NUMERICAL ANALYSIS OF JETS PRODUCED BY INTENSE LASER

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1. INTRODUCTION

Astrophysical jets found in supersonic bipolar flows from black holes, active galactic nuclei (AGNs), protostars, and similar phenomena are subject to active research. Several problems concerning jets remain to be solved, for example, why the jets propagate such a long distance so stably. Many numerical simulations of astrophysical jets have been performed so far. Two-dimensional nonrelativistic calculations have been done for AGN jets by Norman et al. (1982). They found some dimensionless parameters relevant to the propagation of jets over a long distance. Two- and three-dimensional relativistic calculations have been carried out for protostellar jets by Aloy et al. (1997) and Aloy et al. (1999). Protostellar jets have been studied, for example, by Marti et al. (1997) and Aloy et al. (1999). Protoplastellar jets found in AGN jets by Norman et al. (1982) have been performed so far. Two-dimensional nonrelativistic calculations have been done for AGN jets by Norman et al. (1982). They found some dimensionless parameters relevant to the propagation of jets over a long distance. Two- and three-dimensional relativistic calculations have been carried out by Marti et al. (1997) and Aloy et al. (1999). Protostellar jets have been studied, for example, by Marti et al. (1997) and Aloy et al. (1999).

So far these studies are purely theoretical. Recently, however, with the advance of femtosecond laser technology, an experimental investigation of these astrophysical phenomena in the laboratory seems to be possible (Remington et al. 1999, 2000b). This new scientific field is called laboratory astrophysics or laser astrophysics. This progress enables the investigation of radiation hydrodynamics, hydrodynamic instabilities, and similar astrophysical phenomena in model experiments with intense lasers. Such experiments might be able to give us a new and improved insight into those phenomena which could not be obtained by studies of extraterrestrial observations or numerical simulations. Some experiments along these lines have already been carried out, in which hydrodynamic instabilities and the interactions of a strong shock with matter have been measured. Other experiments have been proposed (Remington et al. 2000a). Backups by numerical simulations are indispensable because they can connect the processes occurring in laser-produced plasmas to astrophysical phenomena that are the aim of our studies. This paper is one of those attempts along this line. In this paper we will propose a different experiment to understand astrophysical jets.
expansion velocity of the plasma reaches a few hundreds of kilometers per second. This is of the same order as the velocity of protostellar jets. This indicates that jets which realistically mimic astrophysical ones can be produced in the laboratory. In these experiments, however, the plasma expands into the vacuum, and one cannot investigate the influence of ambient gas, which does exist for the astrophysical jets. As a matter of fact, the ratio of the mass density of the jet to that of the interstellar gas \( \rho_{\text{jet}}/\rho_{\text{ambient}} \) is 1–10 for protostellar jets.

In this paper we investigate the propagation of laser-produced jets not only in the vacuum but also in the ambient gas using the two-dimensional hydrodynamics code developed recently (Mizuta 2001). Although the code has two versions, relativistic and nonrelativistic ones, in the present paper we use only the nonrelativistic one.

This paper is organized as follows. In §§ 2 and 3 the basic formulas, the numerical method, and the initial conditions are presented. The main results are shown in § 4. We discuss some implications for astrophysical jets in § 5 and give a short summary in § 6.

2. BASIC EQUATIONS AND NUMERICAL METHOD

Assuming a plasma having axial symmetry, the two-dimensional Euler equation for a perfect fluid is solved:

\[ \frac{\partial \mathbf{u}}{\partial t} + \frac{1}{r} \frac{\partial (r f)}{\partial r} + \frac{\partial \mathbf{g}}{\partial z} = \mathbf{s}, \]

where the conserved variable vector \( \mathbf{u} \), the flux vectors \( f \) and \( g \), and the source vector \( s \) are defined as

\[ \mathbf{u} = (\rho, \rho v_r, \rho v_z, E)^T, \]

\[ f = [\rho v_r, \rho v_r^2 + p, \rho v_r v_z, (E + p) v_r]^T, \]

\[ g = [\rho v_z, \rho v_r v_z, \rho v_z^2 + p, (E + p) v_z]^T, \]

\[ s = \left( \frac{\rho}{r}, 0, 0, 0 \right)^T, \]

\[ E = \rho \left( \epsilon + \frac{v_r^2 + v_z^2}{2} \right). \]

In the above equations, \( \rho, v_r, p, \) and \( \epsilon \) are mass density, \( r \) component of velocity, pressure, and specific internal energy, respectively. We ignore viscosity and heat conduction of the fluid. These assumptions are justified as follows. As for the collisions, the parameter \( \delta = \lambda_{\text{mfp}}/L_{\text{sys}} \), defined as the ratio of the collisional mean free path \( \lambda_{\text{mfp}} \) of electrons or ions and the characteristic length of the system \( L_{\text{sys}} \), is \( 10^{-4} \) to \( 10^{-6} \) for our models described below. Hence, a hydrodynamic description of the jets is appropriate. Other dimensionless parameters characterizing the dissipative processes are the Reynolds number \( Re \equiv r_j v_j/v \), which is a measure of the local viscosity, and the Peclet number \( Pe \equiv r_j v_j/v_{\text{th}} \), which describes the heat conduction inside the jet. Here \( r_j, v_j, v, \) and \( v_{\text{th}} \) are the radius and velocity of the jet and the viscosity and thermal diffusivity coefficients, respectively. Criteria for the influence of these dissipative phenomena on the flow are \( Re \lesssim 1 \) and \( Pe \lesssim 1 \). In our models we find \( Re \gtrsim 10^9 \) and \( Pe \gtrsim 10^7 \). Thus, the neglect of these dissipative processes is well justified. These numbers are evaluated at \( t = 2 \) ns at the center of the inflow region (see below).

An ideal gas equation is used in the computations,

\[ p = \rho \epsilon (\gamma - 1), \]

where \( \gamma \) is the adiabatic exponent (= 5/3 in this paper).

We adopt Marquina’s flux formula (Donat & Marquina 1996), which is based on an approximate Riemann solver derived from the spectral decomposition of the Jacobian matrix of Euler equations. We have developed a computational code which solves equations (1) and (7). The code uses cylindrical coordinates with 300(\( r \)) \times 1500(\( z \)) grid points. The accuracy of the code is the second order in space due to an MUSCL method (van Leer 1977, 1979) and is the first order in time.

It is assumed that the plasma is optically thin, so the radiation cooling effect is important. In the cooling term \( |J_{\text{rad}}| \), only the bremsstrahlung is included (Zeldovich & Raizer 1966),

\[ |J_{\text{rad}}| = 1.42 \times 10^{-27} Z^2 T^{1/2} n_e n_{\text{ion}} \text{ergs cm}^{-3} \text{s}^{-1}, \]

where \( Z^* \), \( T, n_e, \) and \( n_{\text{ion}} \) are the average ionic charge, temperature, and number densities of electron and ion, respectively. The temperature and average ionic charge are related to the specific internal energy \( \epsilon \) by

\[ \epsilon = \frac{3}{2} \frac{(Z^* + 1) T}{m_i}, \]

where \( m_i \) is the mass of ion. The average ionic charge state \( Z^* \) and temperature \( T \) were calculated iteratively from local instantaneous density and specific internal energy, as obtained from the hydrodynamic calculation by using equation (9) and an approximate expression based on the Thomas-Fermi model for \( Z^* \) (Salzmann 1998). The formulas for the line emission processes were computed in the usual manner and are not included here for simplicity. As already mentioned, the radiation transport is also neglected in this paper. Since we are interested mainly in the qualitative effects of radiation cooling, these approximations are adequate.

We have to distinguish the target matter from the ambient gas in order to make the radiation cooling effective only for the former. For this purpose we introduced another continuity equation for the target matter:

\[ \frac{\partial (f)}{\partial t} + \frac{1}{r} \frac{\partial (r f)}{\partial r} + \frac{\partial (\rho f)}{\partial z} = 0, \]

where \( f \) is the fraction function. This equation is solved simultaneously with the hydrodynamic equations. With \( Z^* \),
Fig. 2.—Initial conditions. The cone target is set at the left end of the computational region (1200 µm × 3000 µm). Hot plasma is put at the surface. Its depth is ∼10 µm.

In order to simulate the experiments, we put a "cone" target at an end of the computational region. At \( t = 0 \), the laser deposits all its energy on the surface. The depth of the surface is ∼10 µm. A Gaussian-shaped laser pulse of duration 100 ps was assumed. This is much shorter than the dynamical timescale, which is more than a few nanoseconds. When the target made of gold is irradiated by an intensity 

\[ I \approx 10^{14} \text{ W cm}^{-2} \neq 0 \], the temperature of the target rises to a few hundred eV up to 1 keV, and the ablation takes place. Because the velocity of the ablated plasma is at most its sound velocity, \( v_{ab} \approx c_s \sim 10^7 \text{ cm s}^{-1} \), the ablated plasma expands only up to ∼10 µm during the laser irradiation ∼100 ps. This is much smaller than the scale of the system we consider in this paper. The laser energy deposited into thermal energy is \( E_L = 526 \text{ J} \) for all simulations. Energy of ionization is not included; therefore, in a real experiment the incident laser energy, of course, has to be significantly larger.

Four simulations have been done in this paper. In cases 1 and 2, the ambient gas is very dilute (\( \rho_a = 10^{-6} \text{ g cm}^{-3} \)). In the first case we neglect the radiation cooling, while in the second case the radiation cooling is turned on. These intend to simulate the experiments. In cases 3 and 4, we allow the plasma expansion in the presence of an ambient gas. Case 3 neglects the radiation cooling, and case 4 is done with the radiation cooling term. No corresponding experiments have been carried out as yet. These models serve to clarify the effect of the ambient gas.

The temperature of the target surface, where the laser energy is deposited, increases to ∼400 eV. The angle of the cone is 126°, which is similar to the experiments (see Fig. 2). The target matter is chosen to be gold (\(^{79}\text{Au}\)) so that we can obtain a large effect of the radiation cooling due to a high-Z plasma. The density of the target \( \rho_t \) is initially uniform, \( \rho_t = 19.2 \text{ g cm}^{-3} \), and equals the solid density of the gold. Table 1 summarizes the parameters of the initial conditions.

### 4. RESULTS

#### 4.1. Vacuum Case

First, we discuss cases 1 and 2 which simulate plasma expansion into the vacuum. This condition is similar to the experiments of Farley et al. (1999) and Shigemori et al. (2000). For numerical reasons, we assumed in these cases as well a cold (\( T_e = 0.026 \text{ eV} \)) and very low density (\( \rho_a = 10^{-6} \text{ g cm}^{-3} \)) ambient gas. As a result of this very low density, this ambient gas had no influence on any of our results. Figure 3 shows the density contours at different times with the radiation cooling off (Fig. 3a) and on (Fig. 3b), respectively.

In the case of no radiation cooling (Fig. 3a), we can see two regions, that is, the inflow region in which matter flows to the axis and the outflow region. This is schematically shown in Figure 4. Most of the laser-heated matter expands from the target, forming plasma with velocity of a few hundreds of kilometers per second; farther away from the target it converges to the symmetry axis, making a nozzle-like structure in the inflow region. Then, the plasma in the nozzle spouts out from the tip of the nozzle, forming the outflow region. Because the ambient gas is very dilute, there is no pressure support to sustain this nozzle structure. As a result, the outflow spreads out in all directions. This flow structure is suitable for low-Z targets because the effect of the radiation cooling is indeed negligible.

An analysis of the nozzle structure in the inflow region is in order. The converging flow from the cone target to the central axis is very similar to the structure considered by Cantó, Tenorino-Tagle, & Różycka (1988). From their theoretical model we can estimate the beam radius \( r_f \) and...
velocity $v_j$ from the converging flow angle $\theta$, the reciprocal of the compression ratio $\xi$ between the density of the nozzle to that of the inflow and the velocity of the converging flow $v_0$ (see Fig. 5). These are $\theta = 26.5^\circ$ and $\xi = \frac{1}{3}$ in our case. Then $r_j$ and $v_j$ are determined as

$$r_j = \frac{\tan \alpha}{\tan \theta + \tan \alpha} y_0,$$  \hspace{1cm} (11) \\
$$v_j = \frac{\cos (\theta + \alpha)}{\cos \alpha} v_0.$$  \hspace{1cm} (12)

where $y_0$ is the width of the inflow and $\alpha$ is the angle between the conical shock and $z$-axis, given as

$$\tan \alpha = \frac{(1 - \xi) + [(1 - \xi)^2 - 4\xi \tan^2 \theta]^{1/2}}{2 \tan \theta}. \hspace{1cm} (13)$$

Applying this theory to our case with $y_0 = 250 \mu m$ and $v_0 \sim 3 \times 10^7 \text{ km s}^{-1}$, we obtain $r_j = 180 \mu m$ and $v_j \sim 1 \times 10^7 \text{ km s}^{-1}$, in good agreement with our numerical results.

For high-Z targets, we have to consider radiation cooling effects. When radiation cooling is taken into account, the
flow structure is changed dramatically (Fig. 3b). The plasma in the inflow region is collimated strongly and shows a very thin nozzle at a later time (the radius of the structure is \( \sim 40 \mu m \)). Although the collimation of the plasma is sustained in the inflow region by converging inflow from the target, the matter again spreads out in the outflow region, resulting in almost the same structure as that for the case without radiation cooling. These experiments are not suitable for the study of the propagation of collimated jets.

Figures 6a, 6b, and 6c are the density, pressure, and temperature profiles along the \( r \)-axis, respectively, at \( z = 1000 \mu m, t = 2.0 \) ns. In the case with radiation cooling, the plasma is cooled efficiently around the symmetry axis, leading to the increase of the pressure gradient. As the density increases, the plasma is further cooled, and as a result a jetlike nozzle is formed, which was actually observed in the experiments. Altogether, our simulations support the experimental results that the radiation cooling is a crucial ingredient for the formation of this nozzle.

4.2. Dense Gas Case

In cases 3 and 4, the ambient gas density was much higher, \( \rho_g = 10^{-3} \) g cm\(^{-3}\). The time evolution of hydrodynamics leading to the formation of the inflow and outflow regions is essentially unchanged. Figure 7 shows the density contours at different times with the radiation cooling off (Fig. 7a) and on (Fig. 7b), respectively. The bow shock, which accelerates the ambient gas, can be seen clearly.

In the case with radiation cooling, the flow structure is dramatically changed as is obvious in Figure 7b. The effect of radiation cooling manifests itself in both the inflow and the outflow regions. In the inflow region, the nozzle structure becomes very thin just like in the vacuum case in the later time. The main difference appears in the outflow region. Unlike all the other cases, the width of the outflow region is very small. Opposed to intuition, this narrow structure is not a direct outcome of the thin nozzle in the inflow region. This is understood from the second panel of Figure 7b, in which there is already a very thin outflow region while there is no collimated structure in the inflow region. When the matter enters the outflow region, it tends to spread out as in the vacuum case. However, the shocked ambient gas slows the expansion in the direction perpendicular to the symmetry axis, and the density and temper-
Fig. 7.—Density profiles in the r-z plane, in the dense gas case. (a) Radiation off (case 3). (b) Radiation on (case 4). Time $t = 0.0, 1.0, 2.0, 3.0$, and 4.0 ns.

The density profiles show the evolution of the plasma structure in the outflow region. The density ratio is $\eta = \rho_{\text{jet}}/\rho_0 \sim 100$, namely, this is a so-called dense jet. From Figure 8b, we see that the narrow structure in the outflow region is sustained by the shocked ambient gas, whose pressure is about 1000 times larger than that of the unshocked ambient matter. The jet is supersonic, $M_{\text{jet}} \sim 70$. The velocity of the jet is $\sim 400 \text{ km s}^{-1}$. This is of the same order as the velocity of protostellar jets. In order to study the propagation of astrophysical jets in the laboratory, we...
think that it is necessary to do experiments with ambient gas.

5. DISCUSSION

Firstly, we consider the scaling law between a laboratory jet and a protostellar jet. The velocity of these two jets is almost the same: a few hundred kilometers per second. The ratio of the timescales is, however, widely different: $10^{18} - 10^{20}$. For the experiment to simulate efficiently the astrophysical jet, the ratio of the lengths of jets has to be $\sim 10^{18} - 10^{20}$. The typical ratio of the radii of the jets ($r_{\text{prot}}/r_{\text{lab}}$) is $10^{15} \text{cm}/10^{-3} \text{cm} = 10^{18}$. Thus, the aspect ratios of the jets are close to each other. The temperature is a few eV for both jets, which is consistent with the scaling of velocity mentioned above. Thus, a good scaling law holds between the protostellar jet and the laboratory jet considered in this paper.

The relevant parameters which characterize the jet are the density ratio $\eta$, the Mach number and velocity of jets $M_{\text{jet}}$ and $v_{\text{jet}}$, the pressure ratio $K$, and the cooling parameter $\chi$. As the magnetic fields were not included in our code, no comparison can be made with MHD simulations (Todt et al. 1993; Cerqueira & de Gouveia Dal Pino 1997; Frank et al. 1998, 2000; Gardiner et al. 2000). As is obvious from the results obtained for both the vacuum and the ambient gas cases, the radiation cooling is the most crucial ingredient in the collimation of the jets.

The cooling parameter is defined from the cooling term which has already been employed in the hydrodynamic simulations, $|J_{\text{rad}}|$ in equation (8). It is the ratio

$$\chi = \frac{v_{\text{jet}} \tau_{\text{rad}}}{r_{\text{jet}}}$$

where $v_{\text{jet}}$ and $r_{\text{jet}}$ are the velocity and radius of the jet, respectively, and $\tau_{\text{rad}}$ is a characteristic radiation cooling time; $\tau_{\text{rad}}$ is defined from the ratio of thermal energy density to emission power as

$$\tau_{\text{rad}} = \frac{e_{\text{thr}}}{|J_{\text{rad}}|}.$$  

Fig. 8.—Density, pressure, and Mach number profiles along the $r$-axis at $z = 2000 \mu\text{m}$, $t = 4.0$ ns. The radiation cooling term is off (dashed line; case 3) and on (solid line; case 4).

The smaller $\chi$ is, the more effective the cooling is. The cooling term adopted in equation (8) is different from those used in other papers (Blondin et al. 1990; Stone & Norman 1993). We assumed in this paper that the jet matter is Au, rather than hydrogen in most other papers. Our choice was influenced by the need to enhance the radiation cooling so that the hydrodynamics is affected within a few nanoseconds. In fact, the average ionic charge state $Z^+$ is about 20–40 for $T = 50–100$ eV. Although the detailed dependence on $T$ is different between our cooling term and others, it is emphasized that the dynamics should be similar as long as the cooling parameter $\chi$ has a similar value during the greater part of the time evolution.

In case 4 (with dense ambient gas and radiation cooling), the typical radiation cooling time is $\tau_{\text{rad}} \sim 10$ ns for $t = 1.0$ ns and $\tau_{\text{rad}} \sim 100$ ns for $t = 4.0$ ns. These were evaluated at the points with the largest emission power. Accordingly, the radiation cooling parameters for each time are $\chi \sim 10$ and 100, respectively. Thus, the radiation cooling is effective for the earlier phase, while the jet is cool enough, while the radiation cooling is ineffective for the later phase. However, it is noted that we underestimate the radiation cooling, since we neglected the line emissions which could be substantial in our cases, and take into account only the bremsstrahlung. This was done mainly to keep numerical simplicity.

The aim of our calculations and future experiments is to understand better the physics of protostellar jets. Although the observational values of physical quantities are rather uncertain, they are typically $\eta \sim 10^3 \text{ cm}^{-3}$, $T_j \sim 10^4 \text{ K}$, $v_j \sim 100 \text{ km s}^{-1}$, and $r_j = 10^{15} \text{ cm}$ (Bally & Reipurth 2001). These values roughly correspond to $\eta \approx 1\times10^3$, $M_{\text{jet}} \approx 1$ with the main uncertainty coming from the lack of data on the density and temperature of the ambient gas. From these parameters, it is estimated that the total emission power is $|J_{\text{rad}}| \sim 10^{-18} \text{ ergs cm}^{-3} \text{ s}^{-1}$ for collisions by excited hydrogen atoms (Dalgarno & McCauley 1972). The emission power due to the bremsstrahlung is $|J_{\text{rad}}| \sim 10^{-19} \text{ ergs cm}^{-3} \text{ s}^{-1}$. Thus, the bremsstrahlung is substantially weaker, and in general the bound-bound emission is dominant in such plasmas. The emission power $|J_{\text{rad}}| \sim$
$10^{-18}$ ergs cm$^{-3}$ s$^{-1}$ gives a cooling time of $\tau_{\text{rad}} \sim 10^9$ s, which then leads to a cooling parameter of $\chi \sim 10$. It is evident from these values that the experimental values used in this paper are appropriate for studying astrophysical jets.

As mentioned in § 1, Blondin et al. (1990) and Stone & Norman (1993) studied the properties of propagating jets by varying $\eta$ between 1 and 10. Their parameters correspond to $\chi$ between 0.1 and 10. We think that our results have an overlapping region with the initial conditions of those papers. This means that the experiments proposed here might give some new insights into the astrophysical jets studied theoretically in their papers.

6. CONCLUSION

In this paper results of four hydrodynamic simulations have been shown, in which jets were generated from a "cone" target irradiated by an intense laser. In all the simulations we have found a common feature, that is, the existence of an inflow region and an outflow region. Depending on whether the radiation cooling is turned on or off, and whether there is some ambient gas or not, the flow structures inside these regions are very different from each other.

When the plasma flows into a very low density ambient gas, the collimation of the plasma jet occurs only as a result of radiation cooling. With radiation cooling, the collimated plasma shows up in the inflow region. However, since the structure is sustained by the converging inflow from the target, it disappears in the outflow region, in which there is no pressure support to maintain it. This feature is very similar to that found in the experiments. We have also found that the nozzle structure with no radiation cooling is well described with the analytical model by Cantó et al. (1988).

In the case with ambient gas, a bow shock appears very clearly. The main difference between the vacuum and dense gas cases is that in the latter case the collimated flow appears not only in the inflow region but also in the outflow region. Particularly, the collimated plasma in the outflow region is sustained by the pressure of the shocked ambient gas.

We have shown the possibility of carrying out a laboratory experiment, using high-intensity laser-plasma interaction, to generate collimated plasma jets propagating into an ambient gas. The parameters of such a plasma jet are similar to those of protostellar jets. If such an experiment is performed, we can get important information about the physics of astrophysical jets.

External magnetic fields were not included in our calculations. Experiments with magnetic fields have already been proposed and will be carried out in the future. We intend to adjust our simulations to the needs of future experiments.

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