SIMULATING PROTOSTELLAR JETS SIMULTANEOUSLY AT LAUNCHING AND OBSERVATIONAL SCALES

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ABSTRACT

We present the first 2.5-dimensional magnetohydrodynamic (MHD) simulations of protostellar jets that include both the region in which the jet is launched magnetocentrifugally at scale lengths $<0.1$ AU and where the propagating jet is observed at scale lengths $>10^3$ AU. These simulations, performed with the new adaptive mesh refinement MHD code AZEuS, reveal interesting relationships between conditions at the disk surface, such as the magnetic field strength, and direct observables such as proper motion, jet rotation, jet radius, and mass flux. By comparing these quantities with observed values, we present direct numerical evidence that the magnetocentrifugal launching mechanism is capable, by itself, of launching realistic protostellar jets.

Key words: ISM: jets and outflows – magnetohydrodynamics (MHD) – stars: formation

1. INTRODUCTION

Jets and outflows from protostellar objects are fundamental aspects of the current star formation paradigm, and are observed anywhere star formation is ongoing. The mechanism proposed by Blandford & Payne (1982), in which jets are launched from accretion disks by gravitational, magnetic, and centrifugal forces, has been extensively studied numerically (e.g., Uchida & Shibata 1985; Meier et al. 1997; Ouyed & Pudritz 1997a, 1997b, 1999; Krasnopolsky et al. 1999; Vitorino et al. 2002; von Rekowski et al. 2003; Ouyed et al. 2003; Porth & Fendt 2010; Staff et al. 2010). By treating the accretion disk as a boundary condition (e.g., Ustyugova et al. 1995), one can study jet dynamics independently of the disk (e.g., Pudritz et al. 2007) though, in order to resolve the launching mechanism, numerical simulations have not followed the jet beyond 100 AU (e.g., Anderson et al. 2005).

In stark contrast, protostellar jets are $\gtrsim 10^4$ AU long (Bally et al. 2007) and only recently have observations reached within 100 AU of the source (e.g., Hartigan et al. 2004; Coffey et al. 2008). This large-scale difference between observations and simulations makes direct comparisons difficult and, in this work, we aim to close this gap. We present axisymmetric (2.5-dimensional) simulations of protostellar jets launched from the inner AU of a Keplerian disk and follow the jet well into the observational domain (2500 AU). These calculations allow us to address the efficacy of the magnetocentrifugal mechanism and to relate conditions near the disk with directly observable properties of the jet.

The simulations presented herein are performed with an adaptive mesh refinement (AMR) version of ZEUS-3D (Clarke 1996, 2010) called AZEuS (Adaptive Zone Eulerian Scheme). The ZEUS-3D family of codes is among the best tested, documented, and most widely used astrophysical magnetohydrodynamic (MHD) codes available, though this is the first attempt to couple ZEUS-3D with AMR. We have implemented the block-based method of AMR detailed in Berger & Colella (1989) and Bell et al. (1994). Significant effort was spent minimizing errors caused by passing waves across grid boundaries, which is of particular importance to this work. A full description of the code and the changes required for AMR on a fully staggered mesh will appear in J. P. Ramsey & D. A. Clarke (2011, in preparation).

2. INITIALIZATION

Observationally, the inner radius of a protostellar accretion disk, $r_i$, is between 3 and 5 $R_*$ (Calvet et al. 2000) and, for a typical T Tauri star ($M = 0.5 M_\odot$, $R_* = 2.5 R_\odot$), $r_i = 0.05$ AU. Thus, following Ouyed & Pudritz (1997a), we initialize a hydrostatic, force-free atmosphere surrounding a 0.5 $M_\odot$ protostar coupled to a rotating disk with $r_i = 0.05$ AU. However, unlike Ouyed & Pudritz we use an adiabatic equation of state that conserves energy across shocks rather than an isentropic polytropic equation of state, as the distinction becomes important for supermagnetosonic flow (Ouyed et al. 2003).

We solve the equations of ideal MHD (e.g., $\gamma = 5/3$) over a total domain of 4096 AU $\times$ 256 AU. To span the desired length scales, nine nested, static grids (refinement ratio 2) are initialized each with an aspect ratio of 4:1 (16:1 for the coarsest grid only) and bottom left corner at the origin. Our finest grid has a domain 4 AU $\times$ 1 AU and a resolution $\Delta z = r_i/8 = 0.00625$ AU which we find sufficient to resolve the launching mechanism. Thus, the effective resolution for the entire domain is $>26$ billion zones.

During the simulations, a thin region of low velocity and high poloidal magnetic field, $B_\phi = \sqrt{B_z^2 + B_\phi^2}$, develops along the symmetry axis, the edge of which is defined by a large gradient in the toroidal magnetic field, $\partial_z B_\phi$. Insufficient resolution of $\partial_z B_\phi$ can lead to numerical instabilities, and grids are added dynamically whenever this gradient is resolved by fewer than five zones.

2.1. The Atmosphere

The atmosphere is initialized in hydrostatic equilibrium (HSE; $v_z = v_\phi = 0$). Because the left-hand side of the equation governing HSE,

$$\nabla p + \rho \nabla \phi = 0,$$

1 ENZO, a hybrid $N$-body Eulerian code (O’Shea et al. 2004), links AMR with the hydrodynamical portion of ZEUS-2D.

2 AZEuS solves either the total or the internal energy equation. We chose the latter because positive-definite pressures trump strict conservation of energy in these simulations; see Clarke (2010).
is not a perfect gradient, differencing it directly on a staggered-mesh can commit sufficient truncation error to render the atmosphere numerically unstable. Thus, we replace $\nabla \phi$ with the corresponding poloidal gravitational acceleration vector,

$$\vec{g} = -\frac{1}{\rho_h} \nabla p_h,$$

where $\rho_h$ and $p_h$ are the hydrostatic density and pressure given by

$$\rho_h = \rho_i \left( \frac{r_i}{\sqrt{r^2 + z^2}} \right)^\gamma \quad \text{and} \quad p_h = \frac{p_i}{\gamma} \left( \frac{\rho_h}{\rho_i} \right)^\gamma. \quad (3)$$

Here, $\rho_i$ and $p_i$ are the initial density and pressure at $r_i$ and $p \propto \rho^\gamma$ is assumed throughout the atmosphere at $t = 0$. In this way, differencing Equation (1) maintains HSE to within machine round-off error indefinitely.

However, Equations (3) as given are singular at the origin where truncation errors are significant regardless of resolution. These errors can launch a supersonic, narrow jet from the origin destroying the integrity of the simulation. To overcome this problem, we replace the point mass at the origin with a uniform sphere of the same mass and a radius $R_0$, thus modifying the first of Equations (3) to

$$\left( \frac{\rho_h}{\rho_i} \right)^\gamma = \begin{cases} \frac{r_i}{\sqrt{r_i^2 + z^2}}, & r^2 + z^2 \geq R_0^2; \\ \frac{r_i}{\sqrt{r_i^2 + z^2}} - \frac{3R_0^2 - r^2 - z^2}{2R_0^2}, & r^2 + z^2 < R_0^2. \end{cases} \quad (4)$$

If $R_0$ is sufficiently resolved (e.g., four zones), the numerical jet is eliminated. The resulting “softened potential” is superior to a “softened potential” since the former has no measurable effects beyond $R_0$. Here, we use $R_0 = r_i$.

The atmosphere is initialized with the force-free magnetic field used by Ouyed & Pudritz (1997a):

$$A_\phi = \frac{B_i}{\sqrt{r^2 + (z + z_d)^2 - (z + z_d)}}; \quad B_z = \frac{1}{r} \frac{\partial (r A_\phi)}{\partial r}, \quad B_{\phi} = -\frac{\partial A_\phi}{\partial z}, \quad B_\psi = 0,$$

where $A_\psi$ is the vector potential, $z_d$ is the disk thickness (set to $r_i$), and $B_i$ is the magnetic field strength at $r_i$, given by

$$B_i = \left( \frac{8\pi p_i}{\beta_i} \right)^{1/2}. \quad (6)$$

Here, $p_i$ and $\beta_i$ (plasma beta at $r_i$) are free parameters.

Finally, to ensure that the declining density and magnetic field profiles do not fall below observational limits, we add floor values $\rho_{\text{floor}} \sim 10^{-6} \rho_i$ and $B_{\text{floor}} \sim 10^{-5} B_i$ (cf. Bergin & Tafalla 2007; Vallée 2003) to Equations (4) and (5). By imposing HSE and the adiabatic gas law at $t = 0$, a floor value on $\rho$ imposes effective floor values on $\vec{g}$ and $p$ as well.

2.2. Boundary Conditions

In the accretion disk ($z \leq 0, r \geq r_i$), $v_\psi = v_K = \sqrt{GM_s/r}$, the Keplerian speed, and $v_c = \xi v_K = 10^{-7} v_K$ is an “evaporation speed” at the disk surface. The disk and atmosphere are initially in pressure balance with a density contrast $\eta = \rho_{\text{disc}}/\rho_{\text{atm}} = 100$, while $\vec{B}$ is initialized using Equations (5).

Following Krasnopolsky et al. (1999), $\rho$, $p$, and $v_\psi$ are held constant, $v_r = v_\psi B_r / B_\psi$, $v_\theta = v_K + v_\psi B_\theta / B_\psi$, $E_z(z)$ where $\vec{E} = \vec{v} \times \vec{B}$ is the induced electric field. $E_r(0) = v_K B_\psi(0)$, $E_\theta(0) = E_z(0) - E_z(z)$, $E_\phi(0) = 0$, and $E_\psi(0) = -E_\psi(z)$. Since $\vec{v}_i$ is subsonic, these conditions are formally overdetermined and $p$ should probably be allowed to float. Indeed, we allowed $p$ to be determined self-consistently in test simulations and found only minor quantitative differences in the jet since the pressure gradient is only about 1% of the net Lorentz force at the disk surface. However, allowing $p$ to float in the boundary caused undue high temperatures in the disk, and thus small time steps. Therefore, the simulation proceeds more rapidly but otherwise virtually unchanged when $p$ is maintained at its initial value.

Inside $r_i$ ($z \leq 0$), we apply reflecting, conducting boundary conditions ($\vec{J} = \nabla \times \vec{B} \neq 0$). Thus, $\rho$, $p$, and $\vec{v}$ are reflected across $z = 0$, and magnetic boundary conditions are set according to $E_z(-z) = -E_z(z)$, $E_r(-z) = E_r(z)$, and $E_\psi(-z) = E_\psi(z)$. At $z = 0$, $E_r$ and $E_\psi$ are evolved using the full MHD equations.

Finally, we use reflecting boundary conditions along the $r = 0$ symmetry axis, and outflow conditions along the outermost $r$ and $z$ boundaries. Since nothing reaches these outermost boundaries, the nature of the outflow conditions is moot.

2.3. Scaling Relations

From Equation (1) and the adiabatic gas law, one can show

$$c_s^2 = \gamma \frac{p}{\rho} = (\gamma - 1) \frac{GM}{R} = (\gamma - 1) v_K^2,$$

where $R$ is the spherical polar radius. From Equations (6) and (7), and the ideal gas law ($p = kT/(m)$, where $(m)$ is half a proton mass), we derive the following scaling relations:

$$\rho_i = (160 \text{ dyne cm}^{-2}) \left( \frac{\beta_i}{40} \right) \left( \frac{B_i}{10 \text{ G}} \right)^2; \quad (8)$$

$$\frac{\rho}{(m)} = (5.4 \times 10^{12} \text{ cm}^{-3}) \left( \frac{\beta_i}{40} \right) \left( \frac{B_i}{10 \text{ G}} \right)^2 \left( \frac{r_i}{0.05 \text{ AU}} \right) \times \left( \frac{0.5 M_\odot}{M_\ast} \right); \quad (9)$$

$$T_i = (2.2 \times 10^5 \text{ K}) \left( \frac{0.05 \text{ AU}}{r_i} \right) \left( \frac{M_\ast}{0.5 M_\odot} \right); \quad (10)$$

$$c_{s,i} = (77 \text{ km s}^{-1}) \left( \frac{0.05 \text{ AU}}{r_i} \right)^{1/2} \left( \frac{M_\ast}{0.5 M_\odot} \right)^{1/2}; \quad (11)$$

$$r_i = \left( \frac{r_i}{c_{s,i}} \right) = (9.7 \times 10^4 \text{ s}) \left( \frac{r_i}{0.05 \text{ AU}} \right)^{3/2} \left( \frac{0.5 M_\odot}{M_\ast} \right)^{1/2}, \quad (12)$$

for $\gamma = 5/3$. Note that $\beta_i$ is the only free parameter varied in this work.

3. RESULTS FOR $\beta_i = 40$

Figure 1 depicts a jet with $\beta_i = 40$ at $t \simeq 100$ yr from the highest resolution grid near the disk surface (bottom panel) to the coarsest grid in which the jet has reached a length of...
just under 2500 AU (top panel). A few features worth noting include the following.

1. When $\theta < 60^\circ$ (angle between $\vec{B}_p$ and disk surface), Blandford & Payne (1982) show that cold gas near the disk is launched into a collimated outflow. Here, $\theta < 60^\circ$ for all $r > r_i$, but significant outflow is limited to inside the point where the slow surface intersects the disk ($r_{j,d} \sim 30$ AU = jet radius at the disk; the second panel from top). Below $r_{j,d}$, cold disk material has moved onto the grid and accelerated into the outflow. Above $r_{j,d}$, the weak magnetic field has yet to drive enough disk material onto the grid to displace the hot atmosphere, and outflow is stifled. While $r_{j,d}$ gradually increases with time, the majority of mass flux originating from the disk is driven within $r_i < r < 10 r_i$ (0.5 AU; bottom panel).

2. Jet material becomes superfast ($M_f \lesssim 5$) within a few AU of the disk, and the boundary between jet and entrained ambient material is defined by a steep temperature gradient (contact discontinuity; second panel). Portions of the original atmosphere, which remain virtually stationary throughout the simulation, are still visible above and ahead of the bow shock (top panel).

3. At large distances from the disk ($\gtrsim 500$ AU; top panel), the dynamics of the jet become dominated by $B_\phi$, and the jet is led by an essentially ballistic, magnetic “nose-cone” with a Mach number of $\sim 10$ (e.g., Clarke et al. 1986). Still, $B_\phi$ is a small fraction ($10^{-3}$) of $B_i$, consistent with Hartigan et al. (2007).

4. The knots dominating the bottom panel (cf., Ouyed & Pudritz 1997b) are produced by the nearly harmonic oscillation of $\vec{B}_p$ in $r_i < r < 2 r_i$, whereby $\theta$ fluctuates between $55^\circ$ and $65^\circ$ with a period $\sim 30 \tau_i$. These oscillations result from the interplay between infalling material along the symmetry axis, and under/over pressurisation near the central mass. The knots are denser and hotter than their surroundings, and bound by magnetic field loops. They occupy a region within $\sim 2$ AU of the symmetry axis and gradually merge to form a continuous and narrow column of hot, magnetized material (third panel). As such, they are unlikely to be the origin of the much larger-scale knots observed in some jets (e.g., HH111; Raga et al. 2002).

Further details of this and other simulations of protostellar jets are left to a future paper, and we focus here on a few properties directly comparable with observations.

### 4. COMPARING SIMULATIONS AND OBSERVATIONS

Table 1 summarizes a few observational characteristics of protostellar jets. To connect these attributes to conditions in the launching region, we have performed a small parameter survey in $\beta_i$ and made numerical measurements of the quantities in Table 1. Variation of other parameters (such as $\zeta$ and $\rho_i$) is left to future work.

Note that $\beta_i$ is the initial value of the plasma beta at $r_i$, and not the average $\beta$ in the jet. Indeed, Figure 2(a) demonstrates that at very early time, $\langle \beta \rangle = \frac{8 \pi \langle p \rangle}{\langle B^2 \rangle} \lesssim \beta_i/5$, where $B^2 = B_p^2 + B_\phi^2$, and where the average is taken over zones

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3 Time-lapse animations are available at http://www.ica.smu.ca/zeus3d/rc10/.

4 The knots are resolved by 10–20 zones when they merge, and thus their merger is unlikely related to the ever-decreasing resolution of the nested grids.
that exceed a certain threshold $v_z$ so that only outflowing jet material is considered. Thus, the magnetic field within the jet is stronger than $\beta_i$ would suggest. Initially, $(\beta)$ is dictated by $B_p$, but becomes dominated by $B_0$ within $\lesssim 10$ yr after launch. As time progresses, $(\beta)$ gradually increases but never rises above unity (at least for $t < 100$ yr), even for $\beta_i \gg 1$. Still, one might speculate from Figure 2(a) that with sufficient time, $(\beta) \rightarrow 1$ regardless of $\beta_i$.

### 4.1. Proper Motion

For $t \gtrsim 10$ yr, the velocity of the tip of the jet, $v_{\text{jet}}$, attains a steady value and, from Figure 2(b) and Table 2, we find $v_{\text{jet}} \propto B_0^{0.44\pm0.01}$. To understand this result physically, we begin with the magnetic forces:

$$F_\parallel = -\frac{B_y}{r} \nabla_\parallel (r B_\phi);$$

$$F_\phi = \frac{B_y}{r} \nabla_\phi (r B_\phi);$$

$$F_\perp = -\frac{B_y}{r} \nabla_\perp (r B_\phi) + J_\phi \nabla_\phi B_\phi$$

(e.g., Ferreira 1997; Zanni et al. 2007), where $\nabla_\parallel$ and $\nabla_\perp$ are the gradients parallel and perpendicular to $\vec{B}_p$, respectively. For a given field line, a stronger $B_\phi$ at its “footpoint” in the disk ($r = r_0$) generates a stronger $B_p$ which leads to stronger gradients in $r B_\phi$ and thus, from Equations (13), greater magnetic forces to accelerate the flow. In practice, we find that most of the acceleration occurs before the fast point (and not the Alfvén point) located at $r = r_f$, where $r_f$ is a weak function of the field strength at the footpoint and thus of $B_i$.

Following Spruit (1996), one can show that as a function of the “fast moment arm” ($\xi \equiv r_f/r_0$), the poloidal velocity at the fast point is

$$v_{p,f} = \sqrt{a_{p,f} v_{K,0}} \left( \frac{\xi^2 + 2}{\xi} - 3 \right)^{1/4} \propto \sqrt{B_i},$$

(14)

since $a_{p,f}$, the poloidal Alfvén speed at the fast point, is roughly proportional to $B_i$, $v_{K,0} = \sqrt{GM_*/R}$ is the Keplerian speed at the footprint of the field line. We note that measured values of $v_{p,f}$ in our simulations vary as $B_i^{0.5}$ and agree with Equation (14) to within 1% so long as the fluid is in approximate steady state.5

After the poloidal force given by Equations (13) decreases to 1% of its maximum value ($\gtrsim$ a few $r_f$), $v_\phi$ still follows a power law in $B_i$ with index 0.52 $\pm$ 0.04 and essentially unchanged from Equation (14). Nearer the head of the jet where steady state is no longer valid, we find $v_{p,i} \propto B_i^{0.45\pm0.02}$ (where the momentum-weighted average is taken across the jet radius), only slightly shallower than Equation (14). Thus, while the conditions in the jet have changed, some memory of the steady-state conditions at $r_f$ persists.

Finally, $v_{\text{jet}}$ (Figure 2(b) and Table 2) is within $\sim 10\%$ of $v_\phi$ near the bow shock and maintains the same power-law dependence on $B_i$. Thus, these jets are essentially ballistic, where the observed jet speed $v_{\text{jet}} \propto B_i^{0.44\pm0.01}$. In short, all measures of jet speed increase with $B_i$, a trend that agrees with Anderson et al. (2005) who find for much less evolved jets, $v_p \propto B_i^{1/3}$.

### 4.2. Toroidal Velocity

Figure 2(b) and Table 2 show $v_\phi$ averaged over time and the jet volume for $z \gtrsim 100$ AU as a function of $B_i$. Like $v_{\text{jet}}$, $v_\phi$ asymptotes to a constant value. The region inside 100 AU is ignored because the torsion Alfvén wave at low $z$ has a non-negligible $v_\phi$, is not part of the jet, and skews our results. By fitting a power law to these data, we find $v_\phi \propto B_i^{0.66\pm0.01}$.

Unlike $v_{\text{jet}}$, we have not uncovered a rationale for this power law, yet it seems plausible that one must exist given the tightness of fit. Eliminating $B_i$ from the power laws for $\langle v_\phi \rangle$ and $v_{p,f}$, we find that $\langle v_\phi \rangle \propto v_{p,f}^{5.0\pm0.04}$.

To render this a useful observational tool, further work is needed to quantify the effects of other initial conditions such as $\xi$ and $\beta_i$ on both the power-law index and the proportionality constant, as well as the effect our simplified

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5 Indeed, all four steady-state functions from Spruit (1996) remain constant in our simulations to within $\lesssim 5\%$ along steady-state field lines, which we take as validation of our numerical methods.
disk model may have on conditions in the jet at observational length scales.

4.3. Jet Radius and Mass Flux

The jet radius, \( r_{\text{jet}} \), is defined by the contact discontinuity (steep temperature gradient in the second panel of Figure 1) between shocked jet and shocked ambient material, which in turn is determined by where the radial jet ram pressure balances all external forces. Since ram pressure increases with \( v_{\text{jet}} \), \( r_{\text{jet}} \) should increase with \( B_i \), as seen in Table 2. Indeed we find that \( r_{\text{jet}} \) varies with \( B_i \) as a reasonable power law though, unlike \( v_{\text{jet}} \) or \( \langle v_{\phi} \rangle \), the power index is not constant and decreases slowly in time, while \( r_{\text{jet}} \) itself increases in time, though at an ever-slowing rate.

The mass flux transported by the jet, \( \dot{M}_{\text{jet}} \), consists of material from both the disk and the atmosphere. Unlike previous simulations where jets are typically evolved long after the leading bow shock has left the grid, no part of any bow shock in our simulations reaches the boundary of the coarsest grid. Thus, each jet continues to entrain material from the atmosphere throughout the simulation at a rate that has a strong dependence on \( B_i \), as seen in Table 2. Indeed we find that \( \dot{M}_{\text{jet}} \) varies with \( B_i \) as a reasonable power law, with the power index decreasing slowly in time. As the atmosphere is depleted, the mass flux contribution from the disk (which, by design, is independent of \( B_i \)) becomes more important and the dependence of \( \dot{M}_{\text{jet}} \) on \( B_i \) diminishes.

5. DISCUSSION

We have presented the first MHD simulations of protostellar jets that start from a well-resolved launching region (\( \Delta x_{\text{min}} = 0.00625 \) AU) and continue well into the observational domain (2500 AU). On the AU scale, each jet shows the characteristic and near steady-state knotty behavior first reported by Ouyed & Podio (1997b), though the origin of our knots is quite different. On the 1000 AU scale, each jet develops into a ballistic, supersonic (8 \( \lesssim M \lesssim 11 \)) outflow led by a magnetically confined “nose-cone” (Clarke et al. 1986) and a narrow bow shock, consistent with what is normally observed.

On comparing Tables 1 and 2, our simulations comfortably contain virtually all observed protostellar jets on these four important quantities. We note that these tables would not have been in agreement had we stopped the jet at, say, 100 AU and measured these values then. It is only because our jets have evolved over five orders of magnitude in time that we can state with some confidence that the magnetocentricrugal launching mechanism is, by itself, capable of producing jets with the observed proper motion, rotational velocity, radius, and mass outflow rate. Indeed, our jets are still very young, having evolved to only 100 yr, and allowing them to evolve over an additional one or two orders of magnitude in time may still be useful. For example, it would be interesting to know whether (\( \beta_i \)) rises above unity for any of the jets (Figure 2(a)), and thus enter into a hydrodynamically dominated regime. It would also be interesting to see how long it takes for the power laws in jet radius and mass flux as a function of \( B_i \) to reach their asymptotic limits.

Our jet widths tend to be higher than those observed, particularly when one considers that the values for \( r_{\text{jet}} \) in Table 2 are at \( t = 20 \) yr, and that \( r_{\text{jet}} \) continues to grow in time (e.g., for the \( \beta_i = 40 \) jet, \( 2 r_{\text{jet}} \sim 100 \) AU by \( t = 100 \) yr). As our jet radii mark the locations of the contact discontinuity while observed radii mark hot emitting regions, our widths should be considered upper limits. That our values contain all observed jet widths is a success of these simulations.

Similarly, our numerical mass fluxes are higher than observed values by at least an order of magnitude. Since observed mass-loss rates account only for emitting material (e.g., in forbidden lines; Hartigan et al. 1994), and thus temperatures in excess of \( 10^4 \) K (Dyson & Williams 1997, p. 104), our mass fluxes are necessarily upper limits as well. Indeed, if we measure our mass fluxes near the jet tip (instead of at 200 AU for Table 2) and restrict the integration to fluid above \( 10^4 \) K, our mass fluxes drop by a factor of 10–100, in much better agreement with Table 1.

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