The role of jets in the formation of planets, stars, and galaxies

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1.0.1 Abstract

Astrophysical jets are associated with the formation of young stars of all masses, stellar and massive black holes, and perhaps even with the formation of massive planets. Their role in the formation of planets, stars, and galaxies is increasingly appreciated and probably reflects a deep connection between the accretion flows - by which stars and black holes may be formed - and the efficiency by which magnetic torques can remove angular momentum from such flows. We compare the properties and physics of jets in both non-relativistic and relativistic systems and trace, by means of theoretical argument and numerical simulations, the physical connections between these different phenomena. We discuss the properties of jets from young stars and black holes, give some basic theoretical results that underpin the origin of jets in these systems, and then show results of recent simulations on jet production in collapsing star-forming cores as well as from jets around rotating Kerr black holes.
1.1 Introduction

The goal of this book, to explore structure formation in the cosmos and the physical linkage of astrophysical phenomena on different physical scales, is both timely and important. The emergence of multi-wavelength astronomy in the late 20th century with its unprecedented ground and space-based observatories, as well as the arrival of powerful new computational capabilities and numerical codes, have opened up unanticipated new vistas in understanding how planets, stars, and galaxies form. Galaxy formation has turned out to be a complex problem which requires a deep understanding of star formation on galactic scales and how it feeds back on galactic gas dynamics. The discovery of significant numbers of massive black holes in galactic nuclei now suggests that most galaxies harbour million to billion mass holes. The formation of these monsters and the effects that they can have on the early evolution of galaxies is still not well understood however. Star formation studies have made huge inroads, but must still take a large step before we can truly understand the origin of stellar masses and star formation rates in molecular clouds and galaxies. The newly discovered planetary systems bear little relation to text-book models of how our solar system is believed to have formed. It is also abundantly clear that planet and star formation are intimately coupled through the physics of protostellar disks that are their common cradles.

Astrophysical jets play an important role in the formation and evolution of stars, black holes, and perhaps massive planets. They appear to be an inescapable multi-scale phenomenon that arises during structure formation. The enormous kinetic luminosity of many jets (often comparable to the bolometric luminosity of the central sources) also implies that they could have important feedback effects on structure formation; from the scale of the giant lobes of radio galaxies stretched out across many Mpc scales in the IGM, to jets from stellar mass black holes, down to the stirring up of molecular gas on sub-pc to pc scales in regions of low mass star formation. Jets are observed to be ubiquitous during the process of star formation and associated with a very broad range of objects; from brown dwarfs to B and perhaps even O stars. It is now well established that stellar mass black holes are associated with jets and that the properties of such "microquasars" scale very well only with the mass of the black hole (e.g. see review, Mirabel, 2005), and scale naturally to the limit of massive black holes. In this regard, galactic jets are being
In most cases, there is clear evidence that accretion disks are an essential part of the engine. The fact that outflows and jets are observed around nearly all astrophysical objects during the early stages of their existence - a time where the central objects (young stars, massive planets, or black holes) undergo significant accretion from surrounding collapsing gas structures and/or associated disks - argues for a deep link between accretion and outflow. A large body of observations and theoretical models increasingly suggests that such jets in all of these systems may be powered by the same mechanism, namely, hydromagnetic winds driven off of magnetised accretion disks (see e.g. Livio, 1999; Königl & Pudritz, 2000; Pudritz et al., 2007, for reviews on this subject). The basic theory for hydromagnetic disk winds was worked out by Blandford & Payne (1982) in the context of models for AGN jets, but was quickly found to be very important for understanding jets in protostellar systems (PN83). For accreting systems, the gravitational binding energy that is released as gas accretes through a disk and onto a central object is the ultimate source of the energy for jets. Therefore, if this energy can be efficiently tapped and carried by the jet, then jet energies and other properties should simply scale with the depth of the gravitational potential well created by the central mass. The torques that can be exerted by outflows upon the underlying disks were shown to be much larger than can be produced by even strongly turbulent disks (Pudritz & Norman, 1983; Pelletier & Pudritz, 1992). Thus, jets can also carry off most of the angular momentum underlying disks, thereby assisting in the growth of central objects - be they stars, black holes, or planets. (It should be noted that large scale spiral waves in accretion disks can also efficiently transport angular momentum through disks). Finally, the feedback of these powerful jets upon their surroundings can be important. In star forming systems, outflows may help to drive turbulence within cluster forming regions in molecular clouds, which would help to sustain such a region against global collapse.

It must be noted that isolated magnetised, spinning bodies such as magnetised A stars can also drive drive outflows. The extraordinary Chanra X-ray observations of Crab pulsar show that it driving off a highly collimated, relativistic jet. The models that best describe the physics of all of these jets demonstrate that outflows can be driven and collimated from spinning bodies (be they disks, stars, or compact objects). In protostellar jets, the slow spin of stars has been attributed
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either to star-disk coupling (Königl, 1991) or X-winds (Shu et al., 2000), but could also be explained as an accretion powered stellar wind from the central star (e.g., Matt & Pudritz, 2005). For black hole systems, the highly relativistic component of jets may be associated with electrodynamic processes and currents that arise from spinning holes in a magnetised environment provided by the surrounding disk (e.g., Blandford & Znajek, 1977).

It is impossible to do justice to the enormous body of excellent work on the physics of jets in these different systems. In this chapter, we give specific examples of these basic themes from our own research in these fields in the context of some of the basic literature. In §2 we review aspects of unified models for non-relativistic, and then relativistic jets in these systems. We follow this in §3 with a discussion of disk winds as an excellent candidate for a unified model, analysed particularly for protostellar jets. We follow this in §4 with simulations of jets during gravitational collapse and disk formation. We examine the role of feedback of jets in the specific context of star formation in §5. We then move on to the basics of relativistic jets from compact objects as well as black holes (§6), and then simulations (§7) of MHD jets in the general relativistic limit (GRMHD). We conclude with a comparison of the physics of relativistic vs non-relativistic jets.

1.2 Jets in diverse systems

Protostellar objects: Stars span about 4 decades in mass. Until recently, the study of outflows was restricted to sources that were easily detectable in millimetre surveys of molecular clouds. These studies show that outflows, over a vast range of stellar luminosities, have scalable physical properties. In particular the observed ratio of the momentum transport rate (or thrust) carried by the CO molecular outflow to the thrust that can be provided by the bolometric luminosity of the central star (e.g., Cabrit & Bertout, 1992); scales as

\[ F_{\text{outflow}} / F_{\text{rad}} = 250(L_{\text{bol}} / 10^3 L_\odot)^{-0.3}, \]

This relation has been confirmed and extended by the analysis of data from over 390 outflows, ranging over six decades up to $10^6 L_\odot$ in stellar luminosity (Wu et al., 2004). This remarkable result shows that that jets from both low and high mass systems, in spite of the enormous differences in the radiation fields from the central stars, are probably driven by a single, non-radiative, mechanism.
1.2 Jets in diverse systems

Jets are observed to have a variety of structures and time-dependent behaviour – from internal shocks and moving knots to systems of bow shocks that suggest long-time episodic outbursts (e.g., review [Bally et al., 2007]). They show wiggles and often have cork-screw like structure, suggesting the presence either of jet precession, the operation of non-axisymmetric kink modes, or both. Numerical simulations have recently shown that in spite of the fact that jets are predicted to have a dominant toroidal magnetic field, they are nevertheless stable against such nonlinear modes (e.g., Ouyed et al., 2003; Nakamura & Meier, 2004; Kigure & Shibata, 2005). The nonlinear saturation of the unstable modes is achieved by the natural regulation of the jet velocity to values near the local Alfvén speed. A second point is that jets will have some poloidal field along their “backbone” and this too prevents a jet from falling apart.

Two types of theories that have been proposed to explain protostellar jets, both of which are hydromagnetic: disk winds (e.g., review [Pudritz et al., 2007]), or the X-wind (e.g., review [Shang et al., 2007]). The latter model posits the interaction between a spinning magnetised stellar magnetosphere and the inner edge of an accretion disk as the origin of jets, while the former envisages jets as arising from large parts of magnetised disks. One of the best ways of testing such theories is by measuring jet rotation - which is a measure of how much angular momentum that a jet can extract from its source. Significant observational progress has been made on this problem lately with the discovery of the rotation of jets from T-Tauri stars. The angular momentum that is observed to be carried by these rotating flows (e.g. DG Tau) is a considerable fraction of the excess disk angular momentum – from 60–100% (e.g., Bacciotti, 2004), which is consistent with the high extraction efficiency that is predicted by the theoretical models. Another important result is that there is too much angular momentum in the observed jet to be accounted for the spinning star or very innermost region of the disk as predicted by the X-wind model. The result is well explained by the disk wind model which predicts that the angular momentum derives from large reaches of the disk, typically from a region as large as 0.5 AU (Anderson et al., 2003).

Jovian planets: The striking similarity of the system of Galilean moons around Jupiter, with the sequence of planets in our solar system, has long suggested that these moons may have formed through a subdisk around Jupiter (Mohanv et al., 2007). Recent hydrodynamical simulations of circumstellar accretion disks containing and building up an orbiting protoplanetary core have numerically proven the existence
of a circum-planetary sub-disk in almost Keplerian rotation close to the planet \cite{Kley+2001}. The accretion rate of these sub-disks is about $\dot{M}_{\text{cp}} = 6 \times 10^{-5} M_{\text{Jup}} \text{yr}^{-1}$ and is confirmed by many independent simulations. With that, the circum-planetary disk temperature may reach values up to 2000 K indicating a sufficient degree of ionisation for matter-field coupling and would also allow for strong equipartition field strength \cite{Fendt2003}. It should be possible, therefore, for lower luminosity jets to be launched from the disks around Jovian mass planets.

The possibility of a planetary scale MHD outflow, similar to the larger scale YSO disk winds, is indeed quite likely because: (i) the numerically established existence of circum-planetary disks is a natural feature of the formation of massive planets; and (ii) the feasibility of a large-scale magnetic field in the protoplanetary environment \cite{QuillenTrilling1998,Fendt2003}. One may show, moreover, that the outflow velocity is of the order of the escape speed for the protoplanet, at about $60 \text{ km s}^{-1}$ \cite{Fendt2003,Machida+2006}.

**Black holes**: Relativistic jets have been observed or postulated in various astrophysical objects, including active galactic nuclei (AGNs) \cite[e.g.,][]{UrryPadovani1995,Ferrari1998}, microquasars in our galaxy \cite[e.g.,][]{MirabelRodriguez1998}, and gamma-ray bursts (GRBs) \cite[e.g.,][]{Piran2005}. Table 1.1 summarises the features of relativistic jets in different large scale sources and demonstrate the wide range in power and Lorentz factors achieved by these systems.

| Source          | $L$         | $\Gamma$ |
|-----------------|-------------|---------|
| AGNs/Quasars    | $10^{43}-10^{48} \text{ erg s}^{-1}$ | $\sim 10$ |
| $\mu$-quasars   | $10^{38}-10^{40} \text{ erg s}^{-1}$ | 1-10 |
| GRBs            | $10^{51}-10^{52} \text{ erg s}^{-1}$ | $10^2-10^4$ |

In the commonly accepted standard model of large scale/relativistic jets \cite{Begelman+1984}, flow velocities as large as 99% of the speed of light (in some cases even beyond) are required to explain the apparent superluminal motion observed in many of these sources (see Table 1.1). Later, considerations of stationary MHD flows have revealed that relativistic jets must be strongly magnetised \cite{Michel1969,Camenzind1986,Li+1992,FendtCamenzind1996}. In that case, the avail-
able magnetic energy can be transferred over a small amount of mass with
high kinetic energy.

Here too, the most promising mechanisms for producing the relativis-
tic jets involve magnetohydrodynamic centrifugal acceleration and/or
magnetic pressure driven acceleration from an accretion disk around
the compact objects (e.g., Blandford & Payne, 1982), or involve the ex-
traction of rotating energy from a rotating black hole (Penrose, 1969;
Blandford & Znajek, 1977). These models have been applied to explain
jets features in the galactic micro-quasars GRS 1915+105 (Mirabel & Rodriguez,
1994) and GRO J1655-40 (Tingay et al., 1995), or a rotating super mas-
sive black hole in an active galactic nucleus, which is fed by interstellar
gas and gas from tidally disrupted stars. In general these studies re-
quire solving special relativistic MHD (SRMHD) or general relativistic
MHD (GRMHD) equations and often require sophisticated numerical
codes. We describe the basic theory behind relativistic MHD jets and
summarise recent GRMHD simulations of jets emanating from the vicin-
ity of accreting black holes (see §6 and 7) after we examine the general
theory of outflows from magnetised spinning disks and objects..

1.3 Theory of disk winds

Given the highly nonlinear behaviour of the force balance equation for
jets (the so-called Grad-Shafranov equation), theoretical work has fo-
cused on tractable and idealised time-independent, and axisymmetric or
self-similar models (e.g., BP82) of various kinds. We briefly summarise
the theory, below (see details in Pudritz et al., 2007; Pudritz, 2004).

1.3.0.1 Conservation Laws and Jet Kinematics

Conservation laws play a significant role in understanding astrophysical
jets. This is because whatever the details, conservation laws strongly
constrain the flux of mass, angular momentum, and energy. What can-
not be constrained by these laws will depend on the general physics of
the disks such as on how matter is loaded onto field lines.

Jet dynamics can be described by the time-dependent, equations of
ideal MHD. The evolution of a magnetised, rotating system that is
threaded by a large-scale field \( \mathbf{B} \) involves (i) the continuity equation
for a conducting gas of density \( \rho \) moving at velocity \( \mathbf{v} \) (which includes
turbulence); (ii) the equation of motion for the gas which undergoes
pressure (\( p \)), gravitational (with potential \( \Phi \)), and Lorentz forces; (iii)
the induction equation for the evolution of the magnetic field in the
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moving gas where the current density is $j = (c/4\pi)\nabla \times B$; (iv) the energy equation, where $e$ is the internal energy per unit mass; and, (v) the absence of magnetic monopoles. These are written as:

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0 \quad (1.2)$$

$$\rho \left( \frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \nabla) \mathbf{v} \right) + \nabla p + \rho \nabla \Phi - \frac{j \times B}{c} = 0 \quad (1.3)$$

$$\frac{\partial \mathbf{B}}{\partial t} - \nabla \times (\mathbf{v} \times \mathbf{B}) = 0 \quad (1.4)$$

$$\rho \left( \frac{\partial e}{\partial t} + (\mathbf{v} \nabla) e \right) + p(\nabla \cdot \mathbf{v}) = 0 \quad (1.5)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (1.6)$$

Progress can be made by restricting the analysis to stationary, as well as 2D (axisymmetric) flows, from which the conservation laws follow. It is useful to decompose vector quantities into poloidal and toroidal components (e.g. magnetic field $\mathbf{B} = \mathbf{B}_p + B_\phi \hat{e}_\phi$). In axisymmetric conditions, the poloidal field $\mathbf{B}_p$ can be derived from a single scalar potential $a(r, z)$ whose individual values, $a = \text{const}$, define the surfaces of constant magnetic flux in the outflow and can be specified at the surface of the disk (e.g., Pelletier & Pudritz 1992, PP92).

Conservation of mass and magnetic flux along a field line can be combined into a single function $k$ that is called the "mass load" of the wind which is a constant along a magnetic field line;

$$\rho \mathbf{v}_p = k \mathbf{B}_p. \quad (1.7)$$

This function represents the mass load per unit time, per unit magnetic flux of the wind. For axisymmetric flows, its value is preserved on each ring of field lines emanating from the accretion disk. Its value on each field line is determined by physical conditions - including dissipative processes - near the disk surface. It may be more revealingly recast as

$$k(a) = \frac{\rho \mathbf{v}_p}{B_p} = \frac{d\dot{M}_w}{d\Phi}, \quad (1.8)$$

where $d\dot{M}_w$ is the mass flow rate through an annulus of cross-sectional area $dA$ through the wind and $d\Phi$ is the amount of poloidal magnetic flux threading through this same annulus. The mass load profile is a function of the footpoint radius $r_0$ of the wind on the disk.

The toroidal field in rotating flows derives from the induction equa-
1.3 Theory of disk winds

\[ B_\phi = \frac{B}{k} (v_\phi - \Omega_0 r), \]  

(1.9)

where \( \Omega_0 \) is the angular velocity of the disk at the mid-plane. This result shows that toroidal fields in the jet are formed by winding up the field from the source. Their strength also depends on the mass loading as well as the jet density. Denser winds should have stronger toroidal fields. We note however, that the density does itself depend on the value of \( k \). Equation (9) also suggests that at higher mass loads, one has lower toroidal field strengths. This can be reconciled however, since it can be shown from the conservation laws (see below) that the value of \( k \) is related to the density of the outflow at the Alfvén point on a field line; \( k = (\rho_A/4\pi)^{1/2} \) (e.g., PP92). Thus, higher mass loads correspond to denser winds and when this is substituted into equation (9), we see that this also implies stronger toroidal fields.

Conservation of angular momentum along each field line leads to the conserved angular momentum per unit mass;

\[ l(a) = rv_\phi - \frac{rB_\phi}{4\pi k} = \text{const.} \]  

(1.10)

The form for \( l(a) \) reveals that the total angular momentum is carried by both the rotating gas (first term) as well by the twisted field (second term), the relative proportion being determined by the mass load.

The value of \( l(a) \) that is transported along each field line is fixed by the position of the Alfvén point in the flow, where the poloidal flow speed reaches the Alfvén speed for the first time (\( m_A = 1 \)). It is easy to show that the value of the specific angular momentum is;

\[ l(a) = \Omega_0 r_A^2 = (r_A/r_0)^2 l_0. \]  

(1.11)

where \( l_0 = v_{K,0} r_0 = \Omega_0 r_0^2 \) is the specific angular momentum of a Keplerian disk. For a field line starting at a point \( r_0 \) on the rotor (disk in our case), the Alfvén radius is \( r_A(r_0) \) and constitutes a lever arm for the flow. The result shows that the angular momentum per unit mass that is being extracted from the disk by the outflow is a factor of \( (r_A/r_0)^2 \) greater than it is for gas in the disk. For typical lever arms, one particle in the outflow can carry the angular momentum of ten of its fellows left behind in the disk.

Conservation of energy along a field line is expressed as a generalised version of Bernoulli’s theorem (this may be derived by taking the dot product of the equation of motion with \( B_\phi \)). Thus, there is a specific
energy $e(a)$ that is a constant along field lines, which may be found in many papers (e.g., BP82 and PP92). Since the terminal speed $v_p = v_\infty$ of the disk wind is much greater than its rotational speed, and for cold flows, the pressure may also be ignored, one finds the result:

$$v_\infty \simeq 2^{1/2} \Omega_0 r_A = (r_A/r_0)v_{\text{esc},0}. \quad (1.12)$$

There are three important consequences for jet kinematics here: (i) that the terminal speed exceeds the local escape speed from its launch point on the disk by the lever arm ratio; (ii) the terminal speed scales with the Kepler speed as a function of radius, so that the flow will have an onion-like layering of velocities, the largest inside, and the smallest on the larger scales, as seen in the observations; and (iii) that the terminal speed depends on the depth of the local gravitational well at the footpoint of the flow – implying that it is essentially scalable to flows from disks around YSOs of any mass and therefore universal.

Another useful form of the conservation laws is the combination of energy and angular momentum conservation to produce a new constant along a field line (e.g., PP92): $j(a) \equiv e(a) - \Omega_0 l(a)$. This expression has been used (Anderson et al., 2003) to deduce the rotation rate of the launch region on the Kepler disk, where the observed jet rotation speed is $v_{\phi,\infty}$ at a radius $r_\infty$ and which is moving in the poloidal direction with a jet speed of $v_{p,\infty}$. Evaluating $j$ for a cold jet at infinity and noting that its value (calculated at the foot point) is $j(a_0) = -(3/2)v_{K,0}^2$, one solves for the Kepler rotation at the point on the disk where this flow was launched:

$$\Omega_0 \simeq v_{p,\infty}^2 / (2 v_{\phi,\infty} r_\infty). \quad (1.13)$$

When applied to the observed rotation of the Large Velocity Component (LVC) of the jet DG Tau (Bacciotti et al., 2002), this yields a range of disk radii for the observed rotating material in the range of disk radii, 0.3–4 AU, and the magnetic lever arm is $r_A/r_0 \simeq 1.8–2.6$.

### 1.3.0.2 Angular Momentum Extraction

How much angular momentum can such a wind extract from the disk? The angular momentum equation for the accretion disk undergoing an external magnetic torque may be written:

$$\dot{M}_a \frac{d(r_0 v_0)}{dr_0} = -r_0^2 B_\phi B_z |_{r_0,H}. \quad (1.14)$$
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where we have ignored transport by MRI turbulence or spiral waves. By using the relation between poloidal field and outflow on the one hand, as well as the link between the toroidal field and rotation of the disk on the other, the angular momentum equation for the disk yields one of the most profound scaling relations in disk wind theory – namely – the link between disk accretion and mass outflow rate (see KP00, PP92 for details):

\[ \dot{M}_a \simeq (r_a/r_0)^2 \dot{M}_w. \] (1.15)

The observationally well known result that in many systems, \( \dot{M}_w/\dot{M}_a \simeq 0.1 \) is a consequence of the fact that lever arms are often found in numerical and theoretical work to be \( r_a/r_0 \simeq 3 \) – the observations of DG Tau being a perfect example.

1.3.0.3 Jet Power and Universality

These results can be directly connected to the observations of momentum and energy transport in the molecular outflows. Consider the total mechanical power that is carried by the jet, which may be written as (e.g., Pudritz, 2003):

\[ L_{\text{jet}} = \frac{1}{2} \int_{r_i}^{r_f} d\dot{M}_w v_w^2 \simeq \frac{GM_* \dot{M}_a}{2r_i} \left( 1 - \frac{r_f^2}{r_i^2} \right) \simeq \frac{1}{2} L_{\text{acc}}. \] (1.16)

This explains the observations of Class 0 outflows wherein \( \dot{L}_w/\dot{L}_{\text{bol}} \simeq 1/2 \), since the main luminosity of the central source at this time is due to accretion and not nuclear reactions. (The factor of 1/2 arises from the dissipation of some accretion energy as heat at the inner boundary). The ratio of wind to stellar luminosity decreases at later stages because the accretion luminosity becomes relatively small compared to the bolometric luminosity of the star as it nears the ZAMS.

This result states that the wind luminosity taps the gravitational energy release through accretion in the gravitational potential of the central object – and is a direct consequence of Bernoulli’s theorem. This, and the previous results, imply that jets may be produced in any accreting system. The lowest mass outflow that has yet been observed corresponds to a proto-brown dwarf of luminosity \( \simeq 0.09L_\odot \), a stellar mass of only 20 – 45\( M_{\text{Jup}} \), and a very low mass disk \( < 10^{-4}M_\odot \) (Bourke et al., 2005).

On very general grounds, disk winds are also likely to be active during massive star formation (e.g., König, 1999). Such outflows may already start during the early collapse phase when the central YSO
still has only a fraction of a solar mass (e.g., Banerjee & Pudritz, 2006; Banerjee & Pudritz, 2007). Such early outflows may actually enhance the formation of massive stars via disk accretion by punching a hole in the infalling envelope and releasing the building radiation pressure (e.g., Krumholz et al., 2005).

1.3.0.4 Jet Collimation

In the standard picture of hydromagnetic winds, collimation of an outflow occurs because of the increasing toroidal magnetic field in the flow resulting from the inertia of the gas. Beyond the Alfvén surface, equation (8) shows that the ratio of the toroidal field to the poloidal field in the jet is of the order $B_\phi/B_p \simeq r/r_\Lambda \gg 1$, so that the field becomes highly toroidal. In this situation, collimation is achieved by the tension force associated with the toroidal field which leads to a radially inwards directed component of the Lorentz force (or “z-pinch”); $F_{\text{Lorentz},r} \simeq j_z B_\phi$. The stability of such systems is examined in the next section.

In Heyvaerts & Norman (1989) it was shown that two types of solution are possible depending upon the asymptotic behaviour of the total current intensity in the jet;

$$I = 2\pi \int_0^r j_z(r', z')dr' = (c/2) r B_\phi. \quad (1.17)$$

In the limit that $I \to 0$ as $r \to \infty$, the field lines are paraboloids which fill space. On the other hand, if the current is finite in this limit, then the flow is collimated to cylinders. The collimation of a jet therefore depends upon its current distribution – and hence on the radial distribution of its toroidal field – which, as we saw earlier, depends on the mass load. Mass loading therefore must play a very important role in controlling jet collimation.

It can be shown (Pudritz et al., 2006, PRO) from this that jets should show different degrees of collimation depending on how they are mass loaded. As an example, neither the highly centrally concentrated, magnetic field lines associated with the initial split-monopole magnetic configuration used in simulations by Romanova et al. (1997), nor the similar field structure invoked in the X-wind (see review by Shu et al., 2000) should become collimated in this picture. On the other hand, less centrally (radially) concentrated magnetic configurations such as the potential configuration of Ouyed et al. (1997) and BP82 should collimate to cylinders.
1.4 Gravitational collapse, disks, and the origin of outflows

Fig. 1.1. Large scale outflow (left panel, scale of hundreds of AU), and small scale disk wind and jet formed (right panel, scale of a fraction of an AU) during the gravitational collapse of a magnetised B-E, rotating cloud core. Cross-sections through the disk and outflows are shown – the blue contour marks the Alfvén surface. Snapshots taken of an Adaptive Mesh calculation at about 70,000 years into the collapse. [Adapted from Banerjee & Pudritz (2006)].

This result also explains the range of collimation that is observed for molecular outflows. Models for observed outflows fall into two general categories: the jet-driven bow shock picture, and a wind-driven shell picture in which the molecular gas is driven by an underlying wide-angle wind component such as given by the X-wind (see review by Cabrit et al. (1997). A survey of molecular outflows by Lee et al. (2000) found that both mechanisms are needed in order to explain the full set of systems observed.

1.4 Gravitational collapse, disks, and the origin of outflows

Jets are expected to be associated with gravitational collapse because disks are the result of the collapse of rotating molecular cloud cores. One of the first simulations to show how jets arise during gravitational collapse is the work of Tomisaka (1998, 2002). These authors used as initial conditions the collapse of a magnetised rotating filament (cylinder) of gas and showed that this gave rise to the formation of a disk from which a centrifugally driven disk wind was produced.

Banerjee & Pudritz revisited the problem of collapsing magnetised molecular cloud cores in Banerjee & Pudritz (2006) (low mass cores, BP06) and Banerjee & Pudritz (2007) (high mass cores, BP07). Here the initial cores are modelled as supercritical Bonnor-Ebert (B-E) spheres.
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Fig. 1.2. Magnetic field line structure, outflow and disk. The two 3D images show the magnetic field lines, isosurfaces of the outflow velocities and isosurfaces of the disk structure at the end of our simulation ($t \approx 7 \times 10^4$ years) at the two different scales as shown in Fig. 1.1. The isosurfaces of the upper panel refer to velocities $0.18 \text{ km s}^{-1}$ (light red) and $0.34 \text{ km s}^{-1}$ (red) and a density of $2 \times 10^{-16} \text{ g cm}^{-3}$ (gray) whereas the lower panel shows the isosurfaces with velocities $0.6 \text{ km s}^{-1}$ (light red) and $2 \text{ km s}^{-1}$ (red) and the density at $5.4 \times 10^{-9} \text{ g cm}^{-3}$ (gray). [Adapted from Banerjee & Pudritz (2006)].

with a slight spin. B-E spheres are well controlled initial setups for numerical experiments and there are plenty of observed cores which can fit by a B-E profile (see Lada et al., 2007, on a compilation of hydrostatic cores). Recent efforts have focused on understanding the evolution of magnetised B-E spheres (e.g., Matsumoto & Hanawa, 2003; Machida et al., 2005; Hennebelle & Teyssier, 2005; Hennebelle & Fromang, 2008). Whereas purely hydrodynamic collapses of such objects never show outflows (e.g., Foster & Chevalier, 1993; Banerjee et al., 2004), the addition of a magnetic field produces them.

Outflows are also implicated in playing a fundamental role in controlling the overall amount of mass that is incorporated into a new star from the original core. Observations of the so-called core mass function (CMF) that describes the spectrum of core masses, show that it has nearly the identical shape as the initial mass function (IMF) that describes the spectrum of initial stellar masses - with one caveat. The CMF needs to be multiplied by a factor of $1/3$ (in one well studied case - see Alves et al., 2007) in order to align them. The suggestion is that protostellar outflows are responsible for removing as much as $2/3$ of a core's mass in the process of collapse and star formation.

In BP06, the FLASH AMR code (Fryxell et al., 2000) was used to follow the gravitational collapse of magnetised molecular cloud cores.
1.4 Gravitational collapse, disks, and the origin of outflows

Fig. 1.3. Snapshots of the central region of a collapsing magnetised massive cloud core ($M_{\text{core}} \sim 170 \, M_\odot$). The left panel shows the situation at $t = 1.45 \times 10^4$ years ($1.08 \, t_{\text{ff}}$) into the collapse and before the flow reversal and right panel shows the configuration 188 years later when the outflow is clearly visible. [Adapted from Banerjee & Pudritz (2007)].

This code allowed the Jeans length to be resolved with at least 8 grid points throughout the collapse calculations. This ensures compliance with the Truelove criterion (Truelove et al., 1997) to prevent artificial fragmentation.

The results of the low mass core simulation are shown in fig. 1.1 and fig. 1.2; they show the end state (at about 70,000 yrs) of the collapse of a B-E sphere that is chosen to be precisely the Bok globule observed by Alves et al. (2001) – whose mass is $2.1 \, M_\odot$ and radius $R = 1.25 \times 10^4$ AU at an initial temperature of 16 K. Two types of outflow can be seen: (i) an outflow that originates at scale of $\approx 130$ AU on the forming disk that consists of a wound up column of toroidal magnetic field whose pressure gradient pushes out a slow outflow; and (ii) a disk-wind that collimates into a jet on scale of 0.07 AU. A tight proto-binary system has formed in this simulation, whose masses are still very small $\leq 10^{-2} \, M_\odot$, which is much less than the mass of the disk at this time $\approx 10^{-1} \, M_\odot$.

The outer flow bears the hallmark of a magnetic tower, first observed by Uchida & Shibata (1985), and studied analytically by Lynden-Bell (2003). Both flow components are unbound, with the disk wind reaching 3 km s$^{-1}$ at 0.4 AU which is quite super-Alfvénic and above the local escape speed. The outflow and jet speeds will increase as the central mass grows.

In Banerjee & Pudritz (2007) a similar setup to the one described above is used to study the formation of massive stars. In an unified pic-
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ture of star formation one would also assume that outflows and jets will also be launched around young massive stars. Due to the deeply embedded nature of massive star formation observations of jets and outflows are much more difficult than for the more isolated low mass stars. Nevertheless, there is increasing observational evidence for outflows around young massive stars (e.g., Arce et al. 2007; Zhang et al. 2007).

The simulations of collapsing magnetised massive cores (∼170 M⊙) show clear signs of outflows during the early stages of massive star formation. Fig. 1.3 shows two 2D snap-shots from this simulations where the onset of an outflow is visible. Early outflows of this kind might have a large impact on the accretion history of the young massive star. Krumholz et al. (2005) showed that cavities blown by outflows help to release the radiation pressure from the newly born massive star which in turn relaxes radiation pressure limited accretion onto the central object.

From these simulations one can conclude that the theory and computation of jets and outflows is in excellent agreement with new observations of many kinds. Disk winds are triggered during magnetised collapse and persist throughout the evolution of the disk. They efficiently tap accretion power, transport a significant part portion of the disk’s angular momentum, and can achieve different degrees of collimation depending on their mass loading.

1.5 Feedback from collimated protostellar jets?
The interstellar medium (ISM) and star forming molecular clouds are permeated by turbulent, supersonic gas motions (e.g. see recent reviews Elmegreen & Scalo 2004; Mac Low & Klessen 2004; Ballesteros-Paredes et al. 2007, and references therein). Supersonic turbulence in molecular clouds is known to play at least two important roles: it can provide pressure support to help support molecular clouds against rapid collapse, and it can also produce the system of shocks and compressions throughout such clouds which fragment the cloud into the dense cores that are the actual sites of gravitational collapse and star formation. One of the major debates in the literature is on the question of just how long this turbulence can be sustained - and with it, star formation. Despite the importance of supersonic turbulence for the process of star formation its origin is still unclear. Norman & Silk (1980) proposed that Herbig-Haro (HH) outflows could provide the energetics to drive the turbulence in molecular clouds that could keep star formation active for a number of cloud free-fall times (i.e., several million years). This is an attrac-
1.5 Feedback from collimated protostellar jets?

In a recent investigation by Banerjee et al. (2007) this idea was addressed within a detailed study of feedback from collimated jets on their supra-core environment. These studies show that supersonic turbulence excited by collimated jets decays very quickly and does not spread far from the driving source, i.e. the jet. Supersonic fluctuations, unlike subsonic ones, are highly compressive. Therefore, the energy deposited by a local source, like a collimated jet, stays localised as the compressed gas either heats up (in the non-radiative case) or is radiated by cooling processes. The re-expansion of the compressed regions excites only subsonic or marginally supersonic fluctuations. This is contrary to almost incompressible subsonic fluctuations. These fluctuations travel like linear waves with little damping into the ambient medium and make up most of the overall velocity excitations.

In the study of Banerjee et al. (2007) a series of numerical experiments were performed with the adaptive mesh refinement (AMR) code FLASH. Here the jet is modelled as a kinetic energy injection from the box boundary. The energy injection could be switched on and off after a certain amount of time, i.e. the jets are either continuously driven or transient. These jets are either interacting with a homogeneous or clumpy environment, where the latter is modelled as a spherical overdensity. Additionally the influence of magnetic fields on jet-excited fluctuations where considered in this study.

The authors quantify the jet-excited motions of the gas using velocity probability distribution functions (PDF) which can be regarded as volume weighted histograms. From these PDF the distinction between subsonic and supersonic fluctuations comes about naturally: the supersonic regime is strongly suppressed compared to the subsonic regime. Additionally, separate decay laws of the kinetic energy for the subsonic and supersonic regimes where derived to quantify their different behaviours with time.

Fig. 1.4 shows that the kinetic energy of the supersonic fluctuations decays much faster than subsonic contributions. Furthermore, one can see from fig. 1.5 that the jet-excited fluctuations are mainly subsonic which is due to the fact that the supersonic excitations do not travel far from the edge of the jet (see right panel of fig. 1.5).

This study shows that supersonic fluctuations are damped quickly because they excite mainly compressive modes. The re-expansion of the compressed overdensities drives mainly subsonic velocity fluctua-
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Fig. 1.4. Time evolution of the kinetic energies, $E_{\text{kin}}$, for a transient jet with a speed of Mach 5. The quantities are divided into a supersonic regime, $v > c$, and a subsonic regime, $v < c$. The decay of supersonic energy contributions is much faster (faster than $\propto t^{-2}$) than the subsonic one. The time shift between the peaks in the kinetic energies shows that the subsonic fluctuations are powered by the decay of the supersonic motions. [Adapted from Banerjee et al. (2007)]

Simulations of magnetised jets in this study show that jets stay naturally more collimated if the magnetic field is aligned with the jet axis and therefore entrain less gas. Furthermore, perpendicular motions are damped by magnetic tension preventing a large spread of high ampli-
Fig. 1.5. Velocity structure and velocity PDFs at different times of a continuously driven Mach 10 jet. Essentially no supersonic fluctuations in the ambient gas get excited by the jet (the peak at $v/c = 10$ is the jet itself). [Adapted from Banerjee et al. (2007)]

This study shows that collimated jets from young stellar objects are unlikely drivers of large-scale supersonic turbulence in molecular clouds. Alternatively it could be powered by large scale flows which might be responsible for the formation of the cloud itself (e.g., Ballesteros-Paredes et al., 1999). Energy cascading down from the driving scale to the dissipation scale will then produce turbulent density and velocity structure in the inertial range in between (Lesieur, 1997). If the large-scale dynamics of the interstellar medium is driven by gravity (as suggested, e.g., by Li et al., 2005; Li et al., 2006) gravitational contraction would also determine to a large extent the internal velocity structure of the cloud. Otherwise, blast waves and expanding shells from supernovae are also viable candidates to power supersonic turbulence in molecular clouds (see e.g., Mac Low & Klessen, 2004).

1.6 Relativistic jets: theory

The theory of relativistic MHD jets is best described in the SRMHD regime. The essential parameter is the magnetisation parameter (Michel, 1969),

$$\sigma = \frac{\Psi^2 \Omega_F^2}{4Mc^3}. \quad (1.18)$$
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The iso-rotation parameter \( \Omega_F \) is frequently interpreted as the angular velocity of the magnetic field lines. The function \( \Psi = B_p r^2 \) is a measure of the magnetic field distribution (see Li 1993), and \( \dot{M} = \pi \rho v_p R^2 \) is the mass flow rate within the flux surface. Equation (1.18) demonstrates that the launch of a highly relativistic (i.e. highly magnetised) jet essentially requires at least one of three conditions – rapid rotation, strong magnetic field and/or a comparatively low mass load.

In the case of a spherical outflow (\( \Psi = \text{const} \)) with negligible gas pressure one may derive the Michel scaling between the asymptotic Lorentz factor and the flow magnetisation (Michel, 1969),

\[
\Gamma_\infty = \sigma^{1/3}
\]  

Depending on the exact magnetic field distribution \( \Psi(r, z) \) (which describes the opening of the magnetic flux surfaces), in a collimating jet the matter can be substantially accelerated beyond the fast point magnetosonic point (Begelman & Li, 1994; Fendt & Camenzind, 1996), as it is moved from infinity to a finite radius of several Alfvén radii. As a result, the power law index in Eq. (1.19) can be different from the Michel-scaling (see Fendt & Camenzind, 1996; Vlahakis & Königl, 2001, 2003).

The light cylinder (hereafter l.c.) is located at the cylindrical radius \( r_l = c / \Omega_F \). At the l.c. the velocity of the magnetic field lines “rotating” with angular velocity \( \Omega_F \) coincides with the speed of light \( c \). The l.c. has to be interpreted as the Alfvén surface in the limit of vanishing matter density (force-free limit). The location of the l.c. determines the relativistic character of the magnetosphere. If the light cylinder is comparable to the dimensions of the object investigated, a relativistic treatment of MHD is required.

While SRMHD can be invoked to study jets at scales larger than the gravitational radius of the black hole, close to the horizon general relativity becomes relevant and the MHD equations need to be coupled to General Relativity. This often requires the use of sophisticated numerical codes in order to capture the complexity of jet physics in GRMHD. In the (non-rotating) black hole’s Schwarzschild spacetime the GRMHD equations are identical to the SRMHD equations in general coordinates, except for the gravitational force terms and the geometric factors of the lapse function.

\[\dagger\] Outside the l.c. the magnetic field lines “rotate” faster than the speed of light. As the field line is not a physical object, the laws of physics are not violated.
Fig. 1.6. Lorentz factor from GRMHD simulations; adopted from De Villiers et al. (2005). In all panels, the axes are in units of $M$ (here $1 \, M \approx 4 \, \text{km}$). The black hole is located at the origin. The top row, left to right, shows plots of Lorentz factor for decreasing black hole rotation $a/M = 0.995, 0.9, 0.0$, respectively. The dotted contour marks the boundary of the jets. The only region where elevated values of $\Gamma$ are found is in the jets (and also in the bound plunging inflow near the black hole, which is not resolved at the scale of this figure). Maximum values of $\Gamma$ are found in knots that appear episodically in the upper and lower parts of the funnel. The bottom three panels show, from left to right, the corresponding time-averaged value of the Lorentz factor. The plots show evidence of spin-dependent collimation: cylindrical collimation is seen in the high-spin models, while the zero-spin model shows no such collimation. These plots also show evidence of an extended acceleration zone: large Lorentz factors are built up over the full radial range.

1.7 SRMHD and GRMHD Simulations

Full GRMHD numerical simulations on the formation of jets near a black hole were first performed by Kudoh et al. (1998). This was followed by a plethora of GRMHD codes with fixed spacetimes used to investigate jets from accreting black holes (e.g., De Villiers & Hawley, 2003; Gammie et al., 2003; Komissarov, 2004; Antón et al., 2006; Anninos et al., 2007). Here we focus on the most recent simulations that have managed to reach the highest Lorentz factors for the study of GRMHD jets.
Recent state-of-the-art GRMHD simulations have shown that the accretion flow launches energetic jets in the axial funnel region (i.e. low density region) of the disk/jet system, as well as a substantial coronal wind (De Villiers et al., 2005). The jets feature knot-like structures of extremely hot, ultra-relativistic gas; the gas in these knots begins at moderate velocities near the inner engine, and is accelerated to ultra-relativistic velocities (Lorentz factors of 50, and higher) by the Lorentz force in the axial funnel. The increase in jet velocity takes place in an acceleration zone extending to at least a few hundred gravitational radii from the inner engine. The overall energetics of the jets are strongly spin-dependent, with high-spin black holes producing the highest energy and mass fluxes. In addition, with high-spin black holes, the ultra-relativistic outflow is cylindrically collimated within a few hundred gravitational radii of the black hole, whereas in the zero-spin case the jet retains a constant opening angle of approximately 20 degrees.

Figure 1.7 shows the Lorentz factor, $\Gamma$, in the funnel outflow from the GRMHD simulations discussed in De Villiers et al. (2005). Elevated values of the Lorentz factor are found in compact, hot, evacuated knots that ascend the funnel radially; the combination of low density and high temperature in the knots. The highest values of Lorentz factor reach

† Ouyed et al. (1997) discuss the appearance of knots in MHD jet simulations; the
the maximum allowed by the code ($\sim 50$, and test runs show that much higher values could be reached with a higher ceiling), and these are only found at large radii, suggesting the presence of an extensive region where the knots are gradually accelerated to higher Lorentz factors.

While the details of the mechanism by which GRMHD jets are launched are yet to be fully understood, the following summary addresses the salient points of the complex dynamics by which unbound the plasma from the vicinity of the compact star (see discussion in De Villiers et al., 2005; McKinney & Narayan, 2007): In the region of the accretion flow near the marginally stable orbit, both pressure gradients and the Lorentz force act to lift material away from the equatorial plane. Some of this material is launched magneto-centrifugally in a manner reminiscent of the scenario of Blandford & Payne (1982), generating the coronal wind; some of this material, which has too much angular momentum to penetrate the centrifugal barrier, also becomes part of the massive funnel-wall jet. There is also evidence that the low-angular momentum funnel outflow originates deeper in the accretion flow. Some of this material is produced in a gravitohydromagnetic interaction in the ergosphere (Punsly & Coroniti, 1990), and possibly in a process similar to that proposed by Blandford & Znajek (1977) where conditions in the ergosphere approach the force-free limit. The material in the funnel outflow is accelerated by a relatively strong, predominantly radial Lorentz force; gas pressure gradients in the funnel do not contribute significantly.

In the context of GRBs, the Lorentz factor $\Gamma$ of the relativistic wind must reach high values ($\Gamma \sim 10^2 - 10^4$) both to produce $\gamma$-rays and to avoid photon-photon annihilation along the line of sight, whose signature is not observed in the spectra of GRBs (Goodman, 1986). As we have seen, GRMHD simulations indicate Lorentz factors of up to $\sim 50$ close to the black hole with indication of higher $\Gamma$ far beyond the central object where SRMHD treatment is sufficient. The extreme Lorentz factors have been confirmed by recent SRMHD simulations which have found solutions with $\gamma$ as high as $\simeq 5000$ as shown in Figure 1.7 (see Fendt & Ouyed, 2004, for more discussion).

knots seen in those simulations consisted of high density material, in contrast to what is observed in the present simulations.
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1.8 Non-relativistic versus relativistic MHD jets

While SRMHD and GRMHD jet simulations show similarities with non-relativistic regime, there are nevertheless some important differences in the underlying physics.

- The existence of the l.c. as a natural length scale in relativistic MHD is not consistent with the assumption of a self-similar jet structure (as is often assumed in non-relativistic MHD jet models). The latter holds even more when general relativistic effects are considered.
- Contrary to Newtonian MHD, in the relativistic case electric fields cannot be neglected. The poloidal electric field component is directed perpendicular to the magnetic flux surface. Its strength scales with the l.c. radius, \( E_p = E_\perp = (r/r_l)B_p \). As a consequence of \( E_p \approx B_p \), the effective magnetic pressure can be lowered by a substantial amount (Begelman & Li, 1994).
- In relativistic MHD the poloidal Alfvén speed \( u_A \) becomes complex for \( r > r_l \), \( u_A^2 = B_p^2 (1 - (r/r_l)^2) = b_p^2 - E_\perp^2 \). Therefore, Alfvén waves cannot propagate beyond the l.c. and only fast magnetosonic waves are able to exchange information across the jet.
- In the relativistic case, break-up of the MHD approximation is an issue. The problem is hidden in the fact that one may find arbitrarily high velocities for an arbitrarily high flow magnetisation. However, an arbitrarily high magnetisation may be in conflict with the intrinsic MHD condition which requires a sufficient density of charged particles in order to be able to drive the electric current system (Michel, 1969). Below the Goldreich-Julian particle density \( n_G \) (Goldreich & Julian, 1969) the concept of MHD as applied to relativistic jet breaks down.

The points above imply that with due regard to the break down of the MHD approximation, it is feasible to scale the physics of jets from non-relativistic to ultra-relativistic MHD regime. GRMHD simulations suggest that a continuous scaling is more likely for slowly rotating black holes (where the radius of marginal stability is rather large and in a comparatively low-gravity region of spacetime) but show that a full general relativistic electrodynamic treatment is required to robustly treat the case of rapidly rotating black holes. GRMHD simulations also show that the situation is even more complex in the “plunge region”. This is the region of the disk within the radius of marginal stability in which the accretion flow is undergoing rapid inwards acceleration (ultimately crossing the event horizon at the velocity of light as seen by a locally
1.8 Non-relativistic versus relativistic MHD jets

non-rotating observer). Unless the magnetic field is extremely strong, this is a region where inertial forces will dominate and the commonly employed force-free approximation breaks down. As a particle gets closer to the central object it crosses three regions defining the MHD, force-free and inertial regimes, respectively. Another crucial aspect of the physics inherent to ultra-relativistic jets emanating from the near vicinity of rapidly rotating central objects is the role played by the dynamically important electric field (e.g., Lyutikov & Ouyed, 2007; Ouyed et al., 2007) in driving plasma instabilities generally leading to pair creation. Pair-creation and subsequent annihilation into radiation has yet to be taken into account, at least consistently, in GRMHD or/and GREMD codes which are mostly based on mass conservation schemes.

In summary, it appears that a universal aspect of all jets is their ability to tap their energy from the accretion of gas into the gravitational potential of the underlying central object. Jets of all stripes may also share common morphological features far beyond the source. However, complete unification of non-relativistic with relativistic regimes for jets probably breaks down close to the central black holes, where jet physics will also depend on the rotation of the hole and the production of relativistic plasmas.

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