3\times10^{39}$ erg/s from Eq. (22). For the region of Sgr A*, the origin of the VHE emission is not quite clear (HESS collaboration 2010). One argument in favor of a disk origin of the VHE radiation is that the observed spectral index and the cut-off energy of VHE radiation (Aharonian et al. 2009a), $\beta \simeq -2.1$ and $E_{\rm c} \simeq 16$ TeV, are close to the values that follow from the observation of IR radiation from the disk, $\beta \simeq -2.2$, $\gamma_{\rm m} \simeq 10^4$, and $E_{\rm m} \simeq 10$ TeV. The argument of Aharonian et al. (2008) that VHE radiation from Sgr A* is likely from the jet rather than the disk is based on the assumptions that the X-ray emission originates in the disk and there is no time variability in the VHE flux during the X-ray flare. However, we discuss below another possible origin of X-ray radiation in the Blandford-Znajek mechanism, which is not in the disk.

Although Centaurus A is unsuitable for a comparison of the discussed model with observations, substituting its data into Eq. (25) provides values of the VHE luminosity not far from previous estimates, i.e.,

$$\log\left(\frac{n_{\rm H}({\rm CenA})}{10^7\ {\rm cm}^{-3}}\right) + {\rm const.} = 7.31,$$

where $L_{\rm VHE} = 2.6\times10^{39}$ erg/s (Aharonian 2009b), $L_{\rm IR} = 1.3\times10^{41}$ erg/s, and $M = 5\times10^7\ M_\odot$. This implies that the re-radiation of VHE emission does not strongly affect its power.

We should also consider whether the observed infrared radiation from Sgr A* and M 87 might originate in the disk. For Sgr A*, the short time variability of the NIR emission, which has a period of the order of 20 min, is strong evidence that it originates in the disk (Genzel et al. 2003; Nishiyama et al. 2009). The measured period corresponds to the rotation of a hot spot on the disk around the black hole at distances close to the black hole horizon. We recall that M 87 is a LINER. Apart from the thermal component of its mid-infrared emission, we also observe its power law synchrotron-like emission of similar intensity $\simeq10^{39}$ erg/s (Perlman et al. 2007). The thermal component of temperature $\simeq$50 K is the radiation of the dust around the central energy source. The power-law component is thought to be the radiation from the disk.

LLAGNs also radiate significant power in the X-ray band (Maoz 2007). Because the Thomson cross-section of the scattering of the electromagnetic radiation for protons is $(m_{\rm e}/m_{\rm p})^2$ times less than that for electrons the process of the inverse Compton scattering, which is important to models of standard AGN radiation, can not be applied to explain the X-ray radiation of LLAGNs in the scheme suggested here. However, a second component of fast protons exists in the jet formed by the Blandford-Znajek mechanism. These are protons accelerated in the disk, then ejected into the black hole magnetosphere that obtain additional energy while crossing the light cylinder surface. This two-step mechanism of proton acceleration up to very high energies was suggested by Istomin & Sol (2009). The Lorentz factor of these particles is $\gamma = (\gamma_{\rm d}\omega_{\rm cL}/\Omega_{\rm F})^{1/2}$. We recall that $\omega_{\rm cL}$ is the non-relativistic proton cyclotron frequency at the light surface, $\Omega_{\rm F}$ is the angular frequency of rotation of magnetic field lines, $\Omega_{\rm F} \simeq \Omega_{\rm H}/2$, and $\gamma_{\rm d}$ is the Lorentz factor of disk fast protons. The ratio $\omega_{\rm cL}/\Omega_{\rm F}$ is a very large number, $\gg\gamma_{\rm d}$, and the energy of these particles is much greater than the energy of fast particles in the disk. They radiate synchrotron emission in the region behind the light surface above the outer part of the disk near the base of the jet. The ratio of frequencies of the synchrotron radiation in both this region and the inner disk is $\nu_{\rm J}/\nu_{\rm d} = \bar{\omega}_{\rm c}\gamma^2/\omega_{\rm cH}\gamma_{\rm d}^2$, where $\bar{\omega}_{\rm c}$ is the proton cyclotron frequency at the jet base averaged over the volume from $r_{\rm L}$ to $r_{\rm J}$. In the jet behind the light surface, the poloidal magnetic field weakens like $\propto\rho^{-2}$, the toroidal field decreases more slowly, $\propto\rho^{-1}$, but initially at $\rho = r_{\rm L}$ B_ϕ is small, $B_\phi \ll B_\rho$. Because of this we can consider $B(\rho) = B_{\rm L}(r_{\rm L}/\rho)^2$ out to the outer jet radius $r_{\rm J}$ and $\bar{\omega}_{\rm c} = 2\omega_{\rm cL}(r_{\rm L}/r_{\rm J})^2\ln(r_{\rm J}/r_{\rm l})$. As a result, we obtain

$$\frac{\nu_{\rm J}}{\nu_{\rm d}} = 8a\frac{\omega_{\rm cH}r_{\rm H}}{c\gamma_{\rm d}}\left(\frac{r_{\rm L}}{r_{\rm J}}\right)^2\ln\frac{r_{\rm J}}{r_{\rm L}}\cdot$$

Using the expression in Eq. (1) for the ratio $\omega_{\rm cH}r_{\rm H}/c \simeq 3\times10^{10}$ and the estimates $\gamma_{\rm d} \simeq 10^4$ and $r_{\rm J}/r_{\rm L} \simeq 10$–$10^2$, we obtain $\nu_{\rm J}/\nu_{\rm d} \simeq 10^4$–$10^5$, which corresponds to $\nu_{\rm J}$ frequencies in the X-ray band when $\nu_{\rm d}$ is in the IR band. Unfortunately, we cannot estimate the X-ray luminosity from the jet base because we do not know the fraction of fast protons escaping the disk and being collected by the jet, which depends on the disk model. The variability of this X-ray emission is expected from plasma instabilities in the jet (Istomin 2010).

In conclusion, we can say that LLAGNs, or at least some of them, could be be extracting the energy from the black hole rotation by means of the Blandford-Znajek mechanism. The black hole spends almost all its energy on the jet production and proton acceleration. These LLAGNs are probably sources of high-energy cosmic rays. Their VHE luminosity should reflect the intensive process of proton acceleration. Our scenario predicts that the value of $L_{\rm VHE}$ increases with the black hole mass as $M^{3/2}$ and with the infrared luminosity of the disk as $L_{\rm IR}^{1/2}$ (see Eq. (25)). The discovery of new bright VHE sources from LLAGNs could confirm our hypotheses.

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