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Relativistic jets and nuclear regions in AGN

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Abstract. The main driving mechanism of relativistic jets is likely related to magnetic fields. These fields are able to tap the rotational energy of the central compact object or disk, accelerate and collimate matter ejecta. To zeroth order these outflows can be described by the theory of steady, axisymmetric, ideal magnetohydrodynamics. Results from recent numerical simulations of magnetized jets, as well as analytical studies, show that the efficiency of the bulk acceleration could be more than $\sim 50\%$. They also shed light to the degree of the collimation and how it is related to the pressure distribution of the environment, the apparent kinematics of jet components, and the observed polarization properties.

Key words. MHD – relativity

1. Introduction

The observed superluminal motion of the components of many AGN jets is a clear indication of their relativistic motion. Unfortunately the apparent speed alone is not enough to give the true velocity of these outflows. It can only give a lower limit of the Lorentz factor $\gamma = \left(1-\beta^2\right)^{-1/2}$, through the relation $\beta_{\rm app} = \beta \sin\theta_n \left(1-\beta\cos\theta_n\right)^{-1}$ which also involves the angle θ_n between the flow direction and the line of sight. In cases where the Doppler factor $\delta \equiv \gamma^{-1}\left(1-\beta\cos\theta_n\right)^{-1}$ can be also found, the two relations give the flow speed as a function of distance from the core. For example Unwin et al. (1997), by combining radio and X-ray flux measurements and interpreting the latter as synchrotron self-Compton, deduced a change in the bulk Lorentz factor of the C7 component in 3C345, from $\gamma \sim 5$ to $\gamma \sim 10$ over a deprojected distance range of $\sim 3-20$pc (they also infer a decrease in the Doppler factor from 12 to 4 and an increase in the angle between the velocity of the component and the line of sight from 2 to 10° during the same period). Similarly Piner et al. (2003) inferred an acceleration from $\gamma = 8$ at $r < 5.8$pc to $\gamma = 13$ at $r \approx 17.4$pc in 3C 279. Another way to estimate the Doppler factor is by comparing the variability timescale with the light-travel time across the emitting region (Jorstad et al., 2005).

The extended (parsec scale) bulk acceleration seems to be a general characteristic of AGN jets. Besides the two examples mentioned above, where an increase of the Lorentz factor was directly deduced, there are cases where there are observations of several superluminal features with the innermost one typically exhibiting the smallest proper motion (e.g., Homan et al., 2001). Sikora et al. (2005) give a more general argument related to the extended acceleration: If the bulk flow near the disk is sufficiently fast ($\gamma \gtrsim 5$) it would Comptonize photons coming from the disk, producing bulk-Compton features. The absence of such features indicate that acceleration up to $\gamma \gtrsim 10$ must take at least 10^3 Schwarzschild radii.[1]

Extended acceleration is unlikely to be driven hydrodynamically in a proton-electron outflow. If T_i is the initial temperature, energy conservation implies that pure-hydrodynamic driving gives Lorentz factors $\gamma \sim k_{\rm B} T_i / m_p c^2$, which is of order unity even if the temperature is as high as 10^{12}K. In addition, the distance on which the Lorentz factor attains its terminal value is of the order of the sonic distance, certainly much smaller than pc-scales where the acceleration is inferred from the observations, since the sonic surface is located close to the event horizon of the central black-hole. Alternatively, a possible heating source makes $\gamma \gg 1$ possible (e.g., Meliani et al. (2004)). It is, however, unclear how such a heating source would be established in a natural way on pc-scales. The most likely alternative is magnetic driving, which is considered in the next sections of this article.

AGN jets are also well collimated; e.g., in the galaxy M87, the jet is seen opening widely in its formation region, at an angle of about 60° nearest the black hole, but is squeezed down to only 6° at 100 Schwarzschild radii (Biretta et al., 2002; see also Krichbaum et al., 2006). The jet opening angle continues to decrease at larger distances. Magnetic self-collimation has long been thought to be the underlying mechanism for the observed jet shape.

Polarization maps are very useful and strongly connected to the magnetic field of the jets. They are in general consistent with a helical magnetic field structure (Marscher et al., 2008). In addition, the observed electric field polarization vectors and the Faraday rotation measure gradients across the jet, support the existence of helical magnetic fields with strong transverse (azimuthal) component (e.g., Gabuzda et al. (2004)).

All the above observed characteristics are related or can be explained with the help of magnetic fields, making the magnetic driving of jets the most plausible explanation. The magnetic field can extract energy (Poynting flux) from the source, as well as angular momentum (helping to solve the accretion problem in the underlying disk). The Lorentz force can transfer this energy to matter kinetic energy, resulting in a relativistic flow if the Poynting-to-mass flux ratio is sufficiently high near the source. It also transfer angular momentum to the matter, giving to the jet a nonzero angular velocity.[2]

The components of AGN jets follow curved (helical) trajectories over the years. There have been various tries to explain the apparent jet kinematics: as Kelvin-Helmholtz instabilities (Hardee & Walker, 2005), or due to source precession (Lobanov & Roland, 2005). However, since a magnetized outflow is rotating it may explain (or at least play some partial role to) the observed kinematics.

The magnetohydrodynamic description of AGN jets, which is related with all the above observed characteristics, will be discussed in the following.

2. The ideal-MHD description

The system of equations of special relativistic, steady, ideal MHD, consists of the Maxwell equations $\nabla \times \boldsymbol{B} = 4\pi \boldsymbol{J}/c$, $\nabla \cdot \boldsymbol{E} = 4\pi J^0/c$, $\nabla \cdot \boldsymbol{B} = 0$, $\nabla \times \boldsymbol{E} = 0$, Ohm's law $\boldsymbol{E} = \boldsymbol{B} \times \boldsymbol{V}/c$, the continuity $\nabla \cdot (\rho_0 \gamma \boldsymbol{V}) = 0$, entropy $\boldsymbol{V} \cdot \nabla \left(P/\rho_0^\Gamma\right) = 0$, and momentum $\gamma \rho_0 (\boldsymbol{V} \cdot \nabla)(\xi \gamma \boldsymbol{V}) = -\nabla P + \left(J^0 \boldsymbol{E} + \boldsymbol{J} \times \boldsymbol{B}\right)/c$ equations. Here $\boldsymbol{V}$ is the velocity of the outflow, γ the associated Lorentz factor, $\boldsymbol{E}, \boldsymbol{B}$ the electric and magnetic fields as measured in the central object's frame, $J^0/c, \boldsymbol{J}$ the charge and current densities, ρ_0, P the gas rest-mass density and pressure in the comoving frame, while ξc^2 is the enthalpy-to-rest-mass ratio; for a polytropic equation of state (with index Γ) $\xi c^2 = c^2 + [\Gamma/(\Gamma - 1)] (P/\rho_0)$.

[1] They also argue that the electron/positron kinetic energy (estimated from the emissivity of blazar events), is too small to support the energetics of blazars and of radio lobes in quasars. Thus, the dynamics of AGN jets is likely dominated by protons rather than leptons.

[2] The angular velocity is small though, because the lever arm of the jet at large distances is much larger compared to its values near the source.

Assuming axisymmetry [$\partial/\partial\phi = 0$, in spherical (r, θ, ϕ) and cylindrical (z, ϖ, ϕ) coordinates with $\hat{z}$ along the rotation axis and the central object at $(\varpi = 0, z = 0)$], partial integration of the equations is possible resulting in the five following integrals of motion (e.g. Vlahakis & Königl, 2003a)[3]:

(i) The mass-to-magnetic flux ratio

$$\Psi_A = \frac{4\pi\gamma\rho_0 V_p}{B_p} . \qquad (1)$$

(ii) The field angular velocity

$$\Omega = \frac{V_\phi}{\varpi} - \frac{V_p}{\varpi}\frac{B_\phi}{B_p} . \qquad (2)$$

(iii) The specific angular momentum

$$L = \xi\gamma\varpi V_\phi - \frac{\varpi B_\phi}{\Psi_A} . \qquad (3)$$

(iv) The adiabat

$$Q = \frac{P}{\rho_0^\Gamma} . \qquad (4)$$

(v) The energy-to-mass flux ratio

$$\mu c^2 = \xi\gamma c^2 - \frac{\varpi\Omega B_\phi}{\Psi_A} . \qquad (5)$$

Note that for $B_p > 0$ equation (1) gives $\Psi_A > 0$ while from equation (2) $B_\phi < 0$ (such that V_ϕ remains subluminal beyond the light surface). Thus, the last terms of the integrals L and μ, which are related to the electromagnetic field, are positive. From equation (5) it is evident that the value μ represents the maximum possible Lorentz factor that a flow can attain after the acceleration phase, corresponding to a cold ($\xi = 1$) flow with zero Poynting flux ($B_\phi = 0$). However, this cannot always happen; an important unknown of the MHD problem is the asymptotic Lorentz factor γ_∞ and the efficiency of the acceleration γ_∞/μ.

The poloidal magnetic flux function

$$A = \frac{1}{2\pi}\iint \boldsymbol{B}_p \cdot d\boldsymbol{S} \qquad (6)$$

is also conserved along the flow and can be used as the label of each field-streamline.

An important dimensionless combination of the integrals is the "Michel's magnetization parameter"

$$\sigma_{\rm M} = \frac{A\Omega^2}{\Psi_A c^3} . \qquad (7)$$

Using equations (1-5) the physical quantities can be expressed as functions of γ, A, and the integrals (which are functions of A):

$$\rho_0 = \frac{\Psi_A^2 \xi\left(\mu - \xi\gamma\right)}{4\pi\left[\xi\gamma\left(x^2-1\right) + \mu\left(1 - x_{\rm A}^2\right)\right]} , \qquad (8)$$

$$\xi = 1 + \frac{\Gamma}{\Gamma - 1}\frac{Q}{c^2}\rho_0^{\Gamma-1} , \quad P = Q\rho_0^\Gamma , \qquad (9)$$

$$\boldsymbol{B} = \frac{\nabla A \times \hat{\phi}}{\varpi} - c\Psi_A\frac{\mu - \xi\gamma}{x}\hat{\phi} , \qquad (10)$$

$$\boldsymbol{E} = -\frac{\Omega}{c}\nabla A , \qquad (11)$$

$$\boldsymbol{V}_p = \frac{\left[\xi\gamma\left(x^2-1\right) + \mu\left(1 - x_{\rm A}^2\right)\right]}{\Psi_A\xi\gamma\left(\mu - \xi\gamma\right)}\boldsymbol{B}_p , \qquad (12)$$

$$\boldsymbol{V}_\phi = c\frac{\xi\gamma - \mu\left(1 - x_{\rm A}^2\right)}{\xi\gamma x}\hat{\phi} . \qquad (13)$$

Here $x = \varpi\Omega/c$ is the cylindrical distance in units of the light surface lever arm on each field-streamline A =constant, and $x_{\rm A} = \left(L\Omega/\mu c^2\right)^{1/2}$ is its value at the Alfvén point. In fact for initially Poynting-dominated flows the Alfvén and light surfaces almost coincide, meaning that $x_{\rm A} \approx 1$.

The two remaining unknowns γ and A obey the two components of the momentum equation on the poloidal plane.

The momentum equation along the flow gives the so-called Bernoulli, or, wind-equation

$$\gamma^2 - 1 = \frac{\sigma_{\rm M}^2}{S^2}\left[\frac{\gamma}{\mu - \xi\gamma} - \frac{\xi\gamma - \mu\left(1 - x_{\rm A}^2\right)}{\xi\left(\mu - \xi\gamma\right)x^2}\right]^2 + \left[\frac{\xi\gamma - \mu\left(1 - x_{\rm A}^2\right)}{\xi x}\right]^2 , \qquad (14)$$

[3] The subscripts p/ϕ denote poloidal/azimuthal components.

where the function $S \equiv \frac{A}{\varpi|\nabla A|} \quad \left(= \frac{A}{\varpi^2 B_p}\right)$.
The Bernoulli is an algebraic equation and gives the $\gamma(x)$ for each field-streamline A =constant, provided that we know the function S. Since S is related to the magnetic flux distribution (which is the solution of the force-balance in the transfield direction), it provides a link between the two components of the momentum equation and emphasizes that the two equations are coupled and must be solved simultaneously. However, it is possible to derive some conclusions on the flow bulk Lorentz factor without solving the transfield force-balance. First, we can simplify the Bernoulli equation as follows. We focus on the super-Alfvénic part of the flow where $x \gg 1$, the flow is cold $\xi \approx 1$, and already relativistic $\gamma \gg 1$. By neglecting terms of order x^{-2} and γ^{-3} equation (14) becomes

$$\frac{\sigma_{\rm M}}{S} \approx \mu - \gamma - \frac{\mu}{2\gamma^2} . \qquad (15)$$

Its derivative along the flow yields

$$\frac{\sigma_{\rm M}}{S^2} \boldsymbol{V} \cdot \nabla S \approx \left(\frac{1}{\mu} - \frac{1}{\gamma^3}\right) \boldsymbol{V} \cdot \nabla \gamma . \qquad (16)$$

The latter gives the conditions on the fast-magnetosonic surface (where the derivative of γ is 0/0): $(\boldsymbol{V} \cdot \nabla S)_f \approx 0$,

$$\gamma_f \approx \mu^{1/3} , \text{ and } \quad S_f \approx \frac{\dfrac{\sigma_{\rm M}}{\mu}}{1 - \dfrac{3}{2\mu^{2/3}}} \approx \frac{\sigma_{\rm M}}{\mu} . \qquad (17)$$

(Note that at the fast-magnetosonic surface the value $\gamma_f \approx \mu^{1/3}$ is much smaller than μ, meaning that the flow is still Poynting flux dominated. This is the reason why the super-fast part of the flow is the most interesting to analyze, since only in this regime the transfer of a significant part of the Poynting flux to kinetic energy flux may take place.) The function S is the effective surface of a de Laval nozzle, and becomes minimum at the critical point. After that point we may further simplify the Bernoulli equation to $\gamma \approx \mu - \sigma_{\rm M}/S$, which clearly shows that the increase of the function S after the fast-magnetosonic surface is directly related to the flow acceleration. A geometrical meaning of the function S can be understood as follows: The area between two neighboring field lines with fluxes A and $A + \delta A$ is $2\pi\varpi\delta\ell_\perp$, where $\delta\ell_\perp$ is the distance between the field lines on the poloidal plane. Since $\delta A = 2\pi\varpi\delta\ell_\perp B_p$ we can write $S = (2\pi A/\delta A)(\delta\ell_\perp/\varpi)$. Thus, increasing S corresponds to expanding field lines in a way such that the $\delta\ell_\perp$ increases faster than ϖ.
The asymptotic Lorentz factor is

$$\gamma_\infty \approx \mu - \frac{\sigma_{\rm M}}{S_\infty} . \qquad (18)$$

From the definition of the function S we can estimate $S_\infty \approx 1$, corresponding to a uniform distribution of the magnetic flux (resulting in $\varpi|\nabla A| \approx A$, or, $B_p\varpi^2 \approx A$). In other words, as the poloidal field lines expand the function S increases and this causes acceleration of the flow. However, since the available solid angle is finite, S_∞ cannot become infinity, but only reach a value ~ 1. For this reason the acceleration efficiency is less than 100%. It can, however, reach values close to 100% if $\sigma_{\rm M} \ll \mu$. In this case $S_f \ll 1$ and the field lines are bunched inside a small solid angle at the fast-magnetosonic surface. This is an extreme case. Nevertheless, efficiencies of the order of 50% can be easily reached. They correspond to $S_f \approx 1/2$, a typical value for a dipolar field. Indeed, existing analytical solutions in the literature Li et al. (1992); Vlahakis & Königl (2003a,b, 2004); Beskin & Nokhrina (2006) as well as simulations Komissarov et al. (2007) show that the efficiency of the magnetic acceleration is 50% or more.[4]

Note that the expansion needed for the decline of the function S, which in turn gives rise to efficient acceleration, goes along with the self-collimation property of magnetized outflows: Field lines that are closer to the rotation axis collimate faster than the outer ones.

However, there must always be an external medium that confines the system at its

[4] On the other hand, if $S_f \approx 1$ (as is the case for a monopolar field near the equatorial plane) the efficiency is very small. This shows the peculiarity of the monopole-field case and explains why the results of Michel (1969) are different compared to the studies mentioned above.

boundaries (e.g., an external wind as suggested by Gracia et al., 2005, or some other external medium). As was found in the simulations of Komissarov et al. (2007), there is a one-to-one correspondence between the external pressure and the shape of the jet. External pressure $P_{\rm ext} \propto r^{-\alpha}$ with $\alpha = 3.5, 2, 1.6$, and 1.1 corresponds to flow shapes $z \propto \varpi^{a}$ with $a = 1, 1.5, 2$, and 3, respectively. As expected, the more collimated jets correspond to the less steep external pressure distributions.

For a typical jet that is launched by a rapidly rotating black hole with $\mu = 16$, Komissarov et al. (2007) found that the equipartition between Poynting and kinetic energy flux is reached at a distance

$$r_{eq} \approx 2 \times 10^{16} \left(\frac{M}{10^8 M_\odot} \right) \left(\frac{\Theta_j}{0.1} \right)^{-1} \ \rm cm\,,$$

where Θ_j is the jet opening half-angle. In the case where the jet originates in a Keplerian accretion disk, the distance where equipartition is reached is

$$r_{eq} \approx 10^{17} \left(\frac{M}{10^8 M_\odot} \right) \left(\frac{\varpi_0}{10\, r_g} \right)^{3/2} \left(\frac{\Theta_j}{0.1} \right)^{-1} \ \rm cm\,,$$

where $r_g \equiv GM/c^2$ and ϖ_0 is the ejection cylindrical radius. Beyond this distance the jet continues to be accelerated, entering the matter-dominated regime. If blazar flux variability is associated with the propagation of strong shocks within the jet then we can expect this behaviour to originate on scales $\gtrsim r_{eq}$. When the simulated jets reach $r \simeq 10 r_{eq}$, their characteristic Lorentz factor becomes ~ 10. These properties of the extended magnetic acceleration region are in very good agreement with the observational inferences summarized in Section 1.

Similar results are found using the self-similar exact solutions of the MHD equations (e.g., Vlahakis & Königl, 2004). As stated in Section 1, the rotation of a magnetized flow could explain the observed helical apparent trajectories of jet components (see Vlahakis, 2006), although one can not rule out other explanations previously mentioned.

Polarization measurements seem also to be in a general agreement with the MHD modeling and in particular the large scale helical magnetic field in the acceleration and collimation zone (Marscher et al., 2008).

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