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Relativistic jets in active galactic nuclei: time variability

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ABSTRACT

We show that a jet in an active galactic nucleus (AGN) is produced by a rotating massive black hole. The jet contains relativistic protons accelerated by the corotating radial magnetic field, which is brought to the black hole horizon by the accreting disc matter. The Lorentz factor of protons at the light cylinder surface can gain a value of γ as high as $\simeq 10^5$. The main proton energy is in the azimuthal rotation, and only a small part is in the radial motion. The jet matter, flowing out slowly behind the light cylinder, $v_r = c\gamma^{-1/2}$, is accelerated moving in the expanding poloidal magnetic field (magnetic nozzle) and it obtains a longitudinal Lorentz factor of $\Gamma \simeq (B/B_L)^{-1/2} \le 10$ (where B_L is the value of the magnetic field at the light cylinder surface, $B_L \simeq 10^2$ G). The observed short time variability from some AGNs, $\tau \simeq 300$ s, can be explained by the excitation of the collective ion plasma waves near the jet base.

Key words: acceleration of particles – galaxies: active.

1 INTRODUCTION

One of the basic questions arising in the physics of active galactic nuclei (AGNs) is where and why a jet originates. It is considered that a jet is formed in the neighbourhood of a massive black hole in the centre of an active galaxy, and it is based on an accretion disc rotating around a black hole. Why the accretion is accompanied by the jet is a puzzle in many cases. What generates a jet: the accreting disc flow or the hole itself? It is clear that the jet matter consists of the disc matter and may partially be of the electronpositron plasma generated nearby the black hole horizon (Beskin, Istomin & Pariev 1992). The jet possesses a relativistic velocity up to the bulk Lorentz factor of the order of 10, $\Gamma \simeq 10$. Such high velocities of the accreting matter exist only near the last stable particle orbit around the black hole. Jet models are based mainly on hydrodynamics (HD) flows, or on magnetohydrodynamics (MHD) flows, propagating from the central machine through a surrounding matter. The theory of MHD flows is well described in several monographs (see, for example, the monograph by Beskin 2010). In HD and MHD models, the inhomogeneities of the jet flow and the time variability are the result of the time-dependent ejection of the matter from the black hole vicinity. It is suggested that different velocities of the flow in a jet create the shock wave in which electrons are accelerated and radiate the synchrotron emission (Raga, Canto & Biro 1993; Mendoza et al. 2009). For the jet to be stable, it must possess large angular momentum. The rotating massive black hole has an angular momentum much larger than the angular momentum of the accretion disc. So, it is natural to consider that the angular

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momentum and the energy are transformed to the jet from the black hole. This mechanism is known as the Blandford–Znajek mechanism (Blandford & Znajek 1977) of extracting energy and angular momentum from the black hole rotation. In this mechanism, the jet has not only relativistic bulk velocity but also relativistic internal energy. As a result, we have a kinetic jet, which can be treated using kinetic theory rather than HD and MHD approximations.

2 KINETIC JET

The MHD simulation of conducting plasma motion near the black hole horizon shows that the magnetic field generated in this region is not expelled by the black hole. In contrast, it obtains a radial structure with the magnetic flux going in and going out of the black hole horizon (Komissarov & McKinney 2007). Istomin & Sol (2009) have shown that for the transition of the mass and the angular momentum of the disc matter to the hole, the disc must push out the component of the magnetic field that is vertical with respect to the disc. This is produced by an electric current flowing in a disc. The rotating black hole creates the voltage U by the mechanism of the unipolar inductor, $U = f \Omega^{\rm H} / 2\pi c$ (Landau & Lifshitz 1984). f is the flux of the poloidal magnetic field on the black hole horizon and $\Omega^{\rm H}$ is the frequency of rotation of the black hole. This voltage U supports the electric current I flowing through the disc. It is the part of the closed current loop arising in the jet-accretion disc-black hole system (see Fig. 1). The radial current in the disc $I_{\rho} = -I$ is produced by the radial density current j_{ρ} , which is the component of the disc current j having azimuthal j_{ϕ} and radial j_{ρ} components. The main current is the azimuthal current j_{ϕ} resulting from the spiral motion of disc particles $(v_{\rho} \ll v_{\phi}), j_{\rho} = \alpha j_{\phi}, \alpha \simeq 10^{-2} \ll 1.$ The power of the central engine is UI. In high-conducting media, the magnetic field begins to rotate with the angular velocity $\Omega^{\rm F}$,

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Figure 1. Schematic diagram of the magnetic field lines and electric currents in the vicinity of the massive black hole. The electric current *I* flows in the inner and outer parts of the jet and in the disc. The disc electric density currents are j_{ρ} and j_{ϕ} ; the azimuthal current j_{ϕ} is larger than the radial current $j_{\rho} = \alpha j_{\phi}$, $\alpha \simeq 10^{-2}$. The light surface is shown as a dashed line.

similar to the angular velocity of the black hole rotation $\Omega^{\rm H}, \Omega^{\rm F} \simeq \Omega^{\rm H}$. Also, the plasma in the black hole magnetosphere begins to rotate with the same velocity $\Omega^{\rm F}$. This means that magnetosphere plasma particles moving along the radial magnetic field will be accelerated by the centrifugal mechanism. This process has been described by Istomin & Sol (2009). Protons crossing the light cylinder surface $R_{\rm L} = c/\Omega^{\rm F}$ gain the energy $m_{\rm p}c^2\gamma$, $\gamma = (\gamma_i\gamma_0)^{1/2}$. Here, the value of γ_i is the initial particle Lorentz factor, and $\gamma_0 = \omega_{\rm c}(\rho = R_{\rm L})/\Omega^{\rm F}$, where $\omega_{\rm c}$ is the ion non-relativistic cyclotron frequency of the particle, $\omega_{\rm c} = e B_{\rm L} / m_{\rm p} c$, in the magnetic field $B_{\rm L}$ at the light cylinder surface. For the magnetic field of $B_{\rm L} \simeq 10^2$ G and the frequency of the black hole rotation of $\Omega^{\rm H} \simeq 10^{-5} \, {\rm s}^{-1}$, the value of γ_0 can be as high as $\gamma_0 \simeq 10^{11}$. The value of γ_0 is the maximum Lorentz factor that can be obtained in this centrifugal mechanism. Electrons can obtain much less energy because of large synchrotron losses. Thus, the black hole rotation in the presence of a magnetic field can accelerate the magnetospheric matter (protons) up to high energy $\gamma = \gamma_0^{1/2} \simeq 10^5 (\gamma_i \simeq 1)$. Almost all the energy is in the azimuthal motion $p_{\phi} = m_{\rm p} c \gamma$, and only a small part is in the radial and vertical components of the particle momentum $p_{\rho} \simeq p_z \simeq m_{\rm p} c \gamma^{1/2}$ (Istomin & Sol 2009). We see that the rotating black hole creates the rotating energetic flux of the matter slowly flowing out from the light cylinder surface with the velocity $v_{\rho} = p_{\rho}/m_{\rm p}\gamma \simeq c\gamma^{-1/2} \ll c$. We note that the energy obtained by a charged particle is the work of the vertical electric field E_z under a particle, and a particle must elevate on the high $z = \gamma m_{\rm p} c^2 / e E_z$, which is just its cyclotron radius $r_{\rm c} = m_{\rm p}c^2\gamma/eB_{\rm L}$ on the light surface $(E_z = B_{\rm L})$. The maximum density of this hot plasma $n_{\rm L}$ is determined by the possibility of the magnetic field rotating particles near the light cylinder surface $R_{\rm L}$

with the velocity $v = c(1 - \gamma^{-2})^{1/2}$:

$$n_{\rm L}\gamma m_{\rm p}c^2 = \frac{B_{\rm L}^2}{8\pi} + \frac{E_{\rm L}^2}{8\pi} = \frac{B_{\rm L}^2}{4\pi}.$$

When the black hole magnetosphere is filled by a more dense plasma, $n_L\gamma_0^{1/2} > B_L^2/4\pi m_p c^2$, then the proton energy $\gamma m_p c^2$ decreases, but the energy density remains constant, $n_L\gamma m_p c^2 = B_L^2/4\pi$. Protons, obtaining the energy, move in the vertical direction and create the charge separation. The vertical electric field is screened on the distance of the Debye length, $r_d = (m_p c^2 \gamma / 4\pi n e^2)^{1/2}$, and the particle energy on the light surface becomes $m_p c^2 \gamma = e B_L r_d$. This relation gives the condition $n_L \gamma m_p c^2 = B_L^2/4\pi$.

This energetic plasma, accelerated near the light cylinder surface, flows out through the stationary fields $B \simeq B_{\rm L}$ and $E \simeq E_{\rm L}$ with small radial velocity, $v_{\rho} \simeq c\gamma^{-1/2}$. Thus, the flux of the kinetic energy $L_{\rm k}$ flowing out from the black hole vicinity is

$$\begin{split} L_{\rm k} &= 2\pi n_{\rm L} \gamma^{1/2} m_{\rm p} c^3 R_{\rm L}^2 = \frac{1}{2} c R_{\rm L}^2 B_{\rm L}^2 \gamma^{-1/2} \\ &= \frac{1}{2} c^3 \left(\frac{B_{\rm L}}{\Omega^{\rm F}} \right)^2 \left(\frac{\omega_{\rm c}}{\Omega^{\rm F}} \right)^{-1/4}, \\ L_{\rm k} &\simeq 2 \times 10^{42} \left(\frac{B_{\rm L}}{\Omega^{\rm F}} \frac{10^{-5} \, {\rm s}^{-1}}{10^2 \, {\rm G}} \right)^{7/4} \, {\rm erg \, s}^{-1}. \end{split}$$

This expression is similar to the usual radio pulsar loss $L \simeq c R_{\rm L}^2 B_{\rm L}^2$ used for the estimation of a neutron star magnetic field. Such a formula for the pulsar energy loss is based on the suggestion that neutron stars spin down by the magnetic dipole radiation. For this radiation, the radial velocity equals the speed of light c. In our case, the proton radial velocity is smaller by a factor of $\gamma^{-1/2}$. A similar expression takes place in the Blandford-Znajek mechanism, $L_{\rm BZ} \propto c r_{\rm g}^2 B_{\rm H}^2$. The structure of such expressions does not mean that the stationary electromagnetic field flows out from an energy source. It means that the energy flowing out has a density equal to the energy density of the electromagnetic field in a source. However, the radial flux of the energy beyond the light surface is mainly the Poynting flux $c E_{Lz} B_{L\phi} / 4\pi = \alpha c B_L^2 / 4\pi = \alpha \gamma^{1/2} (n_L \gamma m_p c^2 v_o), \alpha \gamma^{1/2} > 1$, which carries energy along the jet to its outer part. The rotating energetic plasma, expanding from the light cylinder surface $R_{\rm L}$ up to the jet radius $R_{\rm J}$, will create the jet.

3 PLASMA EXPANSION: MAGNETIC NOZZLE

The ejected plasma consists of energetic protons having almost the same Lorentz factor $\gamma \gg 1$; they are magnetized. Their maximum cyclotron radius in the magnetic field at the light cylinder is $r_{\rm c} = \gamma c / \omega_{\rm c} \simeq 10^9$ cm. This is much less than the light cylinder radius $R_{\rm L}$, $R_{\rm L} > r_{\rm g}$, where $r_{\rm g}$ is the gravitational radius of the massive black hole. Initially drifting in the azimuthal direction, they decrease their azimuthal momentum as a result of the conservation of the generalized angular momentum, $\rho p_{\phi} + ef/c = \text{const}$ (Istomin & Sol 2009). Here, f is the flux of the axisymmetric stationary magnetic field. For averaged motion along the magnetic field lines, f is constant and $p_{\phi} = cm_{\rm p}\gamma R_{\rm L}/\rho$, $\rho > R_{\rm L}$. Then protons begin to rotate in the radial magnetic field with a momentum $p_{\perp} = m_{\rm p} c \gamma$ which is perpendicular to the magnetic field. They move slowly along the magnetic field with the velocity $v_{\parallel} = c\gamma^{-1/2}$. As for electrons, these are 'cold' and play the role of neutralizator of the proton electric charge.

Near the light cylinder surface, the main component of the magnetic field is the radial field. It is α^{-1} times larger than the toroidal magnetic field B_{ϕ} created by the jet electric current *I* (Istomin & Sol 2009; see Fig. 1). Here the value of $\alpha \ll 1$ is the ratio of the radial electric current in a disc j_{ρ} to the toroidal current j_{ϕ} , $\alpha = j_{\rho}/j_{\phi}$. Like the component of the poloidal magnetic field $B_{\rm p}$, the radial magnetic field decreases with the radial distance ρ , $B_{\rm p} \propto \rho^{-2}$. Then, approaching the jet radius $R_{\rm J}$, it becomes parallel to the jet axis *z*. The toroidal magnetic field falls down slowly, $B_{\phi} \propto \rho^{-1}$, and becomes comparable with the poloidal magnetic field, $B_{\rm p}$, at the distance $\rho = \alpha^{-1}R_{\rm L}$. This radius can be considered as the jet radius, $R_{\rm J} \simeq \alpha^{-1}R_{\rm L}$.

Moving along the falling magnetic field, the proton gains longitudinal momentum as a result of the conservation of the adiabatic invariant $J = p_{\perp}^2/\omega_c = \text{const.}$ Its longitudinal momentum can become $p_{\parallel} = m_p c \gamma (1 - B/B_L)^{1/2}$, and all the plasma obtains noticeable longitudinal velocity $u_{\parallel} = c(1 - B/B_L)^{1/2}$. However, before this occurs, the distribution function of protons $f(p_{\perp}, p_{\parallel})$, which initially was strongly anisotropic, $p_{\perp} \gg p_{\parallel}$, becomes isotropic in a short time. Let us investigate this.

The anisotropic distribution of the charged particle is unstable with respect to the excitation of electromagnetic waves, which turn the particle momentum in the phase space and make their distribution function f = f(p) isotropic. Electrostatic waves usually have a larger rate of growth than electromagnetic waves. The dispersion properties of electrostatic low-frequency waves in relativistic plasma are described by the plasma longitudinal permittivity ϵ_1 :

$$\epsilon_{\mathrm{I}}(\omega, \mathbf{k}) = 1 + \frac{4\pi e^2}{k^2} \sum_{s=-\infty}^{\infty} \int \frac{J_{\mathrm{s}}^2 (k_{\perp} p_{\perp} / m_{\mathrm{p}} \omega_{\mathrm{c}})}{\omega - s \omega_{\mathrm{c}} / \gamma - k_{\parallel} v_{\parallel}} \\ \times \left(k_{\parallel} \frac{\partial f}{\partial p_{\parallel}} + \frac{s \omega_{\mathrm{c}} / \gamma}{v_{\perp}} \frac{\partial f}{\partial p_{\perp}} \right) \,\mathrm{d}\mathbf{p}.$$
(1)

Here, k_{\parallel} and k_{\perp} are the components of the wave vector k, which are parallel and perpendicular to the magnetic field, and J_s is the Bessel function of the first order. Electrons are strongly magnetized and their input to the permittivity for low frequencies is of the order of $\delta\epsilon_1 \simeq \omega_{pe}^2/\omega_{ce}^2 = (m_e/m_p)\omega_p^2/\omega_c^2 \simeq m_e/m_p$ and can be neglected. Because the initial distribution function is strongly anisotropic $(p_{\parallel} \ll p_{\perp})$ we can put $p_{\parallel} = 0$ in equation (1) and consider the ion distribution function to be monochromatic, $f(p_{\parallel}, p_{\perp}) = n\delta(p_{\parallel})\delta(p_{\perp} - m_pc\gamma)/2\pi p_{\perp}$. Here, $\delta(x)$ is the Dirac delta function. Performing the partial integration in equation (1) and taking into account the recurrence relations for Bessel functions $2sJ_s(z) = z(J_{s-1} + J_{s+1})$, $2dJ_s(z)/dz = J_{s-1} - J_{s+1}$ we find

$$\epsilon_{1} = 1 - \frac{\omega_{p}^{2}}{\gamma k^{2} c^{2}} \sum_{s=-\infty}^{\infty} J_{s}^{2} (c k_{\perp} \gamma / \omega_{c}) \left[\frac{k_{\parallel}^{2} c^{2}}{(\omega - s \omega_{c} / \gamma)^{2}} + \frac{k_{\perp}^{2} c^{2}}{(\omega - s \omega_{c} / \gamma)^{2} - \omega_{c}^{2} / \gamma^{2}} - \frac{s^{2} \omega_{c}^{2} / \gamma^{2}}{(\omega - s \omega_{c} / \gamma)^{2}} \right].$$

$$(2)$$

The quantity ω_p is the proton plasma frequency, $\omega_p^2 = 4\pi ne^2/m_p$. If the argument of the Bessel function is small, $ck_\perp < \omega_c/\gamma$, then the main input in equation (2) gives the term s = 0, and the dispersion equation $\epsilon_1 = 0$ describes ion plasma waves:

$$1 - \frac{\omega_{\rm p}^2}{\gamma k^2} \left(\frac{k_{\parallel}^2}{\omega^2} + \frac{k_{\perp}^2}{\omega^2 - \omega_{\rm c}^2/\gamma^2} \right) = 0.$$
⁽³⁾

The frequency for the low-frequency branch varies from the frequency $\min(\omega_p/\gamma^{1/2}, \omega_c/\gamma)$ at $k_\perp = 0$ to $\omega = 0$ at $k_\parallel = 0$. We note that $\omega_p^2/\gamma = \omega_c^2/\gamma^2$ at the light cylinder surface $\rho = R_L$. At $\rho > R_L$, the plasma density decreases faster than the magnetic field because of ion acceleration along the magnetic field, $\omega_c^2/\gamma^2 > \omega_p^2/\gamma$. Apart from the plasma wave in a magnetized plasma, there are also ion cyclotron waves, or ion Bernstein waves, whose frequencies are close to the harmonics of the ion cyclotron frequency, $\omega \simeq l\omega_c/\gamma$. Substituting into equation (2) the definition $\omega = l\omega_c/\gamma + \delta\omega, \delta\omega \ll \omega$, we notice that the main contribution to the first and third terms within the square brackets in equation (2) gives s = l, and these terms will be proportional to $\delta\omega^{-2}$. As for the second term within the square brackets in equation (2), this is maximum at $s = l \pm 1$ and is proportional to $\delta\omega^{-1}$. Thus, we obtain the quadratic equation for $\delta\omega$. The roots of this equation are

$$\delta\omega = \frac{\omega_{\rm p}^2}{\omega_{\rm c}} \frac{k_{\perp}^2}{k^2} \frac{lJ_{\rm l}J_{\rm l}'}{z} \pm \frac{\omega_{\rm p}^2}{\omega_{\rm c}} |J_{\rm l}| \\ \times \left[\frac{k_{\perp}^4}{k^4} \left(\frac{lJ_{\rm l}'}{z} \right)^2 + \frac{\omega_{\rm c}^2}{\gamma \omega_{\rm p}^2 k^2 c^2} \left(k_{\parallel}^2 c^2 - l^2 \omega_{\rm c}^2 / \gamma^2 \right) \right]^{1/2}.$$
(4)

Here, z is the argument of the Bessel function, $z = ck_{\perp}\gamma/\omega_c$, where the prime denotes the derivative of the Bessel function over z, $J'_1 = dJ_1/dz$. From equation (4), we can see that at

$$\begin{split} k_{\parallel} &< l\omega_{\rm c}/c\gamma, \\ \frac{k_{\perp}^2}{k^2} &< \frac{\gamma}{(lJ_{\rm I}')^2} \frac{l^2 \omega_{\rm c}^2/\gamma^2 - k_{\parallel}^2 c^2}{\omega_{\rm p}^2} \end{split}$$

there appears a positive imaginary part of the frequency, $\delta\omega \simeq i l \omega_c \omega_p |J_l| / k c \gamma^{3/2}$. Thus, the strong anisotropic distribution of ions $(p_{\perp} \gg p_{\parallel})$ makes the excitation of ion cyclotron waves (Mikhailovskii 1975). These waves, satisfying the condition $|\delta \omega|/k_{\parallel} < v_{\parallel} = c \gamma^{-1/2}$, are absorbed by Landau damping at the resonance $\omega - l\omega_c/\gamma - k_{\parallel}v_{\parallel} = 0$ (see equation 1), and particles obtain longitudinal momentum. During the several times $\delta \omega^{-1}$, the ion distribution function becomes isotropic. The increment $Im(\delta\omega)$ is high; its maximum is achieved at $z \approx l, l|J_1| \approx 1$, and is of the order of $Im(\delta\omega) \approx \omega_p / l\gamma^{1/2} (k \simeq k_\perp)$. Near the light cylinder surface, the plasma frequency is of $\omega_{\rm p} \approx \omega_{\rm c}/\gamma^{1/2}$, and we have $Im(\delta\omega) \approx \omega_{\rm c}/l\gamma = \omega/l^2$. For typical values of the proton cyclotron frequency at the light cylinder surface, $\omega_c \approx 10^6 \text{ s}^{-1}$, and we have $Im(\delta\omega) \simeq 10 \text{ s}^{-1}$. This is much larger than the inverted time of the particle travelling through the jet, $t^{-1} = c/R_{\rm J} < c/R_{\rm J} <$ $c/r_{\rm g} \approx 10^{-5} \, {\rm s}^{-1}$. The time of the anisotropic instability development, $Im^{-1}(\delta\omega) \approx 10^{-1}$ s, is also much less than the observed time of variability ≈ 300 s. Thus, we can consider the ion distribution function in a jet to be isotropic. As for the particle distribution function with $p_{\parallel} > p_{\perp}$, this is known to be unstable also because of the so-called 'fan' instability (Parail & Pogutse 1986). The increment of this instability is of the order of $Im(\omega) \simeq \omega_{\rm p}^2 \omega / \gamma \omega_{\rm c} kc$. We see that in any case the anisotropic distribution function of ions becomes isotropic, f = f(p), during a short time.

However, a charged particle in the expanding magnetic field experiences longitudinal acceleration. The force F, directed along the magnetic field, acts on a proton,

$$\boldsymbol{F} = -\mu \nabla_{\parallel} \boldsymbol{B}.$$

Here, *B* is the magnitude of the magnetic field, B = |B|, and μ is the magnetic moment of a particle, $\mu = p_{\perp}^2/2m_p\gamma B$. Thus, we can write

$$\frac{\mathrm{d}p_{\parallel}^2}{\mathrm{d}t} = -\frac{p_{\perp}^2}{B}\frac{\mathrm{d}B}{\mathrm{d}t}.$$
(5)

If the particle energy is conserved, $dp_{\parallel}^2/dt = -dp_{\perp}^2/dt$, then equation (5) gives the adiabatic invariant conservation, $J = p_{\perp}^2/B =$ const. However, in our case the energy is not conserved because

of the collective anisotropic instability. We have shown that in the reference frame moving with the longitudinal Lorentz factor Γ the ion distribution function is isotropic, $\langle p'_{\perp}^2 \rangle = 2 \langle p'_{\parallel}^2 \rangle$, $\langle \epsilon'^2 \rangle = m_p^2 c^4 + 3 \langle p'_{\parallel}^2 \rangle c^2$. The angular brackets denote averaging over the particle distribution function, and the prime denotes values in the moving frame. The connection between quantities in the laboratory and those in the moving frame is

$$p_{\parallel} = \Gamma p'_{\parallel} + (\Gamma^2 - 1)^{1/2} \epsilon' / c, \ p_{\perp} = p'_{\perp},$$

 $\epsilon = \Gamma \epsilon' + (\Gamma^2 - 1)^{1/2} p'_{\parallel} c.$

Averaging equation (5),

$$\frac{\mathrm{d}\langle p_{\parallel}^2\rangle}{\mathrm{d}t} = -\langle p_{\perp}^2\rangle \frac{\mathrm{d}B/\mathrm{d}t}{B}$$

and substituting this into relations

$$\langle p_{\parallel}^2 \rangle = (4\Gamma^2/3 - 1)\langle \epsilon'^2 \rangle/c^2 - \Gamma^2 m_p^2 c^2/3,$$

and
 $\langle p_{\perp}^2 \rangle = 2\langle \epsilon'^2 \rangle/3c^2 - 2m_p^2 c^2/3,$

we can write the equation describing the evolution of the Lorentz factor Γ :

$$\frac{\mathrm{d}\Gamma^2}{\mathrm{d}t} \left(4\langle \epsilon'^2 \rangle - m_{\rm p}^2 c^4 \right) = -2\frac{1}{B} \frac{\mathrm{d}B}{\mathrm{d}t} \left(\langle \epsilon'^2 \rangle - m_{\rm p}^2 c^4 \right) -\frac{\mathrm{d}\langle \epsilon'^2 \rangle}{\mathrm{d}t} (4\Gamma^2 - 3).$$
(6)

The expanding plasma decreases its density and cools. Because of the strong relativity, the equation of state for protons is $\langle \epsilon' \rangle \propto n'^{1/3}$ or $\langle \epsilon'^2 \rangle \propto n'^{2/3}$. However, $n' = n/\Gamma$, where *n* is the proton density in the laboratory frame. From the particle flux conservation, it follows that $n \propto B/V_{\parallel}$, where V_{\parallel} is the plasma velocity along the magnetic field, $V_{\parallel} = c(1 - \Gamma^{-2})^{1/2}$. Thus, we have

$$\langle \epsilon'^2 \rangle = \langle \epsilon'^2 \rangle_0 \left(\frac{B}{B_0} \right)^{2/3} \left(\frac{\Gamma_0^2 - 1}{\Gamma^2 - 1} \right)^{1/3}.$$

Here, the values B_0 , $\langle \epsilon'^2 \rangle_0$ and Γ_0 are the initial values of B, $\langle \epsilon'^2 \rangle$ and Γ , $B_0 = B_L$, $\langle \epsilon'^2 \rangle_0 = \gamma^2 m_p^2 c^4$ and $\Gamma_0 = (1 - \gamma^{-1})^{-1/2}$. Substituting this expression into equation (6), we find

$$\frac{\mathrm{d}\Gamma^2}{\mathrm{d}t} = -\frac{1}{B} \frac{\mathrm{d}B}{\mathrm{d}t} \frac{8\Gamma^2(\Gamma^2 - 1)^{-1/3}\gamma^{5/3}(B/B_0)^{2/3} - 6}{(8\Gamma^2 - 9)(\Gamma^2 - 1)^{-4/3}\gamma^{5/3}(B/B_0)^{2/3} - 3}.$$
 (7)

Taking into account that $\gamma \gg 1$, we can simplify this equation and finally obtain

$$\Gamma^2 \left(\frac{\Gamma^2}{\Gamma^2 - 1}\right)^{1/8} = \gamma^{1/8} \left(\frac{B}{B_0}\right)^{-1}.$$
(8)

For $\Gamma > 2$, this expression gives

$$\Gamma = \gamma^{1/16} \left(\frac{B}{B_0}\right)^{-1/2}.$$
(9)

This is a remarkable result. The Lorentz factor of the longitudinal motion of the jet Γ at its base is not sensitive to the initial proton energy γ , and is determined by the poloidal magnetic field downfall, $\Gamma \simeq (B/B_0)^{-1/2} = R_J/R_L$. Physically, the reason for the longitudinal acceleration of the jet is the expanding magnetic field. In astrophysics, this acceleration was first discussed by Parker (1958) for the solar wind acceleration. For jets, it was applied by Blandford & Payne (1982). In the laboratory, this phenomenon is used in plasma-based propulsions, where a plasma flows along expanding

magnetic field lines (see, for example, Arefiev & Breizman 2005). This is similar to the gas acceleration in a nozzle and is called the 'magnetic nozzle'.

The internal plasma energy $\langle \epsilon' \rangle$ and the density n' fall with the decrease of the magnetic field:

$$\begin{split} \langle \epsilon' \rangle &= m_{\rm p} c^2 \gamma \left[\gamma^{-3/16} \left(\frac{B}{B_0} \right)^{1/2} \right] \\ n' &= n_{\rm L} \gamma^{-9/16} \left(\frac{B}{B_0} \right)^{3/2}. \end{split}$$

The density of the internal energy $n'\langle\epsilon'\rangle$ becomes much less than the density of the magnetic energy $B^2/8$:

$$n'\langle\epsilon'\rangle = n_{\rm L}m_{\rm p}c^2\gamma \gamma^{-3/4} \left(\frac{B}{B_0}\right)^2 = 2\frac{B^2}{8\pi}\gamma^{-3/4}$$

However, the energy density in the laboratory frame $\langle \epsilon \rangle = \Gamma \langle \epsilon' \rangle$ is only several times less than its initial value:

$$\langle \epsilon \rangle = \gamma^{-1/8} m_{\rm p} c^2 \gamma.$$

Taking into account the plasma density decrease, $n = n_1 \gamma^{-1/2} (B/B_0)$, we find the jet kinetic luminosity at its base

$$L_{\rm K} = \pi n_{\rm L} \gamma^{3/8} m_{\rm p} c^3 R_J^2 \left(\frac{B}{B_0}\right)$$

This expression can be used for the estimation of the magnetic field at the light cylinder surface if we know the jet kinetic energy:

$$B_{\rm L} = \Omega^{\rm F} \left[\frac{2L_{\rm K}}{c^3} \left(\frac{e}{m_{\rm p}c} \right)^{5/16} \right]^{16/27}$$

= 6 × 10² $\left(\frac{\Omega^{\rm F}}{10^{-5} \, {\rm s}^{-1}} \right) \left(\frac{L_{\rm K}}{10^{43} \, {\rm erg} \, {\rm s}^{-1}} \right)^{16/27} \, {\rm G}.$

4 TIME VARIABILITY

The TeV gamma-ray flux observed from blazars with the HESS and MAGIC telescopes has a short time variability $\tau \simeq 150-300$ s (Aharonian et al. 2007; Albert et al. 2007). If this radiation is assumed to generate in the jet base, then the time τ restricts the scale of the causally coupled generated region l by the relation, $l < \tau c \Gamma$. Here, the value of Γ is the Lorentz factor of the jet at its base. As we can see from equation (9), this factor is not very large, $\Gamma \simeq (B/B_L)^{-1/2} \simeq R_J/R_L \le 10$. So, $l \le 10^{14}$ cm is less than the jet radius and even the gravitational radius of the massive black hole with the mass of $10^9 \,\mathrm{M}_{\odot}$. Thus, the emission region is much less than the jet radius, or for larger scales it possesses a unified characteristic time. First of all, these times are proton characteristic time-scales. They are times that are inversely proportional to the ion plasma frequency, $\tau_p = 2\pi/(\omega_p/\langle \gamma \rangle^{1/2})$, and the ion cyclotron frequency, $\tau_c = 2\pi/(\omega_c/\langle \gamma \rangle)$. The value of $\langle \gamma \rangle$ is the mean ion Lorentz factor. Different plasma processes can reveal these timescales. First of all, they are plasma instabilities with large increments. Let us point out one of these; it is a beam instability. The burst of a turbulent motion in the accretion disc, or burst activity in a black hole magnetosphere, for example, as a result of the magnetic reconnection results in a particle acceleration (Dermer, Miller & Li 1996; Birk et al. 1999; Istomin & Sol 2009). These fast protons gain additional energy crossing the light cylinder surface by a centrifugal mechanism. Let us call these particles the ion beam. The energy of the beam particles is much larger than the energy of protons in the jet, $\gamma_b = (\gamma_i \gamma_0)^{1/2} = \gamma_i^{1/2} \gamma$, $\gamma_i \gg 1$. Their cyclotron radius is much larger than the main jet proton cyclotron radius, $r_b = c\gamma_b/\omega_c \gg c\gamma/\omega_c$. For this reason, for the jet plasma the beam can be considered as not magnetized. The input of non-magnetized beam protons to the longitudinal dielectric permittivity $\delta \epsilon_b$ is (compare with equation 1)

$$\delta \epsilon_{\rm b} = \frac{4\pi e^2}{k^2} \int \frac{k}{(\omega - kv)} \frac{\partial f_{\rm b}}{\partial p} \, \mathrm{d}p.$$

Performing partial integration and using the relation

$$\frac{\mathrm{d}v_{\alpha}}{\mathrm{d}p_{\beta}} = \frac{\left(\delta_{\alpha\beta} - p_{\alpha}p_{\beta}/m_{\mathrm{p}}^{2}c^{2}\gamma^{2}\right)}{m_{\mathrm{p}}\gamma},$$

we obtain

$$\delta\epsilon_{\rm b} = -\frac{\omega_{\rm b}^2}{(\omega' - k'c\cos\theta')^2} \left(\frac{\cos^2\theta'}{\langle\gamma_{\rm b}^{\prime-3}\rangle^{-1}} + \frac{\sin^2\theta'}{\langle\gamma_{\rm b}^{\prime-1}\rangle^{-1}}\right). \tag{10}$$

Equation (10) is written in the frame moving with the jet along the magnetic field with the Lorentz factor Γ (index prime). The beam plasma frequency ω_b is $(4\pi n'_b e^2/m_p)^{1/2}$, θ' is the angle between the wave vector \mathbf{k}' and the beam velocity, and the brackets $\langle ... \rangle$ denote averaging over the beam proton distribution. In the laboratory frame, the beam is injected in the azimuthal direction (i.e. in the direction perpendicular to the jet axis and to the poloidal magnetic field). This means that the angle between the magnetic field and the beam velocity in the comoving frame α' is $\cos \alpha' = -(\Gamma^2 - 1)^{1/2}/\Gamma$, $\cos \alpha = 0$. We see that the beam velocity is almost antiparallel to the magnetic field, $\cos \alpha' \simeq -1$. So, we can write the dispersion equation for waves exited in the system 'jet + beam' in the form:

$$1 - \frac{\omega_{\rm p}^{\prime 2}}{\gamma^{\prime}} \left(\frac{\cos^2 \theta^{\prime}}{\omega^{\prime 2}} + \frac{\sin^2 \theta^{\prime}}{\omega^{\prime 2} - \omega_{\rm c}^2 / \gamma^{\prime 2}} \right) - \frac{\omega_{\rm b}^2}{(\omega^{\prime} - k^{\prime} c \cos \theta^{\prime})^2} \\ \times \left(\frac{\cos^2 \theta^{\prime}}{\left\langle \gamma_{\rm b}^{\prime - 3} \right\rangle^{-1}} + \frac{\sin^2 \theta^{\prime}}{\left\langle \gamma_{\rm b}^{\prime - 1} \right\rangle^{-1}} \right) = 0.$$
(11)

Here, γ' is the Lorentz factor of jet protons in the moving frame, $\gamma' = \langle \epsilon' \rangle / m_p c^2 = \gamma^{7/8} \Gamma^{-1}$ and $\gamma'_b = \gamma_b \Gamma$. The values of the effective plasma and cyclotron frequencies in the jet are

$$\begin{split} \frac{\omega_{\rm p}^{\prime\,2}}{\gamma^{\prime}} &= \gamma^{-1/4} \Gamma^{-2} \frac{\omega_{\rm pL}^2}{\gamma} \\ \frac{\omega_{\rm c}^{\prime\,2}}{\gamma^{\prime\,2}} &= \gamma^{1/2} \Gamma^{-2} \frac{\omega_{\rm cL}^2}{\gamma^2}, \end{split}$$

where the quantities ω_{pL} and ω_{cL} are the proton plasma and cyclotron frequencies at the light cylinder surface, $\gamma \omega_{pL}^2 = \omega_{cL}^2$. We see that because of the jet plasma acceleration the plasma frequency becomes much less than the cyclotron frequency. In this case, the dependence of the eigen plasma wave frequency on the angle θ' is

$$\omega' = \frac{\omega_{\rm p}'}{{\gamma'}^{1/2}}\cos\theta'.$$

When the frequency equals ω' and the wave vector equals $k' \cos \theta' = \omega'/c$, the beam instability develops. Substituting the relation $\omega' = \omega'_p \cos \theta' / {\gamma'}^{1/2} + \delta \omega' = k'c \cos \theta' + \delta \omega'$ into equation (11) we obtain

$$\delta\omega^{\prime 3} = \frac{\omega_{\rm p}^{\prime 3}}{2\gamma^{\prime 3/2}} \frac{n_{\rm b}^{\prime}\gamma^{\prime}}{n_{\rm p}^{\prime}} \left(\frac{\cos^{3}\theta^{\prime}}{\left\langle \gamma_{\rm b}^{\prime -3} \right\rangle^{-1}} + \frac{\cos\theta^{\prime}\sin^{2}\theta^{\prime}}{\left\langle \gamma_{\rm b}^{\prime -1} \right\rangle^{-1}} \right). \tag{12}$$

Because $\langle \gamma_b^{\prime -3} \rangle \ll \langle \gamma_b^{\prime -1} \rangle$, the maximum of the beam instability increment is achieved at the angle $\cos \theta' = 3^{-1/2}$,

$$Im(\delta\omega')|_{\max} = \frac{\omega_{\rm p}'}{2\gamma'^{1/2}} \left(\frac{n_{\rm b}'\gamma'}{n_{\rm p}'\langle\gamma_{\rm b}'^{-1}\rangle^{-1}}\right)^{1/2}$$

Thus, we conclude that the beam instability in the jet near its base produces the fast growing oscillations of the electric field with the frequency $\omega' = 3^{-1/2}\omega_{cL}\gamma^{-9/8}\Gamma^{-1}$ and with the wave vector directed at the angle θ' to the jet axis, $\cos \theta' = 3^{-1/2}$, in the comoving frame. These oscillations will be observed in the laboratory frame as the time variations of the jet radiations with the characteristic time of $\tau \simeq 2\pi\Gamma/\omega'$ and under the angle θ , $\cos \theta = [(\Gamma^2 - 1)^{1/2} + \Gamma \cos \theta']/[\Gamma + (\Gamma^2 - 1)^{1/2} \cos \theta'], \theta \simeq ((3^{1/2} - 1)/(3^{1/2} + 1))^{1/2}\Gamma^{-1}$. Using typical values of B_L , γ and Γ , we find

$$\tau = 450 \left(\frac{\gamma}{10^5}\right)^{9/8} \left(\frac{\Gamma}{10}\right)^2 \left(\frac{B_{\rm L}}{10^2 G}\right)^{-1} \text{ s},$$

$$2\theta = 6 \left(\frac{\Gamma}{10}\right)^{-1} \text{ grad.}$$
(13)

These values are in good agreement with observations. Indeed, for example, Aharonian et al. (2007) measured the flare of the very high energy (VHE) gamma-ray flux from blazar BL Lac PKS 2155–304, which lasted \approx 80 min. The light curve was modulated by almost harmonic oscillations with the period of 10 min. This fact suggests the existence of some coherent process. The proposed mechanism of the beam instability explains two facts together: the flare is the injection of fast particles and oscillations are the excitation of the coherent wave. If we adopt the barionic scenario of VHE gamma-ray generation, then their intensity is proportional to the energy density of fast protons $nm_{\rm p}c^2\gamma$. If we add the proton beam, then the photon flux F becomes, $F = F_0(1 + n_b \gamma_b / n \gamma)$. The strong plasma wave is excited at the ion plasma frequency ω' . It modulates the beam energy $\gamma_{\rm b} = \gamma_{\rm b0} + \gamma_{\rm b1} \psi(\omega' t')$, where ψ is a function close to the periodic function. The beam instability is the resonance instability; all particles of the relativistic beam, moving with velocities very close to the speed of the light, are caught into phase resonance with the exciting wave, $\omega' = \mathbf{k}' c$. The wave amplitude grows exponentially and beam particles are trapped into wave potential holes. The trapping stabilizes the wave amplitude, and the wave energy becomes of the order of the beam energy:

$$\frac{E'^2}{4\pi} \simeq n'_{\rm b} \gamma'_{\rm b} m_{\rm p} c^2 \left(\frac{n'_{\rm b} \gamma'}{n'_{\rm p} \left\langle \gamma'_{\rm b}^{-1} \right\rangle^{-1}} \right)^{1/3}.$$

Such a phenomenon was observed experimentally for the electron beam instability in the laboratory plasma (Gentle & Lohr 1973). Thus, the mechanism suggested here does not demand complex modelling calculations.

5 SUMMARY

If the Blandford–Znajek mechanism of extracting energy from a rotating black hole works, then a massive black hole and an accretion disc produce the relativistic jet. The energy and the angular momentum of the jet are fed from the black hole rotation. The accretion is not the source of the energy in this case. However, it is important for the delivery of the strong magnetic field to the black hole vicinity. Besides, the electric current flows through the disc. It is part of the current loop excited by the voltage created by the rotating dynamo. The strong magnetic field in the Blandford–Znajek mechanism results in large synchrotron losses for electrons, and almost all energy is transformed to protons. Protons obtain

relativistic energy crossing the light cylinder surface. Thus, a jet consists of ultrarelativistic protons ($\gamma \gg 1$) and 'cold' electrons. The main proton energy is in azimuthal rotation, and only a small part is in the radial motion. The initial radial expansion of a jet is small, $v_r = c\gamma^{-1/2}$. Behind the light cylinder, jet particles are accelerated along the magnetic field, moving in the expanding poloidal magnetic field $B_{\rm p}$, and the jet obtains the longitudinal Lorentz factor $\Gamma \simeq (B_{\rm p}/B_{\rm L})^{-1/2} = R/R_{\rm L}$. The jet cools, remaining hot; the internal proton energy is $\langle \gamma' \rangle = \gamma^{13/16} (B_{\rm p}/B_{\rm L})^{1/2} \gg 1$. However, the density of the proton energy $n'\langle\epsilon'\rangle$ is smaller than the energy density of the magnetic field in the jet $B^2/8\pi$. This circumstance provides the jet stability (Istomin & Pariev 1994, 1996). Electrons in the jet remain cold except in places where they can be accelerated by different mechanisms (Bisnovatyi-Kogan & Lovelace 1995; Bereznyak, Istomin & Pariev 2003). The classical time of the energy transfer from hot protons to electrons is very large because of the high proton energy, $t = \Gamma \langle \epsilon' \rangle^2 / (\pi c^4 c n' \Lambda) \gg l/c$. Λ is the Coulomb logarithm and *l* is the jet length. The short time variability of the jet, $\tau \simeq 300$ s, can be explained by the excitation of ion plasma waves, and can be observed at small angles to the jet axis (equation 13). The mechanism of the origin of the time variability for the explanation of the change of VHE photon flux observed in blazars, proposed in this paper, has a 'universal' character. Of course, it is possible to obtain even shorter time-scales of variability, suggesting the fast motion of a nonuniform jet (see, for example, Ghisellini et al. 2009). In our mechanism, the time of variability τ is the natural characteristic of the jet plasma on the jet base and it depends only on proton density and energy.

In conclusion, it is worth repeating that here we have discussed the origin of a jet by the Blandford–Znajek mechanism only. It demands a strong enough magnetic field nearby the black hole. The estimations show that, in this case, the magnetic field intensity is inversely proportional to the black hole mass, and has moderate values for the AGN, $B \simeq 10^3 - 10^4$ G. As for black holes of lower masses, much higher magnetic fields near the black hole horizon cannot be achieved. However, this will be the subject of further investigations.

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