MAGNETIC FIELD EFFECTS ON THE STRUCTURE AND EVOLUTION OF OVERDENSE RADIATIVELY COOLING JETS

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ABSTRACT

We investigate the effect of magnetic fields on the propagation dynamics and the morphology of overdense, radiatively cooling, supermagnetosonic jets, with the help of fully three-dimensional smoothed particle magnetohydrodynamic simulations. Evaluated for a set of parameters that are mainly suitable for protostellar jets (with density ratios between those of the jet and the ambient medium \( \eta \approx 3-10 \), and ambient Mach number \( M_a \approx 24 \)), these simulations are also compared with baseline nonmagnetic and adiabatic calculations. Two initial magnetic field topologies (in approximate equipartition with the gas, \( \beta = p_j/p_B \approx 1 \)) are considered: (1) a helical field and (2) a longitudinal field, both of which permeate both the jet and the ambient medium. We find that, after amplification by compression and reorientation in nonparallel shocks at the working surface, the magnetic field that is carried backward with the shocked gas into the cocoon improves the jet collimation relative to the purely hydrodynamic (HD) systems, but this effect is larger in the presence of the helical field. In both magnetic configurations, low-amplitude, approximately equally spaced \( (4 \approx 2 - 4R_\text{p}) \) internal shocks (which are absent in the HD systems) are produced by magnetohydrodynamic (MHD) Kelvin-Helmholtz reflection pinch modes. The longitudinal field geometry also excites nonaxisymmetric helical modes that cause some beam wiggling. The strength and amount of these modes are, however, reduced (by about 2 times) in the presence of radiative cooling relative to the adiabatic cases. Besides, a large density ratio, \( \eta \), between the jet and the ambient medium also reduces, in general, the number of the internal shocks. As a consequence, the weakens of the induced internal shocks makes it doubtful that the magnetic pinches could by themselves produce the bright knots observed in the overdense, radiatively cooling protostellar jets. Magnetic fields may leave also important signatures on the head morphologies of the radiative cooling jets. The amplification of the nonparallel components of the magnetic fields, particularly in the helical field geometry, reduces the postshock compressibility and increases the postshock cooling length. This may lead to stabilization of the cold shell of shocked material that develops at the head against both the Rayleigh-Taylor and global thermal instabilities. As a consequence, the clumps that develop by fragmentation of the shell in the HD jets tend to be depleted in the helical field geometry. The jet immersed in the longitudinal field, on the other hand, still retains the clumps, although they have their densities decreased relative to the HD counterparts. As stressed in our previous work, since the fragmented shell structure resembles the knotty pattern commonly observed in HH objects behind the bow shocks of protostellar jets, this result suggests that, as long as (equipartition) magnetic fields are present, they should probably be predominantly longitudinal at the heads of these jets.

Subject headings: ISM: jets and outflows — MHD — stars: pre-main-sequence — stars: mass loss

1. INTRODUCTION

There is increasing evidence that protostellar jets are driven by circumstellar magnetized disks associated with pre-main-sequence stars (see, e.g., Königl & Ruden 1993; Shu et al. 1995). Efficient mass loss in supersonic, collimated magnetized outflows is the most likely mechanism by which protostars dissipate the angular momentum accumulated during the accretion of the surrounding material. While of fundamental importance in the production and initial collimation of the jets, magnetic fields have generally been neglected in most of the analytical and numerical modeling of the structures of protostellar jets since the inferred estimates of their strength \( (B \approx 10^{-6} \text{ to } 10^{-5} \text{ G}) \) suggested that they could be not dynamically important along the flows (see, e.g., Morse et al. 1993). However, recent observations of the circularly polarized radio emission of the young stellar outflow of T Tauri S (Ray et al. 1997), for example, indicate the presence of a strong, ordered magnetic field in the flow, far away from the source, which has possibly been amplified by compression behind the shocks at the head of the outflow. After amplification and reorientation behind the shocks, such magnetic fields may operate significant changes in the dynamics and collimation of a jet and also in its head structure, as suggested by recent numerical studies (see below).

Great effort has been concentrated in the analytical and numerical study of magnetized, adiabatic, light jets, mostly in the investigation of extragalactic jets (for a review, see, e.g., Birkinshaw 1997). Most of that work has focused on the study of the stability properties of the beam against hydromagnetic and Kelvin-Helmholtz instabilities. In the limit of zero-velocity difference between the jet and the surrounding medium, linear theory predicts that a jet magnetically confined by a toroidal field is unstable to the pinch and kink (or helical) hydromagnetic modes (see, e.g., Chandrasekhar 1961), with the temporal growth rates of the pinching mode increasing with increasing density ratio, \( \eta \), between the jet and the ambient medium. In the presence of a nonzero velocity discontinuity at the boundary layer separating the two fluids, these pure hydromagnetic modes
are modified by the development of the Kelvin-Helmholtz (K-H) instability. A jet confined by a toroidal magnetic field is unstable to the fundamental and reflection pinch and kink modes of the K-H instability (Cohn 1983; Fiedler & Jones 1984). The most unstable pinching mode in this case, at wavelengths $\lambda \sim 2\pi R_j$ (where $R_j$ is the jet radius), has a destabilization length $l \propto (M_j R_j)$ (where $M_j = v_j/c_j$ is the jet Mach number), which is similar to that in the pure hydrodynamical (HD) case for jets with (Cohn jet Mach number), which is similar to that in the pure K-H modes become stable for flows (except for sub-Alfvénic magnetic field residing in an unmagnetized medium, all K-H modes become stable for sub-Alfvénic flows (except for a small region of slow reflection modes; see, e.g., Bodo et al. 1989; Hardee et al. 1992; Hardee, Clarke, & Rosen 1997). In the super-Alfvénic regime, on the other hand, the growth rates of the instability are not very much different from those of pure HD flows. For example, a fundamental kink mode in the pure HD case can be identified with an Alfvén disturbance of long wavelength in the magnetohydrodynamic (MHD) case, and reflection modes at shorter wavelengths can be identified with fast magnetosonic (ms) waves reflecting off the jet boundaries. In this super-Alfvénic regime, if the jet is also supermagnetosonic ($M_{\text{ms}, j} = v_j/(v_A^2 + c_s^2)^{1/2} > 1$, where $v_A$ is the Alfvén velocity and $c_s$ is the sound speed), then it becomes more stable with increasing $M_{\text{ms}, j}$, with the destabilization length varying approximately proportional to $M_{\text{ms}, j} (l \propto M_{\text{ms}, j} R_j)$ (Hardee et al. 1992). Also, strong toroidal fields of strength comparable to the poloidal field can lead to increased jet stability (Appl & Camenzind 1992; Thiele & Camenzind 1997).

The effects of the K-H modes on the survival of the beam, however, cannot be predicted by the linear theory alone. As in pure HD flows, the fastest growing, shortest wavelength reflection modes are expected to saturate with the formation of weak oblique shocks, while the fundamental modes may not saturate and may cause large-scale distortions and even disruption of the flow. Numerical simulations that can determine the end points of the operation of these instabilities confirm these predictions. Hardee et al. (1997) and Hardee et al. (1997), for example, assuming slab and cylindrical jets, respectively, with axial magnetic fields, have focused on the comparison of the scale length of the structures generated during nonlinear evolution and the wavelengths of the maximum growth rate predicted by the linear theory. They find that the jet is not stabilized by nonlinear effects associated with increasing magnetic tension and disrupts near the resonant wavelength of the K-H kink mode. Malagoli, Bodo, & Rosner (1996) and Min (1997a, 1997b) have extended those investigations of the nonlinear development of the K-H instability by including diffusion and magnetic reconnection effects and found that even if the magnetic field intensity is not too large to completely suppress the K-H instability, it is still able to mediate turbulence decay and diffusion of energy and mass across the boundary layer between jet and surrounding medium.

These studies of the stability properties of the beam against MHD K-H pinch and helical modes seem to provide potential mechanisms to explain the formation of structures such as knots and wiggles in adiabatic, supermagnetosonic light jets. Besides these stability analyses, some effort has also been spent in numerical studies of the general effects of $B$-fields on the global evolution and morphology of adiabatic, light jets, still in the context of extragalactic jets, assuming both passive (Clarke, Norman, & Burns 1989; Matthews & Scheuer 1990; Hardee & Norman 1990) and dynamically important toroidal fields (see, e.g., Clarke, Norman, & Burns 1986; Lind et al. 1989; Kössl, Müller, & Hillebrandt 1990). In the latter case, the jet was found to be rapidly decelerated at the Mach disk with the shocked jet material being confined to a slender trans-Alfvénic plug instead of being deposited in the large cocoon observed in the purely HD light jets.

Lately, these numerical MHD studies have been extended to heavy, adiabatic jets (Todo et al. 1993; Hardee & Clarke 1995; Stone, Xu, & Hardee 1997). Hardee & Clarke and Stone et al. have focused on simulations of jets with poloidal fields propagating into an unmagnetized medium, while Todo et al. (see also Thiele & Camenzind 1997) have considered a helical field configuration extending also to the

**Fig. 1.** Longitudinal ($B_l$; left panel) and toroidal ($B_b$; right panel) magnetic field components (see eqs. [3a], [3b], and [3c] in the text) as a function of the radial distance $r = (y^2 + z^2)^{1/2}$, for $\beta = 1$ at the jet axis. The coordinates are in code units (for the magnetic fields 1 c.u. = 21.5 $\mu G$).
Fig. 2.—Midplane magnetic field distribution evolution of the $\eta = 3$ radiatively cooling jets ML3r (right column) and MH3r (left column). The initial conditions are $\eta = n_j/n_a = 3$, $n_a = 200$ cm$^{-3}$, $M_j = 24$, $v_j = 398$ km s$^{-1}$, $\beta = 8\pi \rho / B^2 = 1$, $q_{J} = 8$, and $q_{ch} = 0.3$. The times and the jet head positions are: $t/t_d = 1.40$ and $x \approx 25R_j$ (top row); $t/t_d = 1.60$ and $x \approx 29R_j$ (middle row); $t/t_d = 1.65$ and $x \approx 30R_j$ (bottom row). Note the reorientation (and amplification), in both models, of the magnetic fields that are carried with shocked jet material into the cocoon.

ambient medium. Comparing the development of supermagnetosonic, adiabatic, heavy jets with axial and toroidal magnetic fields propagating into unmagnetized ambient media, Hardee & Stone (1997) have found that the toroidal geometry suppresses the mixing and entrainment of ambient gas that is found to develop in the axial case as a consequence of the K-H instability. Also, they have found that the kink mode has a longer wavelength and smaller amplitude in the toroidal configuration.

Still almost unexplored, however, is the role played by $B$-fields on the propagation dynamics and morphology of radiatively cooling, heavy jets (a scenario that is believed to occur in protostellar jets). In the limit of zero magnetic field, numerical simulations of radiatively cooling, heavy jets (see, e.g., Blondin, Fryxell & Königl 1990, hereafter BFK; de Gouveia Dal Pino & Benz 1993, 1994, hereafter GB93, GB94; Stone & Norman 1993a, 1993b, 1994; Chernin et al. 1994; de Gouveia Dal Pino, Birkinshaw, & Benz 1996, hereafter GBB96; de Gouveia Dal Pino & Birkinshaw 1996, hereafter GB96) have shown that thermal energy losses by the jet system have important effects on its dynamics. These studies have revealed that a cooling jet develops a dense, cold shell of shocked material at the head that is fragmented into clumps by Rayleigh-Taylor instability. Moreover, BFK and GB93 have found that the development of the K-H modes along the beam are inhibited by the presence of cooling. Recently, Hardee & Stone (1997) and Stone et al. (1997) (see also Massaglia et al. 1992) have examined the dynamics of K-H unstable cooling jets. Their linear analysis indicate that the growth of the K-H modes is very sensitive to the assumed form of the cooling function. In particular, if the cooling curve is a steep function of the temperature in the neighborhood of the equilibrium state, then the growth of K-H modes is reduced relative to the adiabatic jet—a result that is consistent with previous numerical simulations (BFK; GB93). With the inclusion of a longitudinal magnetic
field in a supermagnetosonic jet, Hardee & Stone (1997) find that the magnetic field does not strongly affect the differences in the K-H stability properties between an adiabatic and a cooling HD jet, provided that the magnetic pressure does not dominate. Besides, they find that the increase in the magnetic field strength makes the linear stability properties become more like those of an adiabatic jet. The nonlinear analysis of the growth of the K-H modes in cooling jets in the $B = 0$ limit shows similar behavior as in the pure adiabatic case. The jet can be disrupted near the resonant frequency of the fundamental kink K-H mode, and nonlinear, higher frequency, reflection waves tend to produce low-amplitude wiggles, which can result in strong shocks in the jet beam.

In the present work, we attempt to extend these previous investigations by exploring the nonlinear effects of magnetic fields (close to equipartition with the gas) on the global evolution and morphology of radiatively cooling, heavy jets and to test the morphological signatures of two different magnetic field geometries (a longitudinal and a helical configuration) on the dynamics of protostellar jets. A preliminary step in this direction was made in previous work (de Gouveia Dal Pino & Cerqueira 1996; Cerqueira, de Gouveia Dal Pino, & Herant 1997, hereafter Paper I), in which we have mainly focused on the effects of the magnetic fields in the jet head structure. With the help of three-dimensional (3-D) smoothed particle magnetohydrodynamical (SPMHD) simulations, we have found that the presence of a helical magnetic field (in close equipartition with the gas) may suppress the formation of the clumpy structure that is found to develop at the head of HD jets by fragmentation of the cold shell of shocked jet material. A cooling jet immersed in a longitudinal magnetic field, on the other hand, tends to retain the clumpy morphology at its head. In the present work, we perform 3-D smoothed-particle magnetohydrodynamical simulations to address the details of the magnetic field effects over the whole structure of the radiatively cooling, heavy jets, covering a more extensive range of parameters, and compare these with both nonmagnetic and adiabatic systems.

Contemporaneously with this work, Frank et al. (1998; see also Frank et al. 1997) have performed grid-based 2.5 dimensional simulations of magnetized, radiatively cooling jets assuming a toroidal ($B_{\phi}$) magnetic field and concluded that toroidal fields are able to excite the development of strong pinches along the beam. Although in our simulations we have assumed somewhat different initial conditions (see below), the results are qualitatively similar where both investigations overlap, except for the fact that in our simulations the MHD pinch modes are found not to be strong in radiatively cooling, heavy jets. This difference, however, is mainly due to differences in the assumed initial density ratios between the jet and the ambient medium, which are much larger in our analyses (see Paper I and discussion below).

In § 2, the numerical method and the initial conditions used are briefly described. In § 3, we present the results of our simulations, and in § 4, we present our conclusions and address the possible implications of our results.

2. NUMERICAL TECHNIQUE

In our simulations, we consider the MHD conservation equations in the ideal approximation:

$$\frac{dp}{dt} = -\rho \nabla \cdot v,$$

(1a)

$$\frac{dv}{dt} = - \nabla p + \frac{1}{4\pi \rho} (\nabla \times B) \times B,$$

(1b)

$$\frac{du}{dt} = -\frac{p}{\rho} (\nabla \cdot v) - \mathcal{L},$$

(1c)

and

$$\frac{dB}{dt} = -B (\nabla \cdot v) + (B \cdot \nabla) v,$$

(1d)

where the symbols have their usual meaning (i.e., $\rho$ is the density; $B$ is the magnetic field; $v$ is the specific internal energy, $\mathcal{L}$ is the radiative cooling rate, etc.). To close the system (eqs. [1a]–[1d]), an ideal equation of state is assumed:

$$p = (\gamma - 1) \rho u,$$

(1e)

with $\gamma = 5/3$.

The MHD equations above are solved (in Cartesian coordinates) using a modified version of the fully three-dimensional smoothed particle hydrodynamics (SPH) code originally developed by de Gouveia Dal Pino & Benz (1993; see also GB94; Chernin et al. 1994; GB96; GBB96) for the

![Fig. 3.—Density of the shell at the jet axis as a function of the time for the three radiatively cooling jets presented in Fig. 2: HD3r (solid line), ML3r (dotted line), and MH3r (dashed line).](image-url)
are taken at and the density and pressure scales can be calibrated using the markers in the top region of each plot.

panel at different positions along the flow in the head region of the ML3r (model (both sides of the beam are the blobs that are much denser in the ML3r central peak corresponds to the beam region, while the secondary peaks on both sides of the beam are the blobs that are much denser in the ML3r model (lower panel). These profiles are taken at $t/t_\lambda = 1.65$, and the density and pressure scales can be calibrated using the markers in the top region of each plot.

Fig. 4.—Density (solid lines) and pressure (dashed lines) across the flow at different positions along the flow in the head region of the ML3r (upper panel) and MH3r (lower panel) jets. The positions are in units of $R_\lambda$. The central peak corresponds to the beam region, while the secondary peaks on both sides of the beam are the blobs that are much denser in the ML3r model (upper panel) than in the MH3r model (lower panel). These profiles are taken at $t/t_\lambda = 1.65$, and the density and pressure scales can be calibrated using the markers in the top region of each plot.

Investigation of the evolution of purely hydrodynamic (HD) jets.

In the SPH formalism, equations (1a)–(1e) are described by (see, e.g., Benz 1990; Stellingwerf & Peterkin 1990; Monaghan 1992; Meglicki 1995):

$$\frac{d\rho_j}{dt} = - \sum_{j=1}^{N} m_j \left( \frac{p_i}{\rho_i^2} + \frac{p_j}{\rho_j^2} + \Pi_{ij} \right) \nabla_i W_{ij}$$

$$+ \frac{1}{4\pi\rho_i^2} \sum_{j=1}^{N} [m_j (B_i - B_j) \times \nabla_i W_{ij}] \times B_i, \quad (2a)$$

$$\frac{du_i}{dt} = \frac{p_i}{\rho_i^2} \sum_{j=1}^{N} m_j (v_i - v_j) \nabla_i W_{ij}$$

$$+ \frac{1}{2} \sum_{j=1}^{N} m_j \Pi_{ij} (v_i - v_j) \cdot \nabla_i W_{ij}, \quad (2b)$$

and

$$\frac{dB_k}{dt} = \frac{1}{\rho_i} \sum_{j=1}^{N} m_j (B_{i,k} v_i - v_{i,k} B_j) \nabla_i W_{ij} \quad (2c)$$

where $m_j$ is the mass of the particle $j$ (located at the position $r = r_j$); $\rho_j$ and $\rho_i$ are the density of the particles $i$ and $j$, respectively; $p_j$ and $p_i$ are their pressure; $B_i$ is the magnetic field at $r_i$ ($B_{i,k}$ is the $k$ component of $B_i$); $v_{ij} \equiv v_i - v_j$ is the velocity difference, and $v_{i,k}$ is $k$ component of $v_i$. The angle brackets on the left-hand sides of the equations mean that the physical quantities inside them are evaluated at the position $r_i$ (at the particle $i$). The term $\nabla_i W_{ij}$ is the gradient of the Kernel function at the position of the particle $i$, and $\Pi_{ij}$ is the artificial viscosity that allows for appropriate treatment of shock wave dissipation. We here adopt the von Neumann-Richtmyer viscosity (see, e.g., Benz 1990).

The density and pressure are both calculated from the definition of discreteness in SPH (see, e.g., Benz 1990; Monaghan 1992), and are expressed, respectively, by the equations

$$\langle \rho \rangle = \sum_{j=1}^{N} m_j W_{ij} \quad (2d)$$

and

$$\langle p \rangle = (\gamma - 1) \sum_{j=1}^{N} m_j u_j W_{ij}. \quad (2e)$$

As in previous work (GB93; GB94; GB96), the radiative cooling rate $\mathcal{L}$ (due to collisional excitation and recombination) in equation (1c) is implicitly calculated using the cooling function evaluated by Katz (1989) for a gas of cosmic abundances cooling from $T \approx 10^6 \text{ K}$ to $T \approx 10^4 \text{ K}$. (The cooling is suppressed below $T \approx 10^4 \text{ K}$, where the transfer of ionizing radiation becomes important and the assumption of a fully ionized flow breaks down; see, e.g., GB93; GB96).

The choice of a vectorial formalism to write the system of equations above (instead of a tensorial form) is to ensure that the magnetic force vector $[(\nabla \times B) \times B]$, in each point of the system, is strictly perpendicular to the magnetic field vector itself. Tensorial implementations in SPMHD, like those proposed by Monaghan (1992), are subject to the development of nonphysical magnetic force components, the projection of which onto the magnetic field vector has nonzero values. Furthermore, these parallel components of the magnetic force are proportional to $\nabla \cdot \mathbf{B}$ (see Brackbill...
Fig. 5.—Left column: Density (solid lines) and pressure (dashed lines) profiles across the HD3r (top row), ML3r (middle row), and MH3r (bottom row) jets taken in three different pinching positions along the flow (at $t/t_{\text{d}} = 1.65$). All the profiles have been scaled as in Fig. 4. Right column: The corresponding axial density along the jet axis ($y = z = 0$) showing the channel of internal shocks. The position of the three pinches depicted in the left column are labeled with arrows. Note that the origin of the coordinates along the jet axis has been shifted from $-6R_j$ to 0.
The initial velocity was set to \( \mathbf{v} \) [Mach 3.0]. Furthermore, we track the behavior of \( u \) (e.g., Meglicki 1994, 1995 for a detailed discussion), we have set \( u_{\text{init}} = 0 \) in SPMHD. As in Otmianowska-Mazur & Chiba (1995; see also 1996; Otmianowska-Mazur & Chiba 1995), for example, is not applicable yet in the context of the SPMHD, and the best way we have found to diminish the effects of a potential \( \mathbf{V} \cdot \mathbf{B} \neq 0 \) was to avoid those nonphysical accelerations by writing our system of equations in the vector formalism. Furthermore, we track the behavior of \( \mathbf{V} \cdot \mathbf{B} \) by evaluating the following quantity at each time step and position of the system (see, e.g., Otmianowska-Mazur & Chiba 1995):

\[
\omega = \frac{\left| \mathbf{V} \cdot \mathbf{B} \right| \cdot h}{|\mathbf{B}|}.
\]

As in Otmianowska-Mazur & Chiba (1995; see also Meglicki 1994, 1995 for a detailed discussion), we have set \( \omega \leq 10^{-2} \) as a limit to the validity of our calculations. In general, 85\% of the particles of the system keep \( \omega \leq 10^{-2} \) (with almost all of them distributed around the zero value most of the time). For small periods of time, some of them may reach \( \omega > 10^{-2} \), but the scalar product between the magnetic field and the magnetic force remains very small (\( \sim 10^{-5} \) to \( 10^{-7} \), in code units) over the entire evolution of the systems simulated here.

As additional tests to check the validity of the code modified by the \( \mathbf{B} \)-field implementation, we have run Alfvén wave tests, similar to those suggested by Clarke (1996). In the center of a rectangular box, a circular pulse of radius \( r = (y^2 + z^2)^{1/2} = 5h \) (where \( h \) is the smoothing length) and velocity \( \mathbf{v} = \mathbf{v}_{\text{init}} \) is set perpendicularly to the magnetic field. The latter was assumed to be in the \( y \)-direction in one case \( (\mathbf{B} = \mathbf{B}_{\text{init}} j) \) and \( \mathbf{B} = \mathbf{B}_{\text{init}} \cdot (j + k) \) in the other case. The initial velocity was set to \( v_{\text{init}}/v_a = (4\pi \rho)^{1/2} \times v_{\text{init}}/B_{\text{init}} = 10^{-3} \) and \( v_{\text{init}}/v_a = (\sqrt{2}/2)(4\pi \rho)^{1/2} \times v_{\text{init}}/B_{\text{init}} = (\sqrt{2}/2)10^{-3} \), for each case, respectively (see Clarke 1996). As in Clarke’s tests, after several time steps the diffusion of the pulses was found to be strictly confined to the direction of propagation of the Alfvén waves, as depicted in Figures 17 and 18 of Clarke (1996).

2.1. Initial and Boundary Conditions

The computational domain is represented by a 3-D rectangular box of dimensions \(-17R_j \leq x \leq 17R_j, -22R_j \leq y, z \leq 22R_j \) where \( R_j \) is the initial jet radius (\( R_j \) is the code distance unit). The Cartesian coordinate system has its origin at the center of the box and the jet flows through the \( x \)-axis and is continuously injected into the bottom of the box \([\text{at } r = (-17R_j, 0, 0)]\). Inside the box, the particles are initially distributed on a cubic lattice. An outflow boundary condition is assumed for the boundaries of the box. The particles are smoothed out by a spherically symmetric kernel function of width \( h \), and the initial values of \( h \) were chosen to be \( 0.4R_j \) and \( 0.2R_j \) for the ambient and jet particles, respectively.

We consider two different initial magnetic field configurations. One of them is an initially constant longitudinal \( \mathbf{B} \)-field permeating both the jet and the ambient medium \( \mathbf{B} = (B_{\text{eq}}, 0, 0) \). An observational support for this kind of configuration is suggested by the fact that some protostellar jets appear to be aligned with the main direction of the local interstellar magnetic field (see, e.g., Appenzeller 1989). The other adopted configuration is a force-free helical magnetic field distribution.

\[
\mathbf{B} = \mathbf{B}_{\text{hel}} = (B_{\text{hel}}, 0, 0),
\]

\[\text{with } B_{\text{hel}} = (2/3)(1/\gamma - 1)B_{\text{eq}} \text{ and } \gamma = (\gamma - 1)/\gamma.
\]

\( B_{\text{eq}} \) is the field strength in the absence of the jet. Following the development of the spherical jet model by Chef (1994), we choose \( B_{\text{eq}} = 1 \times 10^{-5} \text{ G} \) for \( \gamma = 5/3 \). As a model for the ambient medium, we adopt a uniform magnetic field \( \mathbf{B}_{\text{ambient}} = B_{\text{ambient}} \mathbf{n} \), where \( B_{\text{ambient}} = 1 \times 10^{-5} \text{ G} \).

\( \mathbf{n} \) is a random orientation, which is determined in the code by a function of \( \mathbf{x} \) such that \( 10^{-3} \leq \mathbf{n} \cdot \mathbf{x} \leq 1 \). In the initial conditions, the radial field component is \( B_r = 0 \).

A numerical solution to the MHD equations is obtained by integrating \( \mathbf{v} \) step by step in time and space. The steps are small enough to ensure that all magnetic field lines are reconnected, and the plasma is assumed to be incompressible. The code is run until the jet reaches a steady-state solution (in other words, after \( t \gg t_{\text{jet}} \)).
field that also extends to the ambient medium with a functional dependence given by (see also Todo et al. 1993)

\[ B_0 = 0, \]  

\[ B_\phi(r) = B_0 \left[ \frac{0.5Adr^2}{(r + 0.5d)^3} \right]^{1/2}, \]  

\[ B_x(r) = B_0 \left[ 1 - \frac{Ar^2(r + d)}{(r + 0.5d)^3} \right]^{1/2}, \]  

where \( r = (y^2 + z^2)^{1/2} \) is the radial distance from the jet axis and the (arbitrary) constants \( A \) and \( d \) are given by 0.99 and 3\( R_j \), respectively. In these equations, \( B_0 \) is the maximum strength of the magnetic field and corresponds to the magnitude of the longitudinal component at the jet axis. Figure 1 displays both the longitudinal and the toroidal \( B_\phi \) components as a function of the radial distance (in code units, for \( \beta = 1 \) at the jet axis). The pitch angle at 1\( R_j \) is \( \approx 19^\circ \).

The models are parameterized by the dimensionless numbers: (1) the density ratio between the jet and the ambient medium, \( \eta = n_j/n_a \); (2) the ambient Mach number, \( M_a = v_j/c_a \) (where \( v_j \) is the jet velocity and \( c_a \) is the ambient sound speed); (3) the jet-to-ambient medium pressure ratio at the jet inlet, \( \kappa = p_j/p_a \), which we assume to be equal to unity; (4) the thermal-to-magnetic pressure ratio, at the jet axis \( \beta = p_m/p_j \); and (5) \( d_{\text{cool}}/R_j \), the ratio of the cooling length in the postshocked gas behind the bow shock to the jet radius (see, e.g., GB93).

3. THE SIMULATIONS

As in previous work (see, e.g., GB96), based on typical conditions found in protostellar jets, we have adopted the following initial values for the parameters: \( \eta = 3–10 \) (see, e.g., Morse et al. 1992; Raga & Noriega-Crespo 1993), an ambient number density \( n_a = 200 \) cm\(^{-3} \), \( v_j = 398 \) km s\(^{-1} \) (see, e.g., Reipurth, Raga, & Heathcote 1992), \( M_a = 24 \), and \( R_j = 2 \times 10^{15} \) cm (see, e.g., Raga 1993). In the MHD simulations, we have assumed an initial \( \beta = 1 \) (which corresponds to a maximum initial value \( B_0 = 83 \) \( \mu \)G). The subsections below present the results of the simulations we have performed for both radiatively cooling and adiabatic jets with and without magnetic fields, and Table 1 summarizes the values of the input parameters. In Table 1, \( M_{\text{Alf}} = v_j/v_{A_j} \) and \( M_{\text{mas}} = v_j/(v_{A_j}^2 + c_s^2)^{1/2} \) give the initial Alfvén and magnetosonic Mach numbers, respectively. The purely hydrodynamical models are labeled “HD,” and the MHD models, with an initially longitudinal or helical magnetic field configuration, are labeled “ML” and “MH,” respectively.
3.1. Radiatively Cooling Jets

Figure 2 depicts the time evolution of the magnetic field distribution of two supermagnetosonic, radiatively cooling jets: ML3r (an MHD model with initial longitudinal magnetic field configuration; right column) and MH3r (an MHD model with initial helical magnetic field configuration; left column). Both models have initial $\eta = n_j/n_a = 3$ and ambient Mach number $M_a = 24$ (see Table 1). We find that the field lines are amplified and reoriented across the bow shock in both magnetic field configurations. The shocked jet material that is decelerated at the head and is deposited in the cocoon carries the field lines embedded in it. Part of these lines have their polarity reversed, as we can see in Figure 2.

The cooling parameters behind the bow shock, $q_{bs} \approx 8.0$, and the jet shock, $q_{js} \approx \eta^{-3}q_{bs} \approx 0.3$ (see, e.g., GB93) in Figure 2, imply that, within the head of the jet, the ambient shocked gas is almost adiabatic, whereas the shocked jet material is subject to rapid radiative cooling. The corresponding density contours and velocity field distribution maps of the models above were presented in Figures 1 and 2 of Paper I, where they were also compared with a pure hydrodynamical model (HD3r; see also Table 1). As we have stressed in Paper I, the cold dense shell that develops at the jet head because of the cooling of the shock-heated jet material in the pure hydrodynamical case also appears in the MHD jets. Similarly, it becomes Rayleigh-Taylor (R-T) unstable (see, e.g., GB93) and breaks into blobs that spill into the cocoon and resemble the Herbig-Haro objects that are observed at the head of protostellar jets. As in the HD case, the density in the shell of the MHD jets also undergoes fluctuations with time that are caused by global thermal instabilities of the radiative shock (for details, see, e.g., Gaetz, Edgar, & Chevalier 1988; GB93). Figure 3 shows the time evolution of the shell-density ($n_{sh}$) variations at the jet axes of the three jets, which have a period of the order of the gas cooling time ($t_{cool} \approx 10$ yr, or $t_{cool} \approx 0.3t_d$, where $t_d = R_j/c_a \approx 38$ yr corresponds to the transverse jet dynamical time). We note that although the HD and the MHD jets attain a maximum density of approximately the same magnitude at later on $t/t_d \approx 1.1$, the density enhancement in the MHD jets is inhibited by the presence of the $B$-field, particularly in the jet with a helical field (MH3r).

Figure 2 indicates that the cold blobs that develop from shell fragmentation at the head of the MHD jets detach...
from the beam as they are expelled backward to the cocoon. The survival of these blobs in the cocoon seems to suffer with the presence of the magnetic fields, particularly in the helical case. Compared to that of the blobs of the HD jet, their density is reduced by a factor of \( \sim 2 \) in the longitudinal case (ML3r) at the final time step \( (t/t_d = 1.65) \) and almost vanishes in the helical jet (MH3r). Figure 4 compares the density and pressure profiles across the flow in three different positions along the MHD jets in the head region. The lateral blobs that develop in the cocoon, on both sides of the beam, are clearly less intense in the helical case than in the longitudinal case. In the helical case, for example, we find...
at the heads of the jets is (maximum) is given by black, light gray, white, and dark gray. The tions are the same as in Fig. 6. The gray scale (from minimum to
size against the R-T instability and the clump formation is
density growth. As a consequence, the shell tends to stabil-
the cooling length behind the jet shock and reduces the
Ðed by compression in the shocks at the head (by a factor of
magnetic Ðeld, which is initially less intense than the longi-
tudinal component (by a factor of
that the toroidal (not parallel-to-shock) component of the magnetic field, which is initially less intense than the longitudi-
tional component (by a factor of \( \leq 2.5 \)), is strongly ampli-
fied by compression in the shocks at the head (by a factor of \( \sim 5 \)). As pointed out in Paper I, this amplification increases the cooling length behind the jet shock and reduces the density growth. As a consequence, the shell tends to stabil-
ize against the R-T instability and the clump formation is
inhibited. At the end of the evolution, the dense shell with clumpy structure that was observed to develop in the head region of the pure HD jet (Fig. 1 of Paper I) is replaced in the helical case by an elongated plug (Fig. 2) of low-density material.

Figure 2 also indicates the development of some pinching along both MHD jets. In the pure HD case, constriction occurs only very close to the jet head where the beam is overconfined by the gas pressure of the cocoon (see Figs. 1 and 2 of Paper I; see also right column of Fig. 5). In the MHD jet with a helical field (MH3r), the toroidal component \( (B_\phi) \), which is amplified by compression in the shocks at the head, is advected back with the shocked material to the cocoon. The associated magnetic pressure \( (\sim B_\phi^2/8\pi) \) causes an increase in the total pressure of the cocoon relative to the pure HD case, which collimates the beam and excites the (fastest growing) small-wavelength pinch modes of the MHD K-H instability. These modes overconﬁne the beam and drive the approximately equally spaced internal shocks seen in the MH3r jet (Fig. 2; left column). Along the MHD jet with longitudinal field (ML3r; Fig. 2; right column), the increase in the total confining pressure of the cocoon also drives the development of the MHD K-H instabilities that excite beam pinching and internal shocks. Consistent with linear theory for K-H modes, in the supermagnetosonic regime considered here, they begin to appear at a distance \( \sim M_{ms}R_j \), which is smaller than in the pure HD jet (see Table 1).

The presence of these oblique internal shocks along the beam of the MHD jets can be testiﬁed to by the density and pressure proﬁles across and along the ﬂow depicted in Figure 5, where the pure HD model is also depicted for comparison. We see that the induced internal shocks in the MHD cases have a density contrast of \( n_{in}/n_j \approx 4 \) in the ML3r model (middle row) and \( n_{in}/n_j \approx 5 \) in the MH3r model (or \( n_{in}/n_{sh} \approx 0.08 \) and 0.15, respectively, where \( n_{sh} \) is the density at the shell). We also note a close correlation between the pinching zones and the appearance of more intense reversed fields in the contact discontinuity between the jet and the cocoon. In both magnetic ﬁeld conﬁgurations, the original longitudinal components are reoriented in the nonparallel shocks at the head and advected back to the cocoon. As a consequence, a predominantly toroidal current density distri-

Fig. 10.—Gray-scale representation of the midplane density of adia-
batic jets with \( \eta = 3 \): a hydrodynamical jet (top panel; model HD3a), an MHD jet with initial longitudinal magnetic ﬁeld distribution (middle panel; model ML3a), and an MHD jet with initial helical magnetic ﬁeld distribution (bottom panel; model MH3a), at a time \( t/t_f = 1.65 \). The initial conditions are the same as in Fig. 6. The gray scale (from minimum to maximum) is given by black, light gray, white, and dark gray. The maximum density reached at the heads of the jets is \( n_{in}/n_j \approx 37 \) (top panel), \( n_{in}/n_j \approx 39 \) (middle panel), and \( n_{in}/n_j \approx 42 \) (bottom panel).
the shroud that envelopes the cocoon/beam, while the shocked jet material cools much faster \((q_{js} > 1)\) and the shock is effectively isothermal. The cold shell is thus much thinner in the larger \(\eta\) jets and the head resembles a bullet.

As in the \(\eta = 3\) case, the density in the shell of the \(\eta = 10\) jets also undergoes variations with time (with a period of the order of the jet radiative cooling time), which are caused by the global thermal instabilities of the radiative shocks. Likewise, after reaching a similar maximum density amplitude, the MHD jets have their shell-density growth inhibited because of the decrease of the shock compressibility caused by the presence of the \(B\)-field, particularly in the jet with a helical field (MH10r). We note, however, that in the HD jet, the shell-density variations attain a smaller maximum density amplitude than in the MHD cases. This is possibly because of the smaller total pressure confinement that acts on its head.

In the \(\eta = 10\) MHD jets (ML10r and MH10r), the increase in the total confining pressure in the cocoon (due to the presence of the magnetic field) also overconfines the beam, relative to the HD jet, and drives some beam pinching.

Figure 9 (left column) shows the density and pressure profiles across the flow in different pinch positions along the MHD and HD \(\eta = 10\) jets after they have propagated over a distance \(33R_j\) (at \(t/t_d = 1.65\)); this can be compared with Figure 5 (for \(\eta = 3\)). As in the \(\eta = 3\) models, the pinch collimation is larger in the helical MH10r jet than in the ML10r jet, while the pure hydrodynamic jet, HD10r, has not developed any internal pinches over the timescale depicted. It is interesting to note that in the MHD jets, the fastest growing pinch modes of the K-H instability would be expected to appear only at distances \(~M_{ms_j}R_j \sim 50R_j\) which are beyond the computed scales. Thus the early development of pinches in these cases, is possibly being triggered mainly by the hydromagnetic \(\theta\)-pinch effect (see, e.g., Boyd & Sanderson 1961; Cohn 1983). As in the \(\eta = 3\) case, the jet with longitudinal \(B\)-field also develops a nonaxisymmetric helical (kink) mode close to the head of the jet, which causes some beam twisting.

As expected, after propagating about the same distance, the pinches that develop in the larger \(\eta\) case are less numerous, since the total amount of confining shocked material that deposits into the cocoon is comparatively smaller (see Figs. 5 and 9 for a comparison). Nonetheless, the density contrasts, \(n_{is}/n_j\), attained in the pinch regions of the MHD jet with a longitudinal field, ML10r (\(n_{is}/n_j \approx 3-4.5\), or \(n_{is}/n_{sh} \sim 0.05-0.07\)) are approximately the same as those in the smaller \(\eta\) jet (ML3r). The situation is a little more complex for the MHD jets with helical field. Over the whole jet evolution, some few pinches are found to be stronger in the larger \(\eta\) jet (MH10r; Fig. 8), but in general the density
contrasts are approximately the same in both cases
\( n_{\text{sh}}/n_j \lesssim 5 \), or \( n_{\text{sh}}/n_j \lesssim 0.13 \).

3.2. Adiabatic Jets

Figure 10 depicts the density in the midplane section of the head of three supermagnetosonic, adiabatic jets with \( \eta = 3 \): a purely HD jet (HD3a, top panel); an MHD jet with initial longitudinal \( B \)-field, (ML3a, middle panel); and an MHD jet with initial helical \( B \)-field (MH3a, bottom panel). Figures 11 and 12 show the time evolution of the corresponding velocity and magnetic field distributions. The initial conditions are the same as in Figure 2 (see Table 1). Previous numerical analyses comparing pure HD jets with and without radiative cooling (see, e.g., BFK, GB93) have shown that the presence of radiative cooling tends to reduce the strength and number of internal shocks excited by K-H instability. Consistently, the HD3a jet in Figures 10, 11, and 12 reveals the appearance of a pinching zone (at \( x \approx 1R_j \)), which is absent in its radiatively cooling counterpart (HD3r; Fig. 2; see also Figs. 1 and 2 of Paper I). Besides, the beam constriction that appears close to the head is larger in the adiabatic jet (HD3a). These results are also compatible with the predictions of the linear stability theory (Hardee & Stone 1997), when applied in the context of the cooling function employed in this work (see §1). Similarly, the pinches that develop in the MHD adiabatic jets (Figs. 10, 11, and 12) are generally more intense and numerous than in the radiatively cooling counterparts of Figure 2. Their larger strength can be testified by direct comparison of Figures 13 and 5, which show the transverse profiles of some pinches along the jets. In particular, for the ML3a jet, we find a chain of eight evolving pinches, at \( t/t_d = 1.65 \), against five in the radiatively cooling jet (ML3r), and the corresponding densities \( n_{\text{sh}}/n_j \), are about 2–3 times greater than in the ML3r jet. The helical adiabatic jet MH3a displays pinches that are 40% denser than those found in the radiatively cooling jet (MH3r) at \( t/t_d = 1.65 \).

We find similar results when comparing radiatively cooling and adiabatic jets with larger \( \eta \). Figure 14 shows the midplane density (left column) and velocity field distribution
Fig. 13.—Left column: Density (solid lines) and pressure (dashed lines) profiles across the HD3a (top row), ML3a (middle row), and MH3a (bottom row) jets taken in three different pinching positions along the flow (at \( t/t_d = 1.65 \)). The profiles have been scaled as in Figs. 4 and 5. Right column: Corresponding axial density along the jet axis (\( y = z = 0 \)), showing the channel of internal shocks. The positions of the three pinches depicted in the left column are labeled with arrows.

(right column) for adiabatic jets with \( \eta = 10 \) after they have propagated \( \simeq 33R_j \) with the same initial conditions as the cooling jets of Figures 6, 7, and 8. The corresponding magnetic field distribution is presented in Figure 15. As before, the \( \eta = 10 \) adiabatic jet with longitudinal \( B \)-field (ML10a) develops stronger pinches than its cooling counterpart (ML10r model) by a factor of \( \approx 3 \). The \( \eta = 10 \), adiabatic helical jet (MH10a), on the other hand, has pinches with
densities of the same order of magnitude as those in the cooling jet (MH10r). Also, the pinches are found to be stronger in the adiabatic jets with smaller $\eta$, in both magnetic field configurations.

4. CONCLUSIONS AND DISCUSSION

We have investigated here the effects of magnetic fields on the structure of evolving overdense, radiatively cooling, supermagnetosonic jets with the help of 3-D SPMHD simulations and compared these with purely hydrodynamical and adiabatic calculations. Two initial magnetic field configurations (with magnitude in approximate equipartition with the gas) have been examined: a longitudinal and a helical field permeating both the jet and the ambient medium. Calculated for a set of parameters that are particularly appropriate to protostellar jets (with density ratios between the jet and the ambient medium $\eta \approx 3-10$ and ambient Mach number $M_a \approx 24$), our results indicate that magnetic fields have important effects on the dynamics of radiatively cooling jets. Both magnetic field geometries are able to improve jet collimation relative to the pure hydrodynamical (HD) jets, but this effect is larger in the helical field case.

As we have stressed in Paper I, the cold dense shell that develops at the jet head because of the cooling of the shock-heated jet material in the HD cases, also appears in the MHD jets. Likewise, it becomes Rayleigh-Taylor (R-T) unstable and breaks into clumps that are more visible in the smaller $\eta$ jets where the developed shell is thicker. Also as in the HD case, the shell of the MHD jets undergoes density variations with time, which are caused by global thermal instabilities of the radiative shocks. However the amplification and reorientation of the nonparallel components of the magnetic fields by the shocks at the head, particularly in the helical field geometry, reduces the postshock compressibility and increases the postshock cooling length. This tends to stabilize the shell against both R-T and thermal instabilities. As a consequence, the clumps that are observed to develop by fragmentation of the shell in the HD jets are depleted in the helical field geometry. The jet immersed in the longitudinal field, on the other hand, still retains the clumps, although they have their densities decreased relative to the HD counterparts. The fact that the clumpy shell structure resembles the knotty pattern of the Herbig-Haro objects that are commonly observed at the heads of protostellar jets (see, e.g., Herbig & Jones 1981; Brugel et al. 1985; Reipurth 1989; Heathcote et al. 1996) suggests that a longitudinal magnetic field geometry would be more likely in the outer regions of these jets than a helical field geometry (see also Paper I).
Over the computed time and length scales, internal oblique shocks along the beam are not found to develop in the HD systems examined here. On the other hand, in their MHD counterparts with both magnetic field configurations, the confining total pressure of the cocoon excites (the fastest growing) low-amplitude MHD K-H reflection pinching modes that drive a chain of approximately equally spaced internal shocks along the beam, but these shocks are found to be slightly stronger in the helical-field case (by a factor of \( \approx 20\% \)). Also, as expected, the internal shocks tend to appear in larger numbers in the smaller \( \eta \) jets (because of the larger amount of confining shocked material that is deposited in the cocoon), although their densities are of approximately the same magnitude as those in the larger \( \eta \) jets. A nonaxisymmetric helical mode is also excited close to the head of the radiatively cooling MHD jets with longitudinal fields, causing some beam wiggling.

The number and strength of the internal shocks excited in the MHD adiabatic jets are larger than in the radiatively cooling counterparts (by a factor of \( \approx 2 \) in both number and strength). This result is compatible with the linear stability theory (Hardee & Stone 1997), when applied in the context of the cooling function employed in the present work, and is also compatible with previous numerical work of HD jets, which has shown that the presence of radiative cooling tends to reduce the strength and number of internal shocks along the jet (see, e.g., BFK; GB93). Also, the pinches are found to be more numerous in the adiabatic jets with smaller \( \eta \), in both magnetic field configurations.

The internal shocks are found to propagate downstream with velocities close to that of the jet head \( (v_{\text{sh}} \approx 250 \text{ km s}^{-1}) \). The mean distance between them \( (\approx 2-4R_j) \) is in agreement with the observed knots in the jets. However, the weakness of the shocks in the radiatively cooling jets \( (n_{\text{sh}}/n_j \approx 3-5) \) makes it doubtful that they could produce by themselves the bright knots observed in protostellar jets. Probably, other mechanisms, like intermittence in the jet injection velocity, play a more relevant role in knot formation in those jets (see, e.g., Raga et al. 1990; de Gouveia Dal Pino and Benz 1994; Stone and Norman 1993a).

We should note that the recent numerical study of magnetized radiatively cooling jets by Frank et al. (1998) found that toroidal fields (also in approximate equipartition with the gas) may excite the development of strong pinches along the beam. This apparent contradiction between their analysis and ours is possibly due to differences in the assumed initial conditions. Frank et al. have considered a slower and much lighter jet (with a jet Mach number \( M_j = v_j/c_s \approx 10 \) and \( \eta = 1.5 \)) than in the cases examined here \( (M_j = v_j/c_s \approx 42-76 \) and \( \eta = 3-10 \)). Thus, consistently with our results above, their smaller \( \eta \) jet should be expected to produce a larger amount of more intense pinches. This result has also been confirmed by numerical simulations of \( \eta = 1 \) jets (not presented here) that we have performed, which have produced a larger amount of pinches with slightly larger densities than the larger \( \eta \) jets studied above.

Finally, we should make some remarks on the late evolution of the radiatively cooling magnetized jets. In the magnetic field maps of Figures 2 and 8 above, we have detected the development of (sometimes strong) magnetic field reversals at the contact discontinuity between the jet and the cocoon with intensities up to 5 times their initial magnitudes. As stressed in \( \S \) 3, field reversals occur in both investigated magnetic field configurations, because the field lines are amplified by compression in the nonparallel shocks at the jet head and are forced to flow backward with the shocked plasma into the cocoon. In this process, the lines are reoriented and sometimes have their polarization reversed. Beyond the integrated length scales and timescales depicted in the figures above, however, the increasing strength of the reversed fields due to shear at the contact discontinuity (see, e.g., eq. [2c]) may lead to the development of strong pinching regions, which ultimately may cause jet disruption, particularly in the cases with longitudinal fields. As an example, Figure 16 depicts the time evolution of the velocity (left column) and magnetic field (right column) distributions of a disruption zone that occurred in the late evolution of the ML3r jet of Figure 2. We can clearly distinguish two regions with reversed \( B \)-fields whose strength increases with time, which are correlated with developing pinches. The inner constriction becomes so strong that it finally causes the disruption of the beam. Although shear and compression are expected to enhance \( B \), this amplification is possibly partially due to numerical effects. In fact, along a contact discontinuity with such a magnetic field topology with oppositely directed field lines pressed together, magnetic diffusion and reconnection may have an important role and lead to intense magnetic energy release. Of course, under the ideal-MHD approach considered here, these dissipation effects of the magnetic field have not been appropriately considered, thus leading to possibly anomalous amplification of the reversed components in the late stages of the evolution of some of the jets. Further, we should note also that no jet disruption was
detected in the majority of the adiabatic cases. This fact is consistent with recent numerical studies of reconnection processes in two-dimensional (2-D) current sheets (Oreshina & Somov 1998), which have shown that reconnection rates are smaller in adiabatic than in radiatively cooling plasmas. Although the transposition of these results to the more complex flow geometry we have investigated here is not straightforward, they seem to suggest that the appropriate consideration of magnetic field dissipation in our models will possibly decrease and even suppress the disruptive effects of the magnetic fields found in some of the radiatively cooling cases examined here in their late evolution. The transformation of magnetic energy into thermal energy of the gas will probably have important effects on the structure and emission mechanisms of the beam (see, e.g., Malagoli et al. 1996; Min 1997a, 1997b; Jones et al. 1997) and also in the process of turbulent mixing of the jet and cocoon material. Such finite magnetic resistivity effects will be addressed in a forthcoming paper. Besides, the potential signatures that magnetic fields may leave on the morphology of radiative cooling jets, especially behind the shocks, provide important constraints that can be used in
future observational tests to distinguish among different candidate mechanisms for emission-line production, jet collimation, and turbulent entrainment at the contact discontinuity between the jet and the cocoon.

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