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On a MHD classification of AGN

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Abstract. Some general features of ideal MHD plasma outflows from rotating and magnetized central gravitating objects such as AGN are briefly discussed. The asymptotic structure, morphology and degree of collimation of the outflow are analysed in the light of a family of exact solutions produced via a nonlinear separation of the variables in the full set of the MHD equations. Among the interesting features of this solution is a quantitative energetic criterion for the transition of the morphologies of the outflows from highly collimated jets to uncollimated winds. It is proposed that in a space where the two main variables are the energy of the magnetic rotator and the angle between the line of sight and the ejection axis, some observed characteristics of AGN can be understood. Thus, while the horizontal AGN classification from Type 0 to Types 1 and 2 is mainly an orientation effect, as in the standard model, the vertical AGN classification as those with uncollimated outflows (radio-quiet sources: Seyferts, QSO's) and collimated outflows (radio-loud sources : radio galaxies and Blazars) depends on the efficiency of the magnetic rotator and the environment in which the outflows propagate.

INTRODUCTION

The extremely high luminosity of the emitted radiation in $\gamma-$ray, X-ray and UV continuum from rather small volumes in active galactic nuclei (AGN) is believed to be gravitational energy released by matter spiraling around a supermassive central object, possibly a black hole of about $10^9 M_\odot$. Although such an enigmatic object is hidden from our own direct scrutiny, information about the central engine which produces the enormous observed activity may be obtained from other basic ingredients which have gradually formed what we call the standard picture of an AGN: (i) the accretion disk which emits in the UV and soft X-rays, mostly from about $(2-100)$ gravitational radii R_g, (ii) the broad line optically emitting clouds extending up to a few $10^3 R_g$ from the center, (iii) an obscuring dusty torus with an inner radius of a few $10^3 R_g$, (iv) the narrow line regions extending from about $(10^4 - 10^6) R_g$, and finally (v) powerful jets of plasma detected from the sub-parsec to the Mpc scales, mainly in the radio but also in the optical, UV and X-rays. Note in particular, that ultra high energy γ-rays have also been observed from the central regions of several Blazars, as it is discussed in detail in this volume. Hence, together with the emission of radiation from the immediate neighborhood of an AGN, jets provide a crucial link between the easier observed large Mpc scale and the finer sub-parsec scales where presumably the plasma of the jets is accelerated in a few gravitational radii from the center of the AGN.

AGN are commonly divided into three broad classes, according to the phenomenology of their emission of radiation in the radio and optical/UV parts of the spectrum [1], [39]. Type 1 AGN have bright continua with broad emission lines. In the radio-quiet group belong the low luminosity Seyfert 1 galaxies and the higher luminosity radio-quiet quasars (QSO), while in the radio-loud group belong the broad line radio galaxies at low luminosities and the flat or steep spectrum radio-loud quasars at higher luminosities. On the other hand, the Type 2 AGN have weak continua with narrow emission lines. In the radio quiet group belong the low luminosity Seyfert 2 galaxies and narrow emission line X-ray galaxies while in the radio-loud group belong the narrow-line radio galaxies with the two distinct morphologies of Fanaroff-Riley I (low-luminosities) and Fanaroff-Riley II (higher luminosities). Finally, the Type 0 AGN contains the rest AGN with weak or unusual emission lines, i.e., in the radio-quiet end the broad absorption lines QSO and in the radio-loud end the Blazars (BL Lacs and flat spectrum radio galaxies). Hence, such a unification model organizes the various classes of AGN according to orientation, beaming and obscuration effects [39].

However, as noted by Urry and Padovani [39], "although the classification of AGN in Types 0 - 2 appears to be primarily an orientation effect, the differences within the same type, e.g., as we move from the radio-quiet to the radio-loud galaxies, appears to be due to an as-yet unknown physical effect". It seems that in radio-loud AGN the jet is relativistic while in radio-quiet AGN it is not. Various possibilities have been suggested, such as that radio-loudness is related to host galaxy type [33], or, that it is the black hole's spin that distinguishes between the lower spin radio-quiet galaxies from the higher black hole spin in radio-loud galaxies [5], [38], or, that there are differences in the rate of nuclear feeding [27], [2].Nevertheless, in this paper we shall outline a quantitative physical criterion for such a transition from radio-quiet to radio-loud galaxies based on new results from an analytical modeling of MHD outflows from a magnetized and rotating central gravitating object.

Among the basic questions for understanding the physics and role of jets in AGN are those related to their initial acceleration in the near environment of the AGN, as well as their morphology as they propagate away from the central region. Regarding their initial acceleration, two classes of plasma composition and acceleration mechanisms for outflows associated with AGN are still being pursued. In one, the nonrelativistic or relativistic plasma consists basically of electrons/protons and is magnetocentrifugally accelerated from an accretion disk via the classical Blandford & Payne [4] acceleration mechanism. In the second, the rotational energy of the accreting black hole is extracted by a magnetic field, then converted to Poynting flux and finally to relativistic electron/positron pairs via the Blandford & Znajek mechanism [3]. It is likely that jets associated with quasars are a mixture of electron/proton and electron/positron pairs, since pure electron/positron jets overpredict soft X-ray radiation from quasars, while pure proton-electron jets predict too weak nonthermal X-ray radiation [32]. In fact it has been suggested that jets consist of electron/positron pairs close to their axis, surrounded by an electron/proton plasma [31].

Concerning their morphology, plasma outflows may appear either as collimated beams or as uncollimated winds. However, apart the case of the solar wind, uncollimated flows are hardly observable, while jets are observed in several astrophysical environments, from star formation regions to distant AGN. This is mainly due to the much higher density inside jets as opposed to the loosely collimated winds. Furthermore in AGNs where beams move at relativistic speeds, the emission may be largely amplified by Doppler boosting if the jet is pointed towards the observer. In the following we outline some basic properties of MHD outflows towards understanding their morphologies as they are associated with AGN. In particular, we shall use an exact outflow model based on semi-analytical solutions of the full set of the MHD equations in order to investigate the physical basis of the transition of AGN within the same type (aspect) in a unifying classification scheme.

ACCELERATION, ANGULAR MOMENTUM EXTRACTION AND CONFINEMENT

Acceleration. From a rather general perspective, an observational characteristic of many cosmic plasma outflows is that they seem to be *accelerated* to relatively high speeds which may even reach values close to the speed of light in the most powerful AGN jets. The most often invoked mechanisms to accelerate these flows are of thermal, or, of magnetocentrifugal origin. Thus, if the magnetic field does not play a crucial role in the acceleration, the plasma is thermally driven, as is the case in the low-speed solar wind with the heating provided by the dissipation of waves, Joule heating, etc. In this case the presence of a hot corona around the disk that is responsible for the acceleration is essential and the terminal speed obtained in this way is proportional to the square root of the coronal temperature. Conversely, if the magnetic field and rotation are dominant, then the primary source of acceleration may be either due to the build up of a toroidal magnetic pressure (the so-called 'uncoiling spring' model), or it can be due to magnetocentrifugal forces (the so-called 'bead on a rotating wire' model). In this last case the acceleration comes from the conversion of Poynting energy flux to kinetic energy flux. In both those last cases the gain in kinetic energy is proportional to the energy that brings the magnetic field lines into rotation, i.e., the energy of the magnetic rotator; this energy is the product of the rotational frequency and the total specific angular momentum, ΩL. Such a magnetocentrifugal driving mechanism seems to be efficient in disk winds wherein a hot corona is not an absolute requirement.

Angular momentum extraction. Magnetocentrifugally driven outflows extract also rather efficiently angular momentum from their source allowing thus mass to accrete to the central object. This magnetic removal of the angular momentum of the accreted material is efficient because the magnetic lever arm ϖ_a is larger by a factor of order 10 than the cylindrical radius of the footpoint of a fieldline, ϖ_o [37]. It is interesting that even if a tiny percentage of order of 1% of the accreted mass rate $\dot{\mathcal{M}}_a$ is lost through a jet, $\dot{\mathcal{M}}_j$, the major part of the angular momentum of the infalling gas is removed and so this gas can be freely accreted by the central object. This can be quantitatively seen as follows. Since the outflow practically corotates approximately up to the Alfvén distance ϖ_a, the specific angular momentum carried by this outflow is $\dot{J}_j = \dot{\mathcal{M}}_j \Omega \varpi_a^2$, while the angular momentum that needs to be extracted locally at the footpoint ϖ_o of the same fieldline on the disk level in order that accretion takes place at a rate $\dot{\mathcal{M}}_a$ is $\dot{J}_a = (1/2)\dot{\mathcal{M}}_a \Omega \varpi_o^2$. If $\dot{J}_j = f\dot{J}_a$, while the remaining fraction $(1-f)$ of the angular momentum of the accreted mass may be carried away by viscous stresses, then, $\dot{\mathcal{M}}_j/\dot{\mathcal{M}}_a = (f/2)(\varpi_o^2/\varpi_a^2)$. For $\varpi_a \sim 10\varpi_o$, we see that $\dot{\mathcal{M}}_j \leq 1\%\dot{\mathcal{M}}_a$. It looks then that plasma jets play a doubly crucial role in the AGN picture, (i) by allowing material to accrete close to the gravitational horizon of the supermassive central object wherein, (ii) by converting its gravitational energy into high energy radiation, they power the high luminosity that we observe coming from the central parts of the AGN.

Confinement. Another observational constraint is the *morphology* of the outflow. As it is well known, astrophysical outflows are either confined and collimated into a rather cylindrical jet or they remain expanding into a more or less radial wind. Two basic mechanisms may be responsible for collimation, one is of thermal origin and the other is of magnetic origin. Thermal confinement means that the surrounding medium has a higher pressure than the flow such that there is a pressure gradient forcing the outflow to collimate along its ejection axis. In other words, only outflows underpressured with respect to their surrounding environment may be thermally confined. In fact, such a situation seems to occur in many extragalactic jets, as deduced from X-ray data on the hot plasma surrounding early-type galaxies and clusters of galaxies, [13]. But there also exists a second confining mechanism which works both for under- and over-pressured jets. This is the magnetic confinement of the outflow by a toroidal magnetic field wound around the jet. In fact, in recent observations of highly optically-polarized, compact radio-loud quasars (HPQ) the electric vectors of the polarized 43 GHz radio cores are roughly aligned with the inner jet direction indicating magnetic fields perpendicular to the flow [23]. This implies that such beams carry some electric current that eventually closes at their surface or outside. Note that the building of the toroidal magnetic field is done at the expense of the Poynting flux. Thus, in a pure magnetic jet all the Poynting flux cannot be converted to kinetic energy, if part of it remains to confine the jet. Obviously reality in most cases and particularly in the case of AGN, may involve a combination of thermal and magnetic processes in the acceleration and confinement of the outflow.

NUMERICAL SIMULATIONS OF UNSTEADY MHD OUTFLOWS

Electron/proton or electron/positron plasma outflows are usually modeled to zeroth order via the set of the ideal MHD equations. The time-dependent MHD problem can be treated only by means of numerical simulations in which, for instance, the boundaries do not introduce spurious effects and the system relaxes into a reproducible final state. Another difficulty with the numerical simulations has to do with the fact that AGN jets often extend over lengths more than six orders of magnitude their width, while the available grid sizes are much smaller. One way out of this constraint is to solve the problem via a combination of numerical and analytical techniques. For instance, one may solve the time-dependent problem in the near-zone containing all critical surfaces and obtain a suitable steady state with a superfast outflow. This solution can be then used as the boundary condition for solving in the superfast domain the hyperbolic steady state problem analytically, to unlimited large distances from the central source [6], [35].

Sakurai's [28] simulation obtained for the first time a steady-state solution for the rotating and magnetized solar wind, by using an iterative method to solve the coupled set of the Bernoulli and transfield equations. This study used a split monopole configuration for the initial magnetic field and obtained an asymptotic state where the outflow is logarithmically (i.e., weakly) collimated around the rotational axis. The majority of the other numerical simulations on the temporal evolution of outflows has been devoted to model magnetically driven disk-winds [24], [16], [40], [18], [19]. In several of these cases, after the initial transient phase when a torsional Alfvén wave develops and drives the initial acceleration of the flow like an uncoiling spring, the solutions converge to a weakly collimated structure, where the confinement is done by the toroidal magnetic field. However, for the numerical constraints we mentioned above, the outflow cannot be simulated in regions far from the base to follow its degree of collimation and where the thermal pressure could also contribute to its confinement.

A different approach has been followed by Bogovalov & Tsinganos [6], [35], who investigated the temporal evolution of a solution of an initially radial magnetosphere with a uniform flow along the magnetic lines, a case suited to model a wind emerging from a spherically symmetric central body. After sufficient time has elapsed from the onset of the rotation such that a toroidal magnetic field is generated, a cylindrically collimated steady state is realized. In that asymptotic state, a relation between the strength of the magnetic rotator and the amount of collimation is also found, in agreement with the properties of the steady analytical solutions we shall discuss in the next section. In particular the degree of collimation is determined by a parameter which expresses the ratio of the corotation speed at the Alfvén transition to the Alfvén speed there. The main results of the study have been also found in the simulations of Keppens & Goedbloed [17] who used a relaxation method and also included closed magnetic field lines (dead zones).

ANALYTICAL MODELING OF STEADY MHD OUTFLOWS

On the other hand, the construction of analytical solutions of the *steady* MHD equations is a rather difficult undertaking, even when axisymmetry is assumed, regardless of whether the outflow emerges from a hot corona surrounding a spherically symmetric gravitational body, or from an accretion disk. Basically the core of the difficulty lies in the fact that the set of the steady and axisymmetric MHD equations, reduces to a pair of highly nonlinear and coupled differential equations of mixed elliptic/hyperbolic type. Then, from the causality principle, a physically acceptable solution needs to cross several critical surfaces. However, the exact positioning of those critical surfaces is not known *a priori* but it is only determined simultaneously with the solution. It is for this reason that only a few classes of such exact MHD solutions have been studied so far. All available analytical solutions are employing a separation of the variables technique [36]. In the following section we shall briefly outline through two examples what can be done analytically through this technique to study the collimation of winds into jets in the framework of a classical treatment.

It is an observed fact that, besides the case of radio-quiet AGN, extragalactic outflows are relativistic. However, we may reasonably expect that the basic physical mechanisms at work operate also in the classical regime. So, for the sake of mathematical simplification we shall confine ourselves in studying the simpler non-relativistic set of the MHD equations; a generalisation of the results thus obtained to the relativistic case still remains a challenge for future studies.

Exact steady solutions using a nonlinear separation of the variables approach

To fix our notation, we denote the flow density by ρ, velocity by $\vec{V}$, magnetic field by $\vec{B}$ and pressure by P (or equivalently the temperature by T), as functions of the spherical (r,θ,φ) or cylindrical (ϖ,φ,z) coordinates. In fact we need to use interchangeably spherical and cylindrical coordinates, since spherical coordinates are usually adopted for stellar winds starting from a spherical base (corona) while cylindrical coordinates are more appropriate to study disk-wind problems. Another useful set of coordinates are the arclength s along the poloidal direction $\hat{p}$, the normal to the poloidal field line coordinate n along $\vec{\nabla} A / \mid \vec{\nabla} A \mid = \hat{n}$ and the azimuthal angle φ.

Under the assumption of steadiness $\partial / \partial t = 0$ and axisymmetry $\partial / \partial \varphi = 0$, the magnetic field on the poloidal plane is defined by means of a scalar magnetic flux function A, $\vec{B}_p = (\vec{\nabla} A \times \hat{\varphi})/\varpi$, while the momentum equation splits in the poloidal plane into a component along the poloidal flow and a component across this flow. Momentum balance along the poloidal flow may be combined with momentum balance across the flow (the transfield or Grad-Shafranov equation) to form a system of two coupled partial differential equations for the density ρ and the magnetic flux function A. Irrespectively of using a polytropic equation of state between pressure and density, or not, this system contains integrals that depend only on the magnetic flux distribution, such as the mass to magnetic flux ratio, $\Psi_A(A)$, the total angular momentum, $L(A)$, and the angular velocity or rotational frequency of the footpoints of the magnetic fieldlines anchored in the wind source, star or disk, $\Omega(A)$.

With the azimuthal variable φ ignorable, considerable analytical simplification results by assuming a nonlinear separation of the dependance of the physical quantities on the two remaining spherical coordinates, i.e., the poloidal plane coordinates (r, θ). Two main classes of such solutions have been shown to exist wherein all quantities scale with the spherical radius r or the colatitude θ. The two cases correspond to solutions characterized by radial and meridional self-similar symmetries [36]. It is worth to note that this systematic construction unifies all existing analytical models of cosmic outflows, such as the classical Parker [25] and Blandford and Payne [4] models for the solar wind and disk-winds, respectively, in addition to uncovering new and interesting global models [36].

The *first* family with the radially self-similar symmetry is appropriate to winds emerging from disks [4], [21], [22], [14]. No intrinsic scale length exists in this case and all quantities scale as a power law of the radius, similarly to the Keplerian law for the velocity in the disk. The key assumptions in this class of solutions are that the poloidal Alfvén Mach number M and the cylindrical radius ϖ of a particular poloidal fieldline A=*const.* (in units of the cylindrical radius ϖ_a at the Alfvén point along the same poloidal fieldline), are solely functions of the colatitude θ,

$$M = M(\theta) \, , \qquad \frac{\varpi}{\varpi_a} = G(\theta) \; . \tag{1}$$

With these two assumptions one can show that there exist six sets of analytical solutions comprising the large family of radially self-similar MHD outflows [36]. Each of these six sets of solutions corresponds to a judicious choice of the MHD integrals Ψ_A, $L(A)$ and $\Omega(A)$ such that the radial dependance drops out from the MHD partial differential equations which thus become ordinary differential equations in θ. One of these six sets of solutions has been already analysed by [11] and allows a polytropic relationship between pressure and density. This set also contains as special case the [4] classical Blandford & Payne [4] disk-wind model.

The *second* family which we shall examine here in some more detail is characterized by the meridional self-similar symmetry and is appropriate to winds emerging from a spherical source, although the physical variables are not spherically symmetric and the boundary conditions are function also of the colatitude. Even though these solutions seem to be more natural to describe stellar winds, they do not exclude the presence of a surrounding accretion disk. The key assumptions also in this case are that the poloidal Alfvén Mach number M and the dimensionless cylindrical radius ϖ of a particular poloidal fieldline $A{=}const.$ are solely functions of the spherical radius r:

$$M = M(r)\,. \qquad \frac{\varpi}{\varpi_a} = G(r)\,. \tag{2}$$

With these two assumptions one can show that there exist a few sets of analytical solutions comprising the large family of meridionally self-similar MHD outflows. Again, each of these sets of solutions corresponds to a judicious choice of the MHD integrals Ψ_A, $L(A)$ and $\Omega(A)$ such that the meridional dependance drops out from the MHD partial differential equations which thus become ordinary differential equations in r. One of these sets of solutions has been analysed in detail in [29], [30]. The poloidal components of momentum balance give then that the total volumetric forces in the radial and meridional directions have the forms,

$$f_r = a + b\sin^2\theta\,, \qquad f_\theta = c\sin\theta\cos\theta\,, \tag{3}$$

where the expressions a, b and c are functions of the radial distance r, the Alfvén number $M(r)$, the dimensionless pressure $\Pi(r)$, the dimensionless radius $G(r)$ and expansion factor $F(r)$. By setting the total force equal to zero, $F_r = F_\theta = 0$ we obtain three ordinary differential equations which together with the defining relation between $G(r) - F(r)$, $F/2 = 1 - d\ell nG/d\ell nr$ can be solved for the unknown functions $M(r)$, $\Pi(r)$, $G(r)$ and $F(r)$. In this way all physical quantities of a full solution for the outflow from the base to *unlimited* large distances is obtained. Note that in constructing a solution one has to carefully cross the appropriate critical points encountered at the characteristic MHD speeds, in consistency with the assumed self-similarity symmetry [34], [30].

An energetic criterion for the collimation of outflows

In the previous model of the MHD outflow, each streamline is labeled by the dimensionless magnetic flux function A, with $A = 0$ labeling the polar streamline while A increases as we move towards the equator at a given spherical radius r. The specific energy E_o along the streamline A and at a radial distance r is the sum of the kinetic $V_p^2(r,A)/2$ and Poynting $-\Omega(A)\varpi B_\varphi/\Psi_A$ energy flux densities per unit of mass flux density, the gravitational specific energy $-GM/r$ and finally the heat content, or enthalpy $h(r,A)$. If the plasma were adiabatic this specific energy E_o would be precisely the conserved Bernoulli energy. However, in winds a parcel of plasma is always heated on its journey from the base r_o of the particular streamline A to infinity. Then, denote by $\Theta_{r_o}^{r}(A)$ the specific heating received by the unit of plasma mass as it moves from the base r_o to r along $A = const.$, denote by $\Theta_{r_o}^{\infty}(A)$ the total heating from r_o to infinity along $A = const$, while

$$\Theta_r^{\infty}(A) = \Theta_{r_o}^{\infty}(A) - \Theta_{r_o}^{r}(A)\,. \tag{4}$$

By subtracting from E_o the heating that has been available to the gas from the base r_o to the point under consideration r on the given streamline, $\Theta^{r}_{r_o}(A)$, we obtain the familiar Bernoulli specific energy $E(A)$ which is available along the particular streamline. Then, we can define the energy $E_1(A)$ to be the sum of $E(A)+\Theta^{\infty}_{r_o}(A)$,

$$E_o + \Theta^{\infty}_{r}(A) = E(A) + \Theta^{\infty}_{r_o}(A) = E_1(A)\,. \tag{5}$$

Usually in an outflow the thermal input in the form of internal energy and external heating is not all fully converted into other energy forms, unless the terminal temperature is zero, then remains some asymptotic thermal content, $h(\infty, A)$. By subtracting from the previous energy $E_1(A)$ the heat content at infinity, $h(\infty, A)$, we obtain a new streamline constant, $\tilde{E}(A)$, which will be the total convertable specific energy along the given streamline A, i.e., the energy which can be converted to other forms. Finally, the volumetric total convertable energy is $\rho(r,A)\tilde{E}(A)$.

It turns out that in our model the difference of the volumetric convertable energy between a nonpolar streamline and a polar streamline normalized to the volumetric energy of the magnetic rotator $\rho(r,A)L(A)\Omega(A)$ is a constant ϵ' [30],

$$\frac{\epsilon'}{2\lambda^2} = \frac{\rho(r,A)\tilde{E}(A) - \rho(r,\mathrm{pole})\tilde{E}(\mathrm{pole})}{\rho(r,A)L(A)\Omega(A)}\,. \tag{6}$$

This quantity ϵ' plays a crucial role in the asymptotic shape of the streamlines, in the sense that the necessary condition for cylindrical asymptotics is $\epsilon' > 0$. This may be understood from a broader perspective as follows. The gradient of the convertable volumetric energy $\rho(r,A)\tilde{E}(A)$ with a minus sign represents a force in the spirit of Le Chatelier's principle where a system out of several possibilities it has, chooses to evolve towards that equilibrium state which has the minimum energy. In this sense, a streamline will have an asymptotic shape which is in the direction of $\vec{f}$. This force $\vec{f}$ is in the direction of the unit vector $\hat{n} = -\nabla A/\mid\nabla A\mid$,

$$\vec{f} = -\rho(r,A)\vec{\nabla}\tilde{E}(A) \sim \rho(r,A)\frac{\mathrm{d}\tilde{E}(A)}{\mathrm{d}A}\varpi B_p \hat{n}\,. \tag{7}$$

Thus, if f_n is positive, the streamlines focus towards the axis in the direction of $\hat{n}$, On the other hand, asymptotically the streamlines cannot bend towards the equator [15]. Thus, if $f_n \leq 0$ the only possibility left for the streamline asymptotics is to end up radially or paraboloidally.

Altogether, cylindrical collimation is controlled by a single parameter, ϵ', representing the variation between a fieldline A and the pole, of the sum of the (volumetric) poloidal and toroidal kinetic energies, the Poynting flux, the gravitational potential and the converted thermal content. In our specific model, this parameter can be split into two terms, representing the magnetic and the thermal contribution [30]:

$$\epsilon' \equiv \epsilon + \kappa\frac{V_\infty^2}{V_*^2}\,, \tag{8}$$

TABLE 1. AGN classification according to orientation and efficiency of magnetic rotator

Radio-emission	Type 2	Type 1	Type 0	Magnetic Rotator, ϵ
Radio-quiet	Seyfert 2	Seyfert 1, QSO	BAL QSO	Inefficient, $\epsilon < 0$
Radio-loud	FR I		BL Lac	Intermediate efficiency, $\epsilon \gtrsim 0$
Radio-loud	FR II	BLRG/SSRQ	RQ	Efficient, $\epsilon \gg 0$

where V_∞ and V_* are the polar asymptotic and Alfv'en speeds and κ the variation of the pressure across streamlines [30]. The second term represents the variation across the streamlines of the thermal content that is finally converted into kinetic energy and gives a measure of the thermal pressure efficiency to collimate the outflow. For under-pressured flows ($\kappa > 0$) the gradient is inwards and helps collimation. On the other hand, over-pressured jets ($\kappa < 0$) and iso-pressured jets ($\kappa = 0$) can collimate only magnetically.

The energy of the magnetic rotator ΩL is mainly stored in the form of Poynting flux and is the source of (i) acceleration and (ii) collimation for magneto-centrifugal winds. In an equatorial wind, when this energy dominates we have a fast magnetic rotator and the wind is magneto-centrifugally driven. Conversely, when thermal acceleration is dominant the magnetic rotator is termed slow. The first term ϵ in Eq. [8] is a measure of the efficiency of the magnetic rotator to collimate the flow in the present model. This parameter ϵ, which can be evaluated at the base of the flow r_o, is equal to the excess of the magnetorotational energy on a nonpolar streamline which is not used to drive the flow, in units of the energy of the magnetic rotator. In other words, ϵ measures how much of the energy of the magnetic rotator is not used to escape the gravitational well and is available for magnetic collimation alone. If there is an excess of this energy on non polar streamlines, magnetic forces can collimate the wind into a jet. Thus, when $\epsilon > 0$ we have an *Efficient Magnetic Rotator (EMR)* to magnetically collimate the outflow into a jet, and an *Inefficient Magnetic Rotator (IMR)* if $\epsilon \leq 0$ [30].

A PHYSICAL CLASSIFICATION OF AGN

In the previous model we found that the asymptotic morphology of the outflow is controlled by the constant ϵ', characterizing the efficiency of the magnetic rotator and the pressure gradient across the streamlines, Eq. [8]. The first is related to the magnetic properties of the central object in the AGN while the second to the pressure variation across the streamlines. We may assume that this can be somehow related to the environment through which the jet propagates. Thus, we may discuss the following possibilities in the framework of a classification scheme of AGN, as they are summarized in Table 1:

(I). *Inefficient magnetic rotators, $\epsilon \leq 0$, corresponding to radio-quiet AGN of Type 0, 1 and 2 with uncollimated or loosely collimated outflows.* One possibility is that $\epsilon' < 0$ such that the AGN produces a radially expanding outflow. This may happen if the source is in a rich environment such that latitudinally we have an over-pressured outflow, $\kappa < 0$. The other possibility is that $\epsilon' \gtrsim 0$, i.e. ϵ' is marginally positive such that the AGN produces a 'loosely' collimated jet, i.e., collimation occurs slowly at large distances. This may happen, for instance, if we have a latitudinally under-pressured outflow, $\kappa > 0$. Then, if the angle between the jet and the line of sight $\theta \sim 90^o$ we have a Type 2 Seyfert, if $\theta \sim 0^o$ we have a BAL QSO while in the intermediate case of $0 < \theta < 90^o$ a Type 1 AGN. Hence, if the central source is an IMR ($\epsilon \leq 0$) the density drops quite rapidly with the radial distance and this could be related to the weaker outflows in Seyfert 1 and 2 galaxies and radio-quiet QSO's.

(II). *Efficient magnetic rotators, $\epsilon > 0$, corresponding to radio-loud AGN of Type 0, 1 and 2 with well collimated jets.* In this case, since ϵ obtains high positive values, we have a tightly collimated jet, regardless of the value of κ, i.e., regardless if the jet propagates in a rich or poor environment. As before, orientation effects may furthermore distinguish between a Blazar (radio-loud Quasar with $\theta \sim 0^o$), a BLRG, FSRQ, SSRG for the intermediate case of $0 \lessapprox \theta \lessapprox 90^o$ and finally a FR I type of radio-loud galaxy if $\theta \sim 90^o$.

(III). *Intermediate efficiency magnetic rotators, $\epsilon' > 0$, corresponding to radio-loud AGN of Type 0 and 2 with collimated jets.* The two possibilities are either that ϵ is marginally positive and $\kappa < 0$, or, ϵ is marginally negative and $\kappa > 0$. In this case, we always have $\epsilon' > 0$ and hence the outflows always have asymptotically cylindrical flux tubes. Note also that many extragalactic jets, as deduced from X-ray data on the hot surrounding plasma seem to be propagating in rich environments [13]. For example, this seems to be the case with FR I type of RG [26]. Orientation effects may distinguish between BL Lacs ($\theta \sim 0^o$) and FR I radio galaxies ($\theta \sim 90^o$).

Note that although radio-quiet AGN are believed to have nonrelativistic outflows, superluminal effects in radio-loud AGN are taken as an indication that these outflows are relativistic. On the other hand, despite that the model on which our conclusions are based is clearly nonrelativistic, its basic trends have been shown to be preserved in relativistic cases as well [12].

If the strength of the magnetic rotator reduces, one expects a smooth transition from a jet to a loosely collimated wind and finally to a radial wind. This would correspond to moving from the radio-loud quasars and Blazars of Table 1 to the radio-quiet Seyfert galaxies and QSO's.

CONCLUDING REMARKS

Magnetic collimation of winds into jets appears to be a rather general property of the MHD equations governing plasma outflows. To this conclusion one arrives either through general asymptotic analyses of nonrelativistic MHD, [15] and its relativistic generalisation in [12], or, by the construction of exact jet-type MHD solutions by the method of nonlinear separation of variables, [30], [37]. Those studies show also that pressure gradients could contribute to confine the outflow. The transformation of magnetocentrifugal energy into kinetic energy seems to be the most natural driving mechanism for outflows from accretion disks, but close to the rotational axis also a contribution from the thermal energy is required, with appropriate heating processes occurring in the plasma.

The MHD acceleration/collimation mechanisms can work in very different astrophysical scenarios, whenever we have a rotating magnetized body such as a supermassive black hole surrounded by an accretion disk. In extragalactic jets and on scales of several Kpc, pressure confinement by the environment seems to be present, while close to the center the jet may be either magnetically or thermally confined. Even though the basic thermal/magnetic driving and confining mechanisms discussed here should be qualitatively valid also for relativistic velocities, a detailed modeling of such relativistic jets from AGN is needed at this point. Relativistic MHD solutions have been obtained either by integrating the Bernoulli equation only along a streamline, or by assuming a radial self-similar configuration (see e.g. [20], [8]). In both cases it is shown the possibility of acceleration and collimation of the beams to highly relativistic bulk velocities, as required by the observed superluminal motions and Doppler beaming. However, since these self-similar solutions are subcritical with the related causality problem, then a more consistent modeling of jets in AGN requires crossing all MHD singularities [34].

The standard unification scheme of AGN that some classes of objects transform into other classes with the viewing angle, seems to be secure by now. And, in this article we added that there is a physical criterium separating the various classes among themselves. At the same time however, it must be explained why recent data at optical and X-ray wavelengths have shown that, at least for the FR I/BL Lac case, the standard unification model does not seem to be in full agreement with observations [9], [10]. In fact, it seems unavoidable that in order to fit the data, a structure of the velocity across the jet must be assumed.

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