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Theory of magnetized accretion-ejection structures

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PoS(HEPROVIII)004

Accretion-ejection are interdependent processes linking jet acceleration and collimation physics to the underlying accretion disk physics. In these systems, a large scale vertical magnetic field is assumed to thread the accretion disk, leading simultaneously to jet formation and the onset of a magnetic turbulence inside the disk, both inducing accretion. While the first analytical models have been published more than 25 years ago, the global understanding has constantly progressed, showing the dominant role of jets or winds in driving accretion.

The concept of Jet Emitting Disk (JED) and Wind Emitting Disk (WED) has emerged in the theoretical side, while in the computational side configurations such as Magnetically Arrested Disk (MAD) and Standard And Normal Evolution (SANE) have been the focus of much attention. A direct comparison between these costly 3D numerical experiments and steady-state theory has finally become feasible.

After describing the physics and general properties of JED/WED accretion-ejection configurations, I will argue that they provide the state-of-the-art mathematical description of their numerical counterparts, MAD/SANE. More efforts need however to be done in order to firmly assess this point.

In any case, these two complementary approaches have unveiled the critical role played in astrophysical systems by the radial distribution and its temporal evolution of the local disk magnetization. Magnetic field dragging in accretion disks appears therefore to be the key ingredient allowing to understand hybrid disk configurations and outbursting cycles, such as those seen in X-ray binaries.

High Energy Phenomena in Relativistic Outflows VIII (HEPROVIII)
23-26 October, 2023
Paris, France

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1. Introduction

The existence of jets is commonly observed in a number of astrophysical objects, going from non-relativistic sources such as young stellar objects (YSO) or post-AGB binaries, to relativistic objects such as galactic X-ray binaries or AGN. It has soon been understood that launching bipolar supersonic outflows, with opening angles of a few degrees only, requires the action of a large scale magnetic field anchored into an astrophysical rotating source. By this way, jet acceleration and jet collimation become two interdependent processes, the ejected plasma being channeled by a magnetic funnel and shaping it in return.

Magnetohydrodynamics (hereafter MHD) is key to understand jet physics. It is useful to think of asymptotic jets as cylindrical "screw pinches" studied in thermonuclear fusion devices, namely with both vertical and toroidal magnetic field components. While the vertical magnetic field is generally assumed to be brought in by the plasma accreting onto the central object, the toroidal magnetic field depends on the interplay between the inertia due to the ejected plasma and the rotation of the object the magnetic field is anchored to. The toroidal magnetic field is actually the cornerstone of jet physics: (1) it determines the MHD Poynting flux powering the jet; (2) it is responsible for the hoop-stress that eventually leads to plasma self-confinement and, last but not least, (3) its magnitude at the source leads to its spin down (magnetic braking), as a result of Lenz's law. It is an understatement to say that any simplifying assumption on the toroidal field component near the source will lead to major uncertainties on jet physics.

The toroidal magnetic field can also be seen as being generated by a poloidal electric current, flowing downward (towards the source) along the axis and closing[1] thanks to a return current flowing along the jet itself and/or along its interface with the ambient medium (see Fig.1). The existence of such a poloidal electric current relies on an electromotive force (e.m.f), which arises whenever a rotating conductor is embedded in a vertical magnetic field. Two independent e.m.f are therefore expected, one due to the disk itself and the other due to the rotating central object (black hole or star).

Magnetized jets launched from accretion disks have been proposed by Lovelace [1] and Blandford [2, 3] and are nowadays referred to as BP-jets. These jets tap the rotational energy of the orbiting plasma, allowing for accretion and are thus fed by the released accretion power. The other kind of jets, launched from the BH ergosphere and tapping the rotational energy of the BH, have been proposed by Blandford & Znajek [4] and are referred to as BZ-jets. They have received much attention these last 20 years thanks to the outcome of supercomputers.

It is common to have a favorite model, some people advocating BP process while others preferring BZ process to explain astrophysical jets. However, it should be reminded that (1) both processes require the existence of a large scale vertical field threading the accretion disk; (2) the two e.m.f are actually independent; (3) the jet axis is naturally filled-in by the BZ outflow, while BP-jets naturally account for outflows launched beyond a few gravitational radii r_g. So it seems quite plausible to expect both processes to be at work, namely a BZ relativistic spine surrounded by a sub- or mildly-relativistic magnetized BP-outflow.

Section 2 will be devoted to BP-jets and in particular to the only semi-analytical model published so far that relates disk physics to jet physics. Explaining astrophysical jets by relying

[1] The electric current must be divergence free in steady-state.

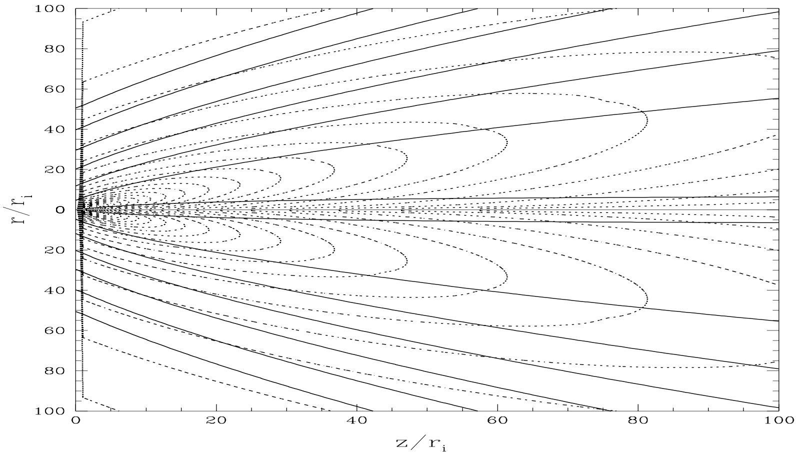

Figure 1: Poloidal magnetic field lines (solid lines) of a cold magnetically-driven jet emitted from a near-keplerian accretion disk [13]. The disk (on the y axis) is too thin and cannot be seen with the vertical scale used here. The dotted lines display the poloidal electric current density. The electric circuit displays a butterfly-like shape, characteristic of a tenuous ejection with $\xi < 1/2$ (see Sect. 2). The current responsible for the collimating hoop-stress flows down along the jet axis, enters the disc at its inner edge r_i and returns along the jet itself. This current configuration is highly dependent on the magnetic field distribution in the disk (see [19] and references therein).

only on accretion disks allows for a universal explanation valid for any central object which is not a BH, such as a forming star, a white dwarf or a neutron star. Moreover, it is noteworthy that even in the case of a BH, there is a recent evidence in M87 that observable jets very close to the source are already much wider than expected from BZ-jets *as obtained from current GRMHD simulations* (see [5, 6]). It is therefore unclear yet if observations of astrophysical jets do actually provide any hint on the BZ process. Nevertheless, even in the eventuality that BP-jets were to dominate BZ jets, the issue of the unavoidable dynamical spine-jet interplay would remain open.

Section 3 is focussed on how global 3D numerical simulations have deeply modified our view of accretion-ejection and how it became nowadays possible to understand them. But this requires digging into the outcome of high resolution, converged simulations and thoroughly comparing them to the semi-analytical models. Unfortunately, this is something very rarely achieved. Nevertheless, I will show that what is popularly referred to as a numerical Magnetically Arrested Disk (hereafter MAD) or a numerical Standard And Normal Evolution (hereafter SANE), can be related to the semi-analytical models of Jet Emitting Disk (hereafter JED) and Wind Emitting Disk (hereafter WEDs). Both situations are actually the numerical and its semi-analytical counterpart of accretion disks at large (MAD/JED) or low (SANE/WED) disk magnetization. I will finally conclude on the critical role played in astrophysical systems by the local disk magnetization.

2. Magnetized Accretion-Ejection Structures

2.1 Assumptions and governing equations

It is assumed that a large scale vertical magnetic field is threading the accretion disk, whose dynamics is mostly controlled by the gravity of the central object. The model assumes a newtonian potential since it is valid for most sources and, for black holes, relativistic effects would be important anyway only below $\sim 10r_g$. As a consequence of newtonian gravity, which is self-similar, it is natural to seek for a mathematical description of accretion and ejection processes sharing that property. This translates into looking for power-law solutions of the cylindrical radius r, namely $A(r,z) = A_o(r/r_o)^{\alpha_A} f_A(z/r)$ for any quantity A, where α_A is the radial exponent, r_o any fiducial radius and $f_A(z/r)$ an unknown function of the variable z/r. As a consequence of this modeling, the *disk aspect ratio* defined as $\epsilon = h/r$, where $h(r)$ is the local isothermal hydrostatic half-thickness, must be a constant of the radius. Actually, any dimensionless parameter is automatically a constant of the radius within this framework. Any physical quantity A can then be analyzed using a function $f_A(x)$, where $x = z/h$ is the self-similar variable [7].

The origin of the vertical magnetic field is not addressed by the theory and it is simply assumed that the steady-state disk magnetic flux follows $a(r,z) = a_o(r/r_o)^{\beta}\Psi(x)$. Obviously, this exponent β must be consistently related to other disk quantities. The local strength of the vertical magnetic field is defined at the disk midplane with the *disk magnetization*

$$\mu = \frac{V_A^2}{C_s^2} = \frac{B_z^2/\mu_o}{P_{tot}} = 2\frac{P_{mag,z}}{P_{tot}} \tag{1}$$

where $P_{tot} = P_{gas} + P_{rad} = \rho C_s^2$ is the total (kinetic + radiation) pressure, C_s the local sound speed and μ_o the vacuum permeability. Note that this definition includes the radiation pressure and is therefore more general than the widely used beta plasma. More importantly, it includes only the vertical *laminar* component, neither the other two laminar components B_r and B_ϕ, nor the turbulent magnetic pressure. This is because all these contributions are actually outcomes of the vertical magnetic field and μ is the real control parameter for the disk magnetic field strength.

Since symmetric bipolar jets or winds (same mass loss from both disk surfaces) are expected, the steady-state disk accretion rate must vary with the radius [7]

$$\xi = \frac{d \ln \dot{M}_a(r)}{d \ln r} \tag{2}$$

where $\dot{M}_a(r) = -\int 2\pi r \rho u_r dz$ is the local disk accretion rate and ξ is the *disk ejection efficiency*[2]. Providing the range of allowed ejection efficiencies ξ as function of the disk properties, in particular $\xi(\mu)$, is the ultimate goal of the theory.

There is an important prescription that needs to be made about field diffusion in accretion disks. In strict steady state, mass needs to accrete onto the central object while the magnetic field remains behind. This can only be done if some very efficient magnetic diffusivity is present in the disk. In fully ionized plasmas, such a diffusivity can only be provided by some self-sustained turbulence. In the very same spirit as Shakura & Sunyaev [8], that assumed the existence of a turbulent viscosity $\nu_v = \alpha_v C_s h$, Ferreira & Pelletier [7] proposed the existence of an anomalous magnetic diffusivity $\nu_m = \alpha_m V_A h$ in the poloidal direction (acting on the electric current density J_ϕ) as well as a toroidal magnetic diffusivity $\nu'_m = \nu_m/\chi_m$ (acting on the poloidal current density), where χ_m is a measure of a possible anisotropy. Since the seminal work of Balbus & Hawley, it is now believed that an MHD

[2]The Standard Accretion Disk (hereafter SAD) solution of Shakura & Sunyaev [8] assumed $\xi = 0$, but also did Narayan & Yi [9] or even Blandford & Payne [3].

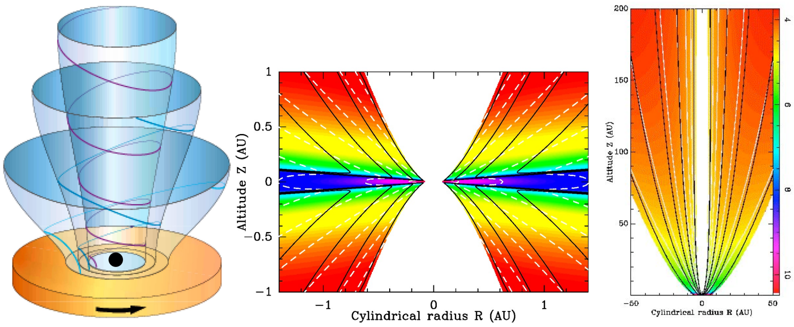

Figure 2: Left: Sketch of axisymmetric magnetic surfaces nested around each other, showing twisted magnetic field lines. Middle: Close-up view of the disk, the density in background, streamlines in black, isocontours of total velocity in white. Right: View of one jet, using the same conventions [13].

turbulence is indeed present and self-sustained by the magneto-rotational instability (hereafter MRI, [10]). So, the theory of steady-state accretion-ejection structures requires to specify at least three dimensionless numbers

$$\alpha_m = \frac{\nu_m}{V_A h} \quad , \quad \chi_m = \frac{\nu_m}{\nu'_m} \quad , \quad \mathcal{P}_m = \frac{\nu_v}{\nu_m} \tag{3}$$

where $\mathcal{P}_m$ is the (effective) magnetic Prandtl number and believed to be of order unity in fully developed turbulence. For the same reason, one should expect χ_m of order unity in 3D turbulence, namely no particular direction and thus no relevant anisotropy in the magnetic field diffusion. This leaves mostly α_m as a free parameter, related to the strength of the MHD turbulence. Remarkably, these relations provide $\alpha_v = \alpha_m \mathcal{P}_m \mu^{1/2}$, which is quite exactly the dependency found in local (shearing box) simulations of MRI-driven turbulent gas-supported disks [11]. However, once the outflowing plasma leaves the accretion disk, such a turbulence must decay and ideal MHD regime then prevails. Ejected plasma is then assumed to be described by ideal MHD and all transport coefficients are set to zero. This requires the prescription of self-similar profiles $f_v(x)$ that decay with height. As a first uneducated guess, these profiles have been chosen identical and a gaussian [12].

To summary, the governing equations that need to be solved are

- Mass conservation

$$\nabla.\rho\mathbf{u} = 0 \tag{4}$$

- Momentum conservation

$$\rho\mathbf{u}.\nabla\mathbf{u} = -\nabla P - \rho\nabla\Phi_G + \mathbf{J}\times\mathbf{B} + \nabla.\mathbf{T} \tag{5}$$

- Ohm's law and toroidal field induction

$$\eta_m J_\phi \mathbf{e}_\phi = \mathbf{u_p}\times\mathbf{B_p} \tag{6}$$

$$\nabla.(\frac{\nu'_m}{r^2}\nabla r B_\phi) = \nabla.\frac{1}{r}(B_\phi \mathbf{u}_p - \mathbf{B}_p\Omega r) \tag{7}$$

where ρ is the density of matter, P the total pressure, $\Phi_G = -GM_*/(r^2+z^2)^{1/2}$ the gravitational potential, $\mathbf{J} = \nabla\times\mathbf{B}/\mu_o$ the current, $\mathbf{T}$ the "viscous" stress tensor (Shakura & Sunyaev), $\eta_m = \mu_o \nu_m$ the anomalous electrical resistivity. The missing energy equation is not reproduced here as several simplifying approximations can be (and have been) used: isothermal magnetic surfaces [13], adiabatic magnetic surfaces [14] or isothermal disk with and additional heat input at the surface leaving them farther up adiabatic [15].

The complete set of PDE can be exactly solved using the method of variable separation within the self-similar ansatz. We stress that, by doing so, *all terms have been kept in the equations*, there is therefore no approximation whatsoever: the 2D solutions shown in Fig.(2) are exact MHD solutions of the problem. The set of PDEs becomes a set of algebraic relations between the exponents α_A plus a set of ODEs with the functions $f_A(x)$. This set of ODEs can be written as $M\cdot Y' = P$, where M is a matrix and P a vector defined on x and the functions f_A while Y' is a vector of the derivatives of the functions f_A. Integration from the disk midplane $x=0$ to "infinity" ($x\longrightarrow\infty$) can then be done with a predictor-corrector numerical scheme for stiff equations, once initial conditions are assumed. This requires however the matrix M to be invertible, which is not the case whenever the ejected plasma achieves a velocity equal to the phase speed of the relevant wave. This critical velocity corresponds to the sound speed in the turbulent disk, but this is a situation that never occurred: starting at $x=0$ in an (anomalous) resistive MHD regime, plasma is being accreted and lifted up until it reaches the ideal MHD regime [12]. Once in that regime, plasma is forced to follow the magnetic surfaces ($\mathbf{u_p}\times\mathbf{B_p}=0$) and the critical speeds become the slow-magnetosonic speed (SM, near the disk surface), the Alfvén point (A, farther away) and the fast-magnetosonic speed (FM, much further out [13, 16]).

Each smooth crossing of a critical point requires a regularity condition that determines one parameter of the model. The resultant parameter space discussed below is thus obtained for a set of 4 free parameters: $(\epsilon, \alpha_m, \chi_m, \mathcal{P}_m)$. For each set, all possible values of the disk magnetization μ are scanned leading to the determination of $\xi(\mu)$. Note that since the disk aspect ration ϵ is free and all terms kept in the equations, the solutions are valid for all types of disks, from geometrically thin $\epsilon \ll 1$ to slim ($\epsilon\sim 0.1$) or even thick disks ($\epsilon > 0.3$).

2.2 Main properties

A strictly steady-state magnetized accretion-ejection structure requires that if $\dot{M}_a \propto r^{\xi}$, then $B_z \propto r^{\xi/2-5/4}$ (or $a\propto r^{3/4+\xi/2}$), $\rho \propto r^{\xi-3/2}$ (or $\Sigma = \int \rho dz \propto r^{\xi-1/2}$) and $I = \int 2\pi r J_r dz \propto r B_\phi \propto r^{\xi/2-1/4}$. It can be readily seen that the value $\xi = 1/2$ plays some interesting role: at $\xi < 1/2$, the disk column density is a decreasing function of the radius and the poloidal electric current *flows out* of the disk from both its surfaces (see Fig. 1), whereas for $\xi > 1/2$, the disk gets denser at larger radii and the poloidal current *flows into* the disk from its surfaces. Since $div\mathbf{J} = 0$, the electric circuit must be closed[3]. So, the low-ξ case has a return current flowing within the jet itself while the jet-confining current, directed towards the disk, must be flowing in along or near the jet axis. On the contrary, the large-ξ case pushes the return current to the outer jet-ambient medium interface.

Which situation is more probable or more stable for astrophysical jets requires to deal with both radial boundary conditions (jet axis and jet boundary) and is of course beyond the reach of

[3]Note that the disk e.m.f gives actually rise to two independent electric circuits, each related to one jet.

PoS(HEPROVIII)004

self-similar solutions (see [19] for numerical simulations). However, they can provide under which conditions the disk allows one situation or the other to occur. The following general properties have been obtained using $\mathcal{P}_m = 1$ and gaussian profiles for the turbulent transport coefficients:

- No solution has been found for a disk thicker than $\epsilon \sim 0.3$. The reason is because both thermal pressure gradient and magnetic field tension combine to significantly lower the disk rotation to sub-keplerian speed and no magneto-centrifugally driven jet can then be launched [14]. Much less powerful thermally-driven winds can however still be produced, in line with ADAF-like solutions.

- The MHD turbulence level must be quite high, with $\alpha_m \geq 1$, otherwise the vertical speed at the disk surface, which is $\propto \nu_m/h \sim \alpha_m V_A$, is too small and the flow cannot become super-SM [13, 17]. Note that such high value for α_m is actually consistent with MRI.

- The torque exerted by the jets or the winds on the disk is always dominant wrt the turbulent torque, their ratio scaling as $1/\epsilon$. This is because each torque roughly scales as $\mu^{1/2}$, leading to a mass-weighted accretion speed also scaling as $\mu^{1/2}$ [14, 17]. As a consequence, supersonic accretion can be achieved at near-equipartition field strength ($0.1 < \mu < 1$) [12, 13].

- For cold magnetically-driven outflows with negligible enthalpy the typical disk ejection efficiency can be as low as $\xi \sim 0.01$ (see Fig.(3) and below). Whenever some additional heat input is deposited at the disk surface, leading to warm magneto-thermal outflows, disk ejection efficiency as high as $\xi \sim 0.5$ becomes possible [15]. Assuming crudely a Bernoulli invariant $E = (\Gamma_\infty - 1)c^2$, one gets that the maximum asymptotic Lorentz factor along a streamline anchored at a radius r_o writes

$$\Gamma_\infty \simeq 1 + \frac{1}{2\xi}\left(\frac{r_o}{r_g}\right)^{-1/2} \tag{8}$$

  showing that cold magnetically-driven jets from keplerian accretion disks can in principle reach relativistic speeds.

- All self-similar jet or wind solutions obtained so far undergo a recollimation towards the jet axis, due to the hoop-stress [13]. This important property, which should be generic whenever $\xi < 1/2$ according to Contopoulos & Lovelace [18], turns out to not be a bias of self-similarity as it is also observed in 2D simulations of jets [19]. The observational consequences of intrinsic MHD recollimation remains to be investigated (see e.g. [20]).

2.3 JEDs and WEDs

Magnetized accretion-ejection structures describe both the disk and its two jets as a single, mathematically connected, system. The smooth crossing of the Alfvén point determines the position of this critical point as well as the strength of the toroidal field at the disk surface. This, in turn, fixes not only the torque allowing accretion, but also the vertical magnetic compression acting on the disk, hence the disk magnetization μ that is consistent with the disk ejection efficiency ξ (which determines the radial distributions) and the smooth crossing of the SM critical point.

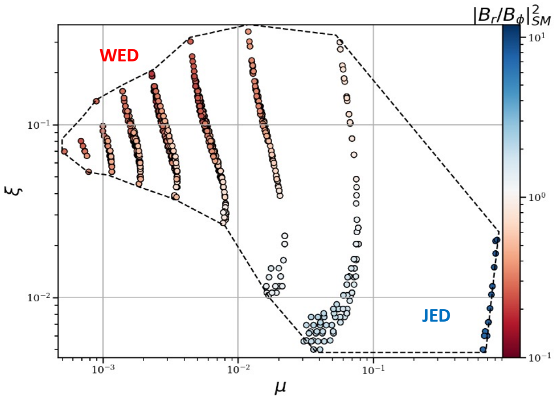

Figure 3: Parameter space $\xi(\mu)$ of isothermal accretion-ejection structures for $\epsilon = 0.1, \alpha_m = \mathcal{P}_m = \chi_m = 1$. The transition from WEDs (with magnetic tower-like winds) to JEDs (with centrifugally-driven jets) occurs roughly around μ a few times 10^{-2} [17].

The first solutions were found for isothermal [13] or adiabatic [14] cold outflows. The disk magnetization μ was found to lie in a very narrow range, between 0.1 and 0.8, with a disk ejection efficiency $\xi \sim 0.01$ leading to very fast and tenuous jets. The disk associated with these solutions has thus been termed *Jet Emitting Disk* [21]. In a JED, accretion is supersonic and most of the released accretion power P_{acc} is shared between the power feeding the two jets $2P_{jet}$ and the power advected onto the central object P_{adv} ($\sim \epsilon^2 P_{acc}$). Indeed, the global energy budget writes [22, 23]

$$P_{acc} = 2P_{jet} + P_{adv} + P_{rad} \tag{9}$$

and it turns out that $b = 2P_{jet}/P_{acc}$ can be close to unity for thin disks, leading to weakly dissipative and thus "radiatively inefficient" JEDs ($P_{rad} << P_{acc}$). As the disk thickens so does the disk rotation rate and the fraction b decreases until no cold jet solution can be found anymore.

What determines in JEDs the maximum value of the vertical magnetic field is the disk vertical equilibrium, not the radial one. Despite a near-equipartition magnetic field ($\mu \sim 1$), rotation remains sub- but near-keplerian as long as the disk remains thin or slim. But even in the case of a very thin disk, the magnetization μ can never be larger than unity. For larger μ, the vertical compression of the disk becomes too strong and breaks down the magneto-hydrostatic balance [12].

On the contrary, it was unclear until recently what determined the minimum value of the disk magnetization. At small μ, MRI sets in despite the presence of anomalous diffusivities and solutions display spatial (vertical) oscillations. Such oscillations are actually non-linear channel modes due to the MRI [17]. Despite a smaller magnetic field, these oscillations help to drive the outflow by building up a stronger toroidal field at the disk surface. Since the vertical pinching of the disk is smaller, these solutions allow for denser cold outflows with a typical ejection efficiency $\xi \sim 0.1$ (see Fig.3). In this case, the denser outflow is also much slower, so the disk associated with these solutions has been termed *Wind Emitting Disk*. In WEDs, accretion is always subsonic and the magnetic configuration is reminiscent of the so-called "magnetic towers".

The reason why WEDs were discarded in previous semi-analytical studies was precisely their oscillatory behavior. As discussed in Jacquemin-Ide et al [17], it is doubtful that oscillations would survive within the disk since secondary instabilities (such as Kelvin-Helmholtz or Rayleigh-Taylor) would most probably be triggered, leading to enhanced anomalous transport within the turbulent region. That is to say that low-μ solutions require to be computed with a non-gaussian vertical profile of the transport coefficients. This is a good incentive to look at global 3D numerical solutions where turbulence is fully resolved.

3. Global 3D MHD simulations

3.1 GRMHD simulations: MADs and SANEs

I won't make a review of all works devoted to GRMHD simulations of magnetized accretion onto BH (see Yosuke Mizuno and Krzysztof Nalewajko's contributions). I will focus instead on differences and similarities with the previous semi-analytical framework.

For some obscure reason, simulations are nowadays classified according to the MAD/SANE terminology. This terminology goes back to the work of Narayan [24] and Igumenshchev [25], who found that if the magnetic field was too large, the disk would not rotate anymore and be therefore "magnetically arrested" (MAD). On the contrary, whenever the field was small enough so that rotation remains nearly keplerian, the numerical outcome has been termed "normal" (SANE), whatever that means. The transition between one state to another depends mostly on two things: (1) initial conditions (how much initial magnetic flux is available in the computational domain) and (2) how long the simulation lasts (since the flux appears to be advected inwards, leading to a growth of the magnetization of the inner regions).

Not only the understanding of these numerical states remains to be done, but it turns out that MADs are actually rotating, turbulent and launching MHD outflows, facts that heavily question the chosen terminology [26, 27]. What seems quite clear however, is the major role played by the magnetic field strength (measured in the disk by μ). It is therefore tempting to associate numerical MADs and SANEs to JEDs and WEDs, respectively. However, the qualification of the disk state (MAD or SANE) in simulations is done by measuring the amount of normalized magnetic flux ϕ_{BH} onto the BH, *not by looking at the local disk magnetization* μ [28, 29]. Nevertheless, steady state JEDs and WEDs both verify

$$B_z r_g^2 \propto \mu^{1/4} \sqrt{\dot{M}_a r_g^2 c} \tag{10}$$

so it seems normal to find that the normalized magnetic flux ϕ_{BH} reaches a constant value once the innermost disk itself has reached a constant magnetization μ. But the radial distribution of μ is never provided in simulations. When the disk magnetic field strength is provided, it is usually done through the beta plasma, at best computed at the midplane, but almost always including also the turbulent magnetic pressure (see however [30]).

The problem with current descriptions of GRMHD simulations, is that studies have been mostly focused on the BH close environment (below $20r_g$) and on the BZ-process, whose axial outflows are called by definition "jets", whereas those coming out of the disk are called "winds" and barely analyzed. Because of this bias, the comparison between simulations around BH and accretion-ejection theory remains to be done.

Nevertheless, a comparison with current MAD simulations is numerically challenging, since it would require to analyse converged simulations beyond say, $10r_g$ or so. Below this radius, spacetime dragging and the necessity that the accretion speed reaches the speed of light give birth

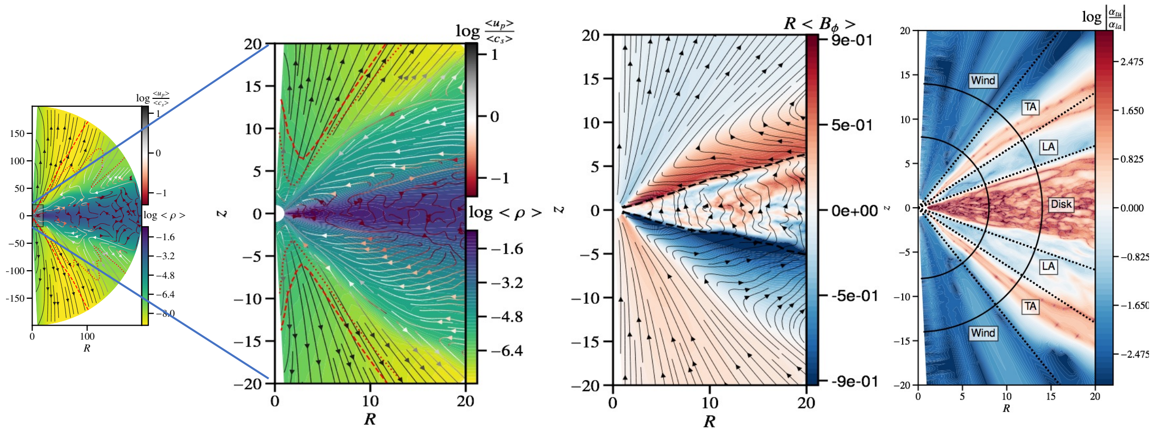

Figure 4: Global 3D simulations of a WED with initial $\mu = 10^{-4}$, only the innermost regions achieved a steady-state for the duration of the simulation. Left panels: background is density, colored lines are streamlines normalized to sound speed. Third panel: solid lines are poloidal field lines, background electric current, dashed line shows $z = h(r)$. Right: the four distinct zones (turbulent disk, laminar atmosphere, turbulent atmosphere, wind) are shown [31].

to the so-called plunging region. In that region, the inevitable plunging of the plasma leads to a dramatic magnetic field dragging producing radially oriented poloidal magnetic field lines. These are thus able to efficiently transfer angular momentum outwardly in the radial direction, connecting the near horizon zone (but beyond the disk FM point) to the region lying outside [32]. This boundary effect, which is obviously out of reach of self-similar solutions, is expected to give rise to an energy input (BH rotational energy) in addition to the local release of binding energy.

3.2 Newtonian simulations: JEDs and WEDs

Newtonian simulations have focused on understanding accretion-ejection processes. They suffer however from another bias which is related to MRI (and numerical difficulties). Since accretion was believed to be mainly due to turbulence and jets/winds only an epiphenomenon, most MHD simulations were done at very weak initial disk magnetization ($\mu \sim 10^{-4}$ or less), a regime only relevant to WEDs.

The first truly global 3D simulation was performed by Zhu & Stone [33], who obtained results that puzzled them: while the disk had a thermal scale height of $h = 0.1r$, MRI-driven turbulence lead the disk to become puffy, with transonic accretion taking place at the disk surface ($z \geq h$) and super-FM winds launched beyond $z \sim r$. It turns out hat puffy disks seems to be a generic situation also in MHD simulations of weakly magnetized disks around BH [34, 35].

This situation was reproduced and thoroughly analyzed by Jacquemin-Ide et al [31] and is displayed in Fig.(4). At low-μ, MRI-driven turbulence lifts plasma vertically out of the disk, in a laminar atmosphere devoid of turbulence (too strong stabilizing magnetic tension). However, the field is too weak there to launch an outflow and material falls in towards the central object, dragging in the field lines. This leads to the generation of a large toroidal magnetic field component allowing MRI to be reactivated. This levitating, near-equipartition turbulent upper atmospheric layer, located around ~ 6 and $9h$ looses angular momentum both sideways: below, by transferring it back to the

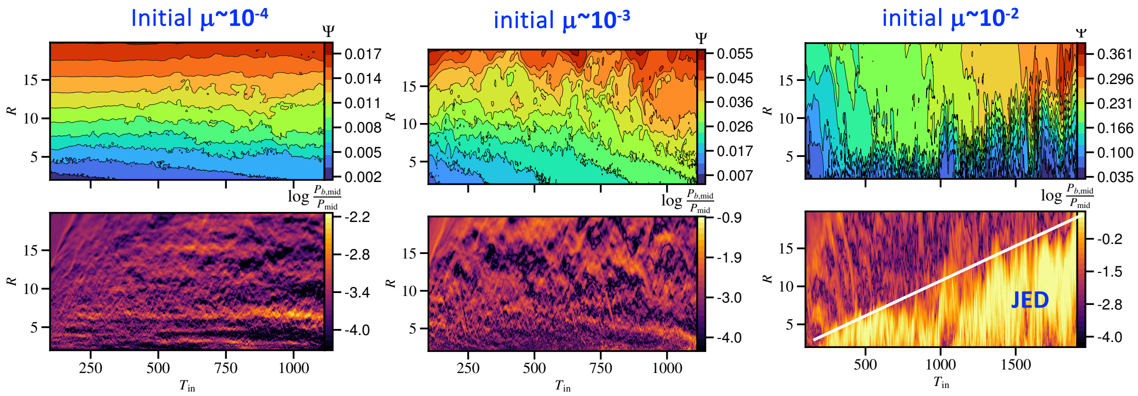

Figure 5: Magnetic field dragging in WEDs and JED generation in the innermost disk regions. Top panels: colors show how magnetic flux is advected inward as function of time (in inner keplerian orbital time). The advection speed gets stronger as μ increases. Bottom panels: maps of log of disk magnetization as function of time. The white line for initial $\mu = 10^{-2}$ shows approximatively the growth of the inner JED [31].

turbulent disk and above, by launching super-FM winds. As a result, accretion there is supersonic as in a JED. This physical picture shows that WEDs are structurally very different from SADs or, to express it differently, *there is no such thing as a SAD when a vertical large scale magnetic field is present*: the disk is always puffy and winds (aka dense and slow super-FM outflows) are unavoidable.

The authors also report another very important result related to magnetic field dragging. As previously realized in the literature, advection of the vertical field is not directly dependent on density but on accretion speed [36], while outward diffusion depends on the toroidal electric current density, hence on the vertical scale of magnetic flux variation [37]. In a WED, that scale is $z \sim r$ and the upper layers are accreting very fast. As a consequence, the field is always advected inwards at a speed very similar to the mass weighted accretion speed, proportional to $\mu^{1/2}$. This implies therefore an almost runaway process, which has been indeed observed in simulations (see Fig.5): starting with a WED initially with $\mu \sim 10^{-2}$, the system converged very quickly to an innermost JED with $\mu \sim 1$. As in numerical MADs, the size of the inner JED region keeps on increasing as long as more magnetic flux is brought in from the outer regions.

4. Final remarks

Modern computational power allows deep investigations of magnetized accretion disks physics. A direct comparison between these costly 3D numerical experiments and steady-state theory has finally become feasible. It appears that mathematical JED/WED solutions reproduce most properties of numerical MAD/SANE configurations. However, there are many diagnostics from the simulations that have not been provided yet, forbidding thorough comparisons. Among those, the radial profiles of the disk magnetization μ and disk accretion rate $\dot{M}_a$ (going beyond $10r_g$) would allow to derive the local dependence $\xi(\mu)$, as well as check the consistency of the local ejection efficiency ξ with the jet MHD invariants. Another quite important feature is the turbulence related parameters, such as α_m for instance, or the magnitude of the turbulent magnetic field with respect to the vertical laminar field. Indeed, numerical experiments show that the turbulent magnetic pressure plays a major role in shaping the disk vertical balance, while it has been neglected so far in analytical studies. Although this work is currently under progress (Zimniak et al, submitted), a thorough analysis of MHD turbulence (values, profiles) in global 3D simulations must be undertaken.

These two complementary approaches have unveiled the critical role played in astrophysical systems by the radial distribution and its temporal evolution of the disk magnetization $\mu(r,t)$. Indeed, not only it is expected to adapt (on accretion time scales) to any variations imposed at the disk outer regions, but it defines also the possible existence of hybrid disk configurations, with the emergence of a magnetically saturated JED state in the innermost regions, as proposed for X-ray binaries [21, 22, 38]. This is a very promising route as such a process should be generic to any accretion-ejection system.

Acknowledgments

This long-term work on the theory of accretion-ejection benefited from the highly valuable contributions of Guy Pelletier, Fabien Casse, Pierre-Olivier Petrucci, Gilles Henri, Jonatan Jacquemin-Ide, Geoffroy Lesur, Grégoire Marcel, Thomas Jannaud, Nathan Zimniak. I acknowledge financial support from the CNES French space agency and PNHE program of French CNRS.

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