Adiabatic–radiative shock systems in YSO jets and novae outflows

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ABSTRACT

Context. The termination regions of non-relativistic jets in protostars and supersonic outflows in classical novae are non-thermal emitters. This has been confirmed by radio and gamma-ray detection, respectively. A two-shock system is expected to be formed in the termination region where the jet, or the outflow material, and the ambient medium impact. Radiative shocks are expected to form in these systems given their high densities. However, in the presence of high velocities, the formation of adiabatic shocks is also possible. A case of interest is when the two types of shocks occur simultaneously. Adiabatic shocks are more efficient at particle acceleration while radiative shocks strongly compress the gas. Furthermore, a combined adiabatic–radiative shock system is very prone to developing instabilities in the contact discontinuity, leading to mixing, turbulence, and density enhancement. Additionally, these dense non-thermal jets and outflows are excellent candidates for laboratory experiments as demonstrated by magnetohydrodynamics scaling.

Aims. We aim to study the combination of adiabatic and radiative shocks in protostellar jets and novae outflows. We focus on determining the conditions under which this combination is feasible together with its physical implications.

Methods. We performed an analytical study of the shocks in both types of sources for a set of parameters by comparing cooling times and propagation velocities. We also estimated the timescales for the growth of instabilities in the contact discontinuity separating both shocks. We studied the hydrodynamical evolution of a jet colliding with an ambient medium with 2D numerical simulations, confirming our initial theoretical estimates.

Results. We show that for a wide set of observationally constrained parameters, the combination of an adiabatic and a radiative shock is possible at the working surface of the termination region in jets from young stars and novae outflows. We find that instabilities are developed at the contact discontinuity, mixing the shocked materials. Additionally, we explore the magnetohydrodynamic parameter scaling required for studying protostellar jets and novae outflows using laboratory experiments on laser facilities.

Conclusions. The coexistence of an adiabatic and a radiative shock is expected at the termination region of protostellar jets and novae outflows. This scenario is very promising for particle acceleration and gamma-ray emission. The parameters for scaled laboratory experiments are very much in line with plasma conditions achievable in currently operating high-power laser facilities. This provides a new means for studying novae outflows that has never been considered before.

Key words. shock waves – ISM: jets and outflows – stars: jets – novae, cataclysmic variables – instabilities – radiation mechanisms: non-thermal

1. Introduction

Astrophysical jets and outflows from parsec to megaparsec scales are very common in the Universe (e.g. Livio 1999; de Gouveia Dal Pino 2005). Shocks in non-relativistic jets and outflows such as those in protostars and novae are expected to be radiative given their large densities. Radiative shocks are not efficient particle accelerators, but there is evidence of non-thermal emission. Synchrotron radio emission has been detected in both protostellar jets (e.g. Carrasco-González et al. 2010; Feeney-Johansson et al. 2019) and novae (e.g. Vlasov et al. 2016), whereas gamma-rays have only been detected in the latter (Ackermann et al. 2014a). Accelerating synchrotron radio-emitting electrons is possible for almost any kind of shock. However, accelerating TeV particles for gamma-ray emission is not trivial for radiative shocks. Steinberg & Metzger (2018) showed that a working surface composed of two radiative shocks can accelerate ions with an efficiency of ~ 0.01 when non-linear thin shell instabilities take place at the contact discontinuity. In contrast, non-relativistic adiabatic shocks have a ten times greater efficiency (Caprioli & Spitkovsky 2014).

We are interested in the combination of adiabatic and radiative shocks in the termination region of non-relativistic jets and outflows. The self-similar dynamics of this configuration of shocks was studied for the first time by Gintrand et al. (2021). The advantage of this configuration of shocks is that, whereas particles are accelerated in the adiabatic shock, the downstream region of the radiative shock acts as a dense target for gamma-ray and neutrino emission by inelastic proton–proton (pp) collisions and via relativistic Bremsstrahlung. In addition, the radiation field from the downstream region of the radiative shock can ionise the plasma of the jet or outflow, increasing the efficiency of particle acceleration from the adiabatic shock.
termination region is prone to the growth of hydrodynamic instabilities. It is now possible to study some of these instabilities through high-energy-density laboratory experiments; for example Rayleigh-Taylor (Kuranza et al. 2018), Kevin-Helmholtz (Doss et al. 2015), and thermal instabilities (Suzuki-Vidal et al. 2015).

In this work, we focus on protostellar jets and classical novae outflows. For simplicity, we refer to both outflows as jets hereafter. In both cases the initial jet flow moves with velocities \( v_j \sim 100 - 1000 \text{ km s}^{-1} \), but in ambient media with different densities (see Table 1). We show that for certain combinations of jet velocities and ambient \((n_a, n_j)\) and jet \((n_j, n_j)\) densities, the working surface in the termination region is composed of an adiabatic and a radiative shock. We also show that this combination leads to fast growth of hydrodynamical instabilities and therefore a significant level of mixing. This situation is very promising for gamma-ray emission through pp inelastic collisions and relativistic Bremsstrahlung.

The paper is organised as follows: in Sect. 2 we describe the properties of the shocks in the termination region. In Sect. 3 we perform a stability study of the leading working surface. In Sect. 4 we perform numerical simulations with the freely distributed code PLUTO. In Sect. 5 we discuss the results and implications for the expected gamma-ray emission. In Sect. 6 we discuss magnetohydrodynamical (MHD) scaling of the jets of young stellar objects (YSOs) and nova outflows to laboratory experiments. Finally, in Sect. 7 we present the summary and conclusions of this work.

### 2. Shocks in the termination region

Protostellar jets and novae outflows are supersonic and therefore their termination region is made of a bow shock (or forward shock) in the ambient medium and a Mach disc (or reverse shock) moving into the initial jet flow (Fig. 1). In the unidimensional flow approximation, the bow shock moves into the ambient medium at \( v_{bs} \sim v_j/(1 + 1/\sqrt{\chi}) \), where \( \chi \equiv n_j/n_a \) is the jet-to-ambient-density contrast (e.g. Raga et al. 1998; Hartigan 1989). The reverse shock in the jet moves at \( v_{rs} = v_j - 3v_{bs}/4 \). In ‘heavy’ jets (\( \chi > 1 \)), the bow shock is faster than the Mach disc, whereas in ‘light’ jets (\( \chi < 1 \)), the reverse shock is faster than the bow shock. In particular, \( v_{bs} < v_j \) and \( v_{bs} < v_j \sqrt{\chi} \) when \( \chi \ll 1 \), whereas \( v_{bs} > v_j \) when \( \chi \gg 1 \).

The jet density in the termination region can be calculated from conservation of mass as (Rodríguez-Kamenetzky et al. 2017)

\[
\frac{n_j}{cm^3} \approx 150 \left( \frac{M_i}{10^{-6} M_\odot \text{ yr}^{-1}} \right) \left( \frac{v_j}{1000 \text{ km s}^{-1}} \right)^{-1} \left( \frac{R_j}{10^{16} \text{ cm}} \right)^{-2},
\]

where \( M_i \) is the ionised mass loss rate and \( R_j \) is the radius of the section of the jet in the termination region.

In the strong shock approximation, the plasma is compressed by a factor of four and the temperature is \( T_{ps} \sim 2 \times 10^5 (v_{bs}/100 \text{ km s}^{-1})^2 \text{ K} \) immediately after the shock front. Establishing the nature of the shocks, whether they are adiabatic or radiative, can be done by comparing the advection (escape) timescale \( t_{esc} \sim R_j/(v_{bs}/4) \) to the cooling timescale

\[
t_{cool} = \frac{(3/2)k_B T}{n_a \Lambda(T)},
\]

where \( \Lambda(T) \) is the cooling function, which depends strongly on the temperature \( T \). For the case of novae, which are characterised by high shock velocities and thus high temperatures, we use free-free emission which is given by \( \Lambda(T_{ps}) = 2 \times 10^{-27} T_{ps}^{1/2} \text{ erg cm}^{-3} \text{ s}^{-1} \). Shocks in YSO jets have lower velocities and therefore metal-line cooling dominates (e.g. Sutherland & Dopita 1993).

Equivalently, we can compare the (thermal) cooling length \( l_{cool} = t_{cool} v_{bs}/4 \) and \( R_j \). The condition \( l_{cool} > R_j \) (Heathcote et al. 1998) for a YSO shock to be adiabatic can be rewritten as \( v_{sh} > v_{sh,ad} \), where

\[
v_{sh,ad} = 650 \left( \frac{n_a}{10^3 \text{ cm}^{-3}} \right) \left( \frac{R_j}{10^{16} \text{ cm}} \right)^{1/2}. \tag{3}
\]

In order to study the nature of the two shocks in the termination region, i.e. whether they are adiabatic or radiative, we sample the parameter space for \( n_j \) and \( n_a \) for the cases of YSO jets and novae outflows presented in Table 1.

#### 2.1. Young stellar objects

Stars are formed within dense molecular clouds, accreting matter onto the central protostar with the formation of a circumstellar disc and bipolar jets. These ejections are collimated flows of disc and stellar matter accelerated by magnetic field lines (Blandford & Payne 1982; Shu et al. 1994), and moving with speeds of \( v_j \sim 100 - 1000 \text{ km s}^{-1} \) into the ambient molecular cloud. Molecular matter from the cloud is entrained by the jet, forming molecular outflows.

The ionised mass-loss rate in YSO jets is \( 10^{-8} \leq M_i \leq 10^{-5} \text{ M}_\odot \text{ yr}^{-1} \), and the width of the jet in the termination region is characteristically \( R_j \sim 10^{16} \text{ cm} \). By inserting these values into Eq. (1), we find \( n_j \) in the range \( 1 - 10^3 \text{ cm}^{-3} \). The ambient medium is the molecular cloud where the protostar is embedded, and typical values of \( n_a \) are in the range \( 10 - 10^3 \text{ cm}^{-3} \). This results in values \( 10^{-5} \leq \chi \leq 10 \). We consider \( v_j = 300 \) and \( 1000 \text{ km s}^{-1} \) and use the condition given by Eq. (3) to classify the shocks. The results are shown in Figure 2.

Figure 2 indicates when the shocks are radiative or adiabatic, the left plot corresponds to \( v_j = 300 \text{ km s}^{-1} \) and the right plot to a faster jet with \( v_j = 1000 \text{ km s}^{-1} \). The particular case for \( n_a = 500 \text{ cm}^{-3} \) and \( n_j = 5 \text{ cm}^{-3} \) for \( v_j = 300 \text{ km s}^{-1} \),

![Fig. 1. Schematic diagram showing the formation of forward and reverse shocks in the termination region of YSO jets and nova outflows interacting with an ambient medium. The arrow indicates the fluid velocity direction in the laboratory frame.](image_url)
marked with a white dot in the figure, is investigated in detail with 2D numerical simulations in Section 4. For faster YSO jets, \( v_j = 1000 \text{ km s}^{-1} \), the green region is larger, and in the range of densities studied here there is no region where both shocks are radiative (no pink region). We can conclude that for a large range of parameters for the ambient medium and the jet, we can expect the formation of the adiabatic–radiative shocks that are of particular interest in this work.

2.2. Classical novae

In order to study the shocks produced in classic novae, we follow the model of Metzger et al. (2015); see also Metzger et al. (2016). The shocks in novae outflows are formed from the collision of a slow shell ejection with velocity \( V_{ej} \sim 1000 \text{ km s}^{-1} \), produced in the thermonuclear runaway, and a faster outflow or continuous wind of velocity \( V_w \sim 2V_{ej} \) that follows the ejecta within a few days. The collision, as in the case of YSO jets, produces a system of two shocks: forward and reverse.

The forward shock propagates through the slow shell and the reverse shock moves back through the wind. The shock velocities depend on the density ratio of the colliding media. The expressions are the same as the ones presented at the beginning of this section with \( v_j = V_w - V_{ej} \). The slow shell density is \( n_{ej} \approx M_{ej}/(4\pi R_{ej}^2 f_{\Delta M} m_p) \), where \( M_{ej} \) is the mass of the ejecta, \( R_{ej} = V_{ej} t_{ej} \) its expansion radius, and \( f_{\Delta M} \sim 0.5 \) is a filling factor related to the geometry of the model, i.e. the fraction of the total solid-angle subtended by the outflow. By considering typical values for the slow ejecta, we obtain

\[
\frac{n_{ej}}{\text{cm}^{-3}} \approx 10^{10} \left( \frac{f_{\Delta M}}{0.5} \right)^{-1} \left( \frac{M_{ej}}{10^{-4} M_\odot} \right) \left( \frac{V_{ej}}{1000 \text{ km s}^{-1}} \right)^{-3} \left( \frac{t_{ej}}{1 \text{ wk}} \right)^{-3}.
\]

(4)

The typical mass-loss rate of the fast wind is \( \dot{M}_w = 10^{-5} M_\odot \text{ wk}^{-1} \), giving a wind density for an outflow with radius \( R_{ej} \) of

\[
\frac{n_w}{\text{cm}^{-3}} = 2 \times 10^9 \left( \frac{M_{ej}}{10^{-5} M_\odot \text{ wk}^{-1}} \right) \left( \frac{V_{ej}}{1000 \text{ km s}^{-1}} \right)^{-1} \left( \frac{R_{ej}}{6 \times 10^{13} \text{ cm}} \right)^{-2}
\]

(5)

(see Eq. (1)).

We study a parameter space for \( n_{ej} \) and \( n_w \) varying between \( 10^{-5} \leq M_{ej}/M_\odot \leq 10^{-2} \) and \( 10^{-6} \leq M_w/M_\odot \text{ wk}^{-1} \leq 10^{-3} \). We consider \( V_{ej} = 1000 \text{ km s}^{-1} \) and 3000 km s\(^{-1}\). To classify the shocks we compare the cooling time given by Eq. (2) assuming free-free cooling and the typical duration of this phenomena of \( \sim 2 \) weeks. The results are shown in Figure 3.

In the case of a slow ejecta with \( V_{ej} = 1000 \text{ km s}^{-1} \), the forward shock is always radiative, and therefore the nature of the shock system is given by the reverse shock, which is radiative for \( M_{ej} > 10^{-5.6} M_\odot \) (pink region on the left plot). The case of interest in this work, an adiabatic reverse shock with a radiative forward shock, corresponds to the smaller green region. In the case of a faster ejecta, \( V_{ej} = 3000 \text{ km s}^{-1} \), we obtain a similar plot to that in the case of YSO jets in Fig. 2. The area of the right plot in Figure 3 is divided into: both adiabatic in the lower left (yellow), both radiative in the upper right (pink), an adiabatic forward shock and a radiative reverse shock in the top left (grey), and the opposite situation in the bottom right (green). In this case, the condition for a reverse shock to be adiabatic with an efficiently cooling forward shock is dominant in the given parameter space.

The very high densities in novae make the shocks very prone to efficient radiative cooling. However, we can conclude that under the model adopted, the possibility of having adiabatic shocks in these systems is plausible, especially for \( V_{ej} > 1000 \text{ km s}^{-1} \).

3. Instabilities at the contact discontinuity

The density of the plasma downstream of the radiative forward (bow) shock is

\[
\frac{n'_w}{n_w} = 233 \left( \frac{v_{fs}}{100 \text{ km s}^{-1}} \right)^2 \left( \frac{T}{10^4 \text{K}} \right)^{-1},
\]

(6)

(e.g. Blondin et al. 1990) when the plasma is cooled down to a temperature \( T \), making the density contrast at the contact discontinuity \( 4n'_w/n_w < 1 \). As a consequence, the contact discontinuity is vulnerable to dynamical and thermal instabilities. A dense layer of density \( n'_w \) located at distance \( l_{fs} \) downstream of the bow shock fragments into several clumps (e.g. Calderón et al. 2020). However, we note that the component of the magnetic field in the ambient medium parallel to the bow shock front \( (B_{\perp}) \) limits the compression factor to a maximum value of

\[
\frac{n'_w}{n_w} \sim 78 \left( \frac{n_w}{10^4 \text{ cm}^{-3}} \right) \left( \frac{v_{fs}}{100 \text{ km s}^{-1}} \right) \left( \frac{B_{\perp}}{0.1 \text{ mG}} \right)^{-1}
\]

(Blondin et al. 1990).

This indicates that significant enhancement in the plasma density downstream of the reverse shock is feasible if instabilities grow quickly enough to fragment the dense shell and form such clumps.

Rayleigh-Taylor instabilities

The Rayleigh-Taylor (RT) instability can grow in the contact discontinuity because of the velocity shear and the force exerted by the downstream material of the reverse shock on the forward shock.

If the forward shock is radiative, the formation of a shell much denser than \( n_{ej} \) and \( n_w \) at the contact discontinuity makes the

Table 1. Typical flow and ambient parameters for YSO jets and novae outflows.

| Parameter                  | YSO          | Novae         |
|----------------------------|--------------|---------------|
| Velocity \( v_j \) [km s\(^{-1}\)] | 100 – 1000   | 1000          |
| Mass loss rate \( \dot{M} \) \[M_\odot \text{ yr}^{-1}\] | \( 10^{-9} \) – \( 10^{-5} \) | \( 10^{-6} \) – \( 10^{-3} \) [M_\odot \text{ wk}^{-1}] |
| Radius \( R_{ej} \) [cm]   | 10\(^{16}\)   | 6 \times 10\(^{13}\) |
| Density \( n_{ej} \) [cm\(^{-3}\)] | 1 – 10\(^{3}\) | 10\(^{9}\) |
| Ambient density \( n_w \) [cm\(^{-3}\)] | \( 10^{-5} \) – 10\(^{-1}\) | 10\(^{11}\) |
working surface unstable even in the case of light jets. Following the analysis in Blondin et al. (1990), the acceleration of the dense shell with a width $W_{sh}$ can be written as $a \sim n_j v_j^2 / n'_j W_{sh}$.

The growth time of RT instabilities is $t_{RT} \sim 1/\sqrt{\kappa}$, where $k$ is the wavenumber. By considering the characteristic dynamical timescale $t_{dyn} = 4R_j/v_j$ we obtain

$$
t_{RT}/t_{dyn} \sim \left(\frac{\sqrt{\kappa}}{4 \sqrt{k} + 1}\right) \left(\frac{\lambda}{R_j}\right)^{1/2} \left(\frac{W_{sh}}{R_j}\right)^{1/2} \left(\frac{T}{10^4K}\right)^{-1/2} \left(\frac{v_j}{1000 \text{km s}^{-1}}\right).
$$

The condition $t_{RT}/t_{dyn} < 1$ leads to

$$
\left(\frac{v_j}{1000 \text{km s}^{-1}}\right) > \left(4 \frac{\sqrt{k} + 1}{\sqrt{\kappa}}\right) \left(\frac{W_{sh}}{R_j/3}\right)^{-1} \left(\frac{T}{10^4K}\right)^{1/2},
$$

where we assume $\lambda = 2\pi/k \sim W_{sh}$.

### 4. Numerical study

We performed 2D hydrodynamic simulations with the PLUTO code (Mignone et al. 2007) to illustrate the physical processes mentioned in the previous sections. We are interested in simulating the physics in the interaction between the incoming material, the shocked incoming material, the contact discontinuity, the target shocked material, and the target material. Therefore, we study the fluid collisions in a 2D rectangular box of size $L_x \times L_y$, with $L_x = 4R_j$ and $L_y = 8R_j$ in a uniform Cartesian grid of resolution $1024 \times 2048$. Boundary conditions are periodic in the $x$-direction (horizontal) and outflow in the $y$-direction (vertical). The fluids are assumed to follow an ideal equation of state with an adiabatic index of $\gamma = 5/3$.

In the presence of cooling, we use the tabulated cooling function, which includes metal-line cooling from Schure et al. (2009); see Section 2. Calculations are performed using a HLLC solver with parabolic reconstruction. The time integration is performed using a Runge-Kutta 2 algorithm controlled by a Courant-Friedrichs-Lewy number of 0.4.

We simulate the case of a protostellar jet ($R_j = 10^{16}$ cm) impinging upon an ambient molecular cloud. The results can be extrapolated to the case of novae. Initially, we have a fluid of density $n_j = 5$ cm$^{-3}$ and velocity $v_j = 300$ km s$^{-1}$ at $T_j = 1000$ K for $y \leq 4R_j$ colliding with a material at rest of density $n_a = 500$ cm$^{-3}$ and temperature at $T_a = 100$ K for $y > 4R_j$ with $n_a/n_j = 100$. According to our analytical estimates in Sect. 2, this configuration should develop a working surface composed of an adiabatic reverse shock and a radiative forward shock. The white dot in Figure 2 shows the position of these parameters in the shock diagnostic map.
4.1. Results

Figure 4 shows the density (top) and temperature (bottom) profile of the system along the $y$–direction, averaged in $x$, for a time $t = 4 \times 10^9$ s, where $f_{\text{phy}} \equiv R_j/10^5 \text{ cm s}^{-1} = 10^{11} \text{ s}$ is the physical time of the simulation. Both shocks are strong, producing a density jump of the order of 4. The jet and ambient shocked gases reach temperatures of $1.8 \times 10^9$ and $1.8 \times 10^9$ K, respectively. Using the post-shock temperature, we measure the shock velocities in the simulations ($V_s = 100 \sqrt{T_s}/1.38 \times 10^5 \text{ km s}^{-1}$), while calculating the displacement of the shock fronts as time evolves. The shock front cells are established with good accuracy by searching for the position of a strong gradient in the temperature profile. The velocity of the contact discontinuity is the fluid velocity of the shocked material, which is directly extracted from the simulation by measuring the velocities of the shocked jet and the shocked ambient material. The measured velocities of the shocks and the contact discontinuity are in very good agreement with the values predicted by the theory. Using the equations in Sect. 2 we find that $V_s = 362.72 \text{ km s}^{-1}$ and $V_{\text{cd}} = 36.27 \text{ km s}^{-1}$. The contact discontinuity is expected to move with a velocity of $V_{\text{cd}} \approx 3V_s/4 = 27.27 \text{ km s}^{-1}$.

In the presence of cooling, the system acts as predicted from theory: the shocked jet material experiences negligible cooling, while the shocked ambient material cools down. From the density (upper panel) and temperature (lower panel) plots in Figure 4 (dashed lines), we see the typical profile of a radiative forward shock. The material behind the forward shock cools, i.e. the temperature drops, and compresses reaching values of $n'_j \gg 4n_j$. Even when the radiative losses do not affect the reverse shock dynamics, a small drop in temperature can be seen, with $T_{\text{cool}}/T_s \sim 0.8$. The cooling of the forward shock changes the contact discontinuity velocity (see the displacement of the contact discontinuity in Figure 4 with respect to the case without cooling). This also slightly slows the reverse shock down, changing the shock Mach number $M$ and hence the jump conditions across the shock (see e.g. Gintrand et al. 2021). The reverse shock velocity measured in the simulation with cooling is $V_{\text{rs}} \sim 324 \text{ km s}^{-1}$, giving $M_{\text{cool}} \sim 87.5$; in the case without cooling $M \sim 98$. The temperature jump condition can be written as $T_s/T_{\text{cool}} \sim 5/16M^2$ for a strong shock with $\gamma = 5/3$. Therefore, $T_{s,\text{cool}}/T_s \approx M_{\text{cool}}^2/M^2 \sim 0.8$. A difference in the reverse post-shock density is not observed and also not expected because for strong shocks the density jump condition is practically independent of the shock Mach number.

The fluid profile propagates without disturbances in this simple configuration, without mixing of the jet and cloud materials. However, in a real situation the system will suffer perturbations produced by local inhomogeneities for example. In order to study the instabilities that might arise in the evolution of the system (see previous section), we consider a sinusoidal interface $y = y_0 + 0.01 \sin(\frac{\pi}{L}x)$, with $y_0 = 4$ and $L = 2$ between the jet and cloud material, which mimics a perturbation. Figure 5 shows a sequence of density maps as time evolves. Instabilities develop in the contact discontinuity as predicted by the theory, which induce mixing and turbulence. The instability also heats the shocked ambient material locally. The density structures in the unstable layer resemble the typical finger structures developed when the RT instability is operating.

In order to analyse the mixing of the two materials, we include a tracer field in the simulation that is advected with the fluid. Initially, the jet and ambient material have tracer values of 1 and 0, respectively. Figure 6 shows the tracer map at $t = 7 \times 10^9$ s zoomed into the forward shock region. The denser shocked ambient material penetrates the shocked jet material, and it is clear that mixing occurs as expected.

**Magnetic field**

In order to illustrate the effects of a magnetic field, we include a field $B_1 = 5 \mu G$ and $B_2 = 5 \mu G + 5 \mu G \hat{J}$ in the MHD simulation setup. This gives a thermal-to-magnetic pressure ratio $\beta$ of 0.7 for the jet, and 3.5 for the ambient medium. We also consider perturbations of a magnetic field, we include a tracer field in the simulation that is advected with the fluid. Initially, the jet and ambient material have tracer values of 1 and 0, respectively. Figure 6 shows the tracer map at $t = 7 \times 10^9$ s zoomed into the forward shock region. The denser shocked ambient material penetrates the shocked jet material, and it is clear that mixing occurs as expected.

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\footnote{1} In a real situation, perturbations are frequent, arising for example from density inhomogeneities, velocity gradients, and irregular geometry among other effects.
Fig. 6. Tracer map at $t = 7 \times 10^9$ s, zoomed at the forward shock region. The tracer indicates the advection of the jet and ambient materials as the system evolves. Initially the value of $1$ is designated to the jet material (blue) and $0$ to the ambient material (white).

Fig. 7. Magnetic field intensity map at $t = 7 \times 10^9$ s. The red arrows show the magnetic field direction in the given grid point.

Fig. 8. Density maps zoomed into the forward shock region at $t = 7 \times 10^9$ s for the HD (top) and MHD (bottom) regimes.

sider the case in which the magnetic field has no components perpendicular to $V_{sh}$ (no $i$-component). The ambient magnetic field component parallel to the shock front is compressed by the forward shock. Without cooling, the magnetic field compression is approximately $4$, the same factor as the density. In the presence of cooling, further compression is expected and the field is amplified by a factor of about eight. The magnetic field map at $t = 7 \times 10^9$ s is shown in Figure 7. The red arrows indicate the magnetic field direction in each cell plotted over the magnetic field intensity.

The presence of a magnetic field inhibits or hinders the development of the instabilities discussed in Sect. 3. The effect of the magnetic field in the unstable layer can be seen in Figure 8, where we show the density map zoomed into the forward shock region at $t = 7 \times 10^9$ s for hydrodynamic (top) and MHD (bottom) scenarios. Although the presence of the magnetic field reduces the development of instabilities and material mixing in the contact discontinuity, this effect is highly dependent on the magnetic field orientation. In the other case considered (only parallel $B$), the density structures developed by the instabilities exhibit a similar shape and level of material mixing as the non-magnetic case. In an actual astrophysical scenario, establishing the direction of the magnetic field is highly difficult and a number of assumptions are required.

4.2. Power spectrum

We compute the power spectrum in the $x$ direction for fixed height $y$ at $t = 7 \times 10^9$ s. Figure 9 shows power spectra for the density (top) and the velocity (bottom) for two values of $y$. These heights correspond to different regions where density perturbations appear due to instabilities (see upper plot of Fig. 8). The mixing coefficient is defined as $D_{mix} = \sqrt{k_{max} V_{sh}}$, where $k_{max}$ is the scale of the fastest growth rate of the unstable mode in the velocity. From Fig. 8 we estimate $k_{max} \sim 30$ and $V_{sh} \sim 8.5$. In physical units, this yields $D_{mix} \sim 2.4 \times 10^{21}$ cm$^2$ s$^{-1}$. We can estimate the mixing time for a given size $L_0$ by computing

$$t_{mix} = \frac{L_0^2}{D_{mix}} \sim 4.2 \times 10^{10} \left( \frac{L_0}{R_j} \right)^2 \left( \frac{D_{mix}}{2.4 \times 10^{21} \text{cm}^2 \text{s}^{-1}} \right)^{-1} \text{s}. \quad (10)$$

In the following section, we discuss the relevance of efficient mixing for enhancement of the gamma-ray emission by the interaction of relativistic particles with matter fields.

5. Gamma-ray emission

A system of two shocks, one radiative and the other adiabatic, can be present in YSO jets and novae outflows, as is demonstrated in Section 2. The fast reverse shock is more efficient for accelerating particles, as the acceleration time for electrons and protons with energy $E_e$ and $E_p$, respectively, is

$$t_{acc} \sim 2.4 \times 10^{7} \left( \frac{E_e,p}{\text{TeV}} \right) \left( \frac{B}{\text{mG}} \right)^{-1} \left( \frac{V_{sh}}{1000 \text{km s}^{-1}} \right)^2 \quad (11)$$

in the Bohm diffusion regime. Furthermore, the adiabatic reverse shock has a luminosity

$$\frac{L_{sh}}{\text{erg s}^{-1}} \sim 3 \times 10^{35} \left( \frac{R_j}{10^{16} \text{cm}} \right)^2 \left( \frac{n}{10^3 \text{cm}^{-3}} \right)^3 \left( \frac{V_{sh}}{1000 \text{km s}^{-1}} \right)^3, \quad (12)$$

which is higher than the forward radiative shock and can power the shock-accelerated population of particles. Also, the lower densities in the jet in the case of YSOs could avoid injection problems produced by ionisation and collision losses (e.g. O’C Drury et al. 1996).

Accelerated electrons and protons are injected into the shock downstream region with a luminosity $L_{e,p} = f L_{sh}$, where $f \sim 0.1$ and $0.005$ for the case of adiabatic and radiative shocks, respectively (Caprioli & Spitkovsky 2014). These particles must reach the contact discontinuity in order to interact with the cooled compressed layer. In the transport of the particles, various physical ingredients play a role. For simplicity, we only consider the
spatial diffusion (energy dependent) and the advection by the large-scale gas velocity (energy independent). We only consider Bohm diffusion in the shock acceleration process, but far from the shock we expect a faster diffusion regime. Assuming a diffusion coefficient of \( D = 10^{25} (E_{\nu, p}/10 \text{ GeV})^{0.5} \) cm\(^2\) s\(^{-1}\) (slower than the typical value in the ISM due to the instabilities) we obtain
\[
\frac{t_{\text{diff}}}{s} \sim 10^7 \left( \frac{R_j}{10^{16} \text{ cm}} \right)^2 \left( \frac{E_{\nu, p}}{10 \text{ GeV}} \right)^{-2}.
\] (13)

Advection downstream of the adiabatic shock over a distance \( R_j \) can be written as
\[
\frac{t_{\text{adv}}}{s} \sim 4 \times 10^8 \left( \frac{R_j}{10^{16} \text{ cm}} \right) \left( \frac{V_{\text{sh}}}{1000 \text{ km s}^{-1}} \right)^{-1}.
\] (14)

We define the residence time of particles in the downstream region as \( T = \min\{t_{\text{adv}}, t_{\text{diff}}\} \).

The instabilities produced in the contact discontinuity (see e.g. Fig. 8) act to facilitate the interaction of particles with the denser, cool material. In addition, the mixing facilitates the transport and the collision of particles with denser material radiating more non-thermal emission via pp collisions (for hadrons) and relativistic Bremsstrahlung (for leptons). The timescale is very similar in both cooling processes. The simplest form of the relativistic Bremsstrahlung and pp cooling time reads
\[
\frac{t_{\text{Brems, pp}}}{s} \sim 2 \times 10^{11} \left( \frac{n}{10^4 \text{ cm}^{-3}} \right)^{-1}.
\] (15)

We note that \( t_{\text{Brems, pp}} \propto n^{-1} \) and therefore cooling through pp inelastic collisions and relativistic Bremsstrahlung becomes very efficient in the cooling layer where the density is significantly larger than \( 4n_j \).

The synchrotron cooling time of electrons in a magnetic field \( B \) is
\[
\frac{t_{\text{syn}}}{s} \sim 4 \times 10^8 \left( \frac{E_e}{\text{TeV}} \right)^{-1} \left( \frac{B}{\text{mG}} \right)^{-2}.
\] (16)

Inverse Compton (IC) scattering of IR and optical photons from the central object with luminosity \( L_* \) and energy density \( U_{\text{ph}} = L_*/(4\pi z^2 c) \) has a characteristic timescale of
\[
\frac{t_{\text{IC}}}{s} \approx 1.6 \times 10^{10} \left( \frac{E_e}{\text{TeV}} \right)^{-1} \left( \frac{L_*}{10^{44} \text{L}_\odot} \right)^{-1} \left( \frac{z}{10^4 \text{cm}} \right)^2 \frac{c}{E_{\nu, p}},
\] (17)

where \( z \) is the distance from the photon source.

5.1. Young stellar objects

Gamma-ray emission from YSO jets has been modelled in some recent papers (e.g. Araudo et al. 2021), but its detection has not been claimed to date. We show here that efficient mixing by RT instabilities in the contact discontinuity can significantly enhance the predicted gamma-ray emission levels, making them detectable in the GeV domain.

By considering \( B = 1 \) mG, \( V_{\text{sh}} \sim 1000 \text{ km s}^{-1} \), and a typical length scale of \( R_j \sim 10^{16} \) cm, we plot in Figure 10 the timescales of the processes mentioned above. We also plot \( t_{\text{max}} \) for comparison only, given that this value was not computed at the steady state. We can see that the acceleration is very efficient, together with the transport of particles by diffusion, giving a maximum energy of both electrons and protons of about 100 GeV. This transport ensures that the accelerated particles in the reverse shock reach the denser regions where materials start to mix and non-thermal emission is enhanced. On the contrary, the transport by advection drag the particles away downstream of the shocked jet material, which is subdominant in this case.

Even though the real scenario might be much more complicated, these timescales indicate the dominating processes. However, we note that if the magnetic field is locally amplified by non-resonant hybrid instabilities, the maximum energy of protons will probably be determined by the available amplification time (Araudo et al. 2021). A detailed model is beyond the scope of the current study, but will be presented in a future work.

5.2. Novae

Gamma-ray emission has been detected from more than a dozen novae in the GeV range with most of the sources being classical novae (see Chomiuk et al. 2020, for a recent review). For example, Fermi and the High Energy Stereoscopic System (H.E.S.S.) recently detected high-energy and very-high-energy gamma-ray emission from the recurrent nova RS Ophiuchi (ATel#14834 and #14857, respectively); see further details below. The gamma-ray emission spans several orders of magnitude among the detected sources (Franckowiak et al. 2018). This difference might arise
simply because the physical parameters are slightly different in each source. For example, we can see from the shock diagnostic maps in Figure 3 that a change of a factor of three in velocity radically changes the possibilities of having adiabatic shocks in the system. Furthermore, changes in metallicity not considered here can modify the cooling function and might change the shock radiative efficiency.

A correlation between the optical and the gamma-ray emission has been claimed (e.g. Li et al. 2017; Aydi et al. 2020). This correlation appears to be strong in some systems but does not behave equally in all detected sources. In the systems where a correlation exists, it is highly probable that the emission is coming from the same spatial region, i.e. a radiative shock. However, this is not in conflict with our claims. Firstly, not all the possible physical parameters in novae result in adiabatic shocks and some sources might have an adiabatic–radiative shock combination. Secondly, even if the optical and gamma emission are coming from the same radiative shock, this does not confirm that the emitting non-thermal particles have been accelerated in that same shock. Indeed, the emission might arise when an underlying population of relativistic particles —accelerated in the reverse shock for example— is enhanced by the strong radiative shock compression (responsible for the optical emission) and even suffers reacceleration (e.g. Blandford & Cowie 1982). An interesting case that supports our findings is that of the nova V959 Mon (e.g. Fujikawa et al. 2012). This source has been detected as a GeV gamma-ray transient by Fermi (Ackermann et al. 2014b). X-ray data analysis indicates that the reverse shock should be non-radiative (see, Nelson et al. 2021).

The hadronic scenario is one of the most favourable for explaining the high-energy radiation (e.g. Li et al. 2017; Martin et al. 2018). In this scenario, for protons to produce a gamma-ray of energy $E$ through proton–proton collisions, they need to have energies of at least ten times $E$, which is not the case for electrons emitting through relativistic Bremsstrahlung. This last case favours the scenario of particles being accelerated at a reverse shock and radiating elsewhere, given that the most energetic particles diffuse more efficiently. A detailed modelling is needed to quantify the viability of the adiabatic–radiative shock scenario and this will be addressed in future works. Below we analyse the case of RS Oph, which was recently detected at gamma-rays.

### The case of RS Oph

RS Oph is a symbiotic recurrent nova system that undergoes thermonuclear outbursts approximately every 20 years (e.g. Anupama 2008). The last detected optical outburst was during August 2021 (vsnet-alert 26131\(^2\)). The binary system, located at $d = 1.6\, \text{kpc}$ (Bode 1987), is composed of a white dwarf and a red giant (RG) companion (Dobrzycka & Kenyon 1994). Here we assume that in this source the shocks are produced in the collision of a fast wind with the dense and slow wind of the RG star (Vaytet et al. 2007, 2011). However, other models for RS Oph exist in the literature. In particular, given that the white dwarf in RS Oph is close to the Chandrasekhar limit, some authors modelled the system similarly to a supernova expanding in a wind medium (e.g. Walder et al. 2008; Booth et al. 2016).

The fast wind has an inferred velocity of $V_w > 6000\, \text{km}\, \text{s}^{-1}$ and a mass-loss rate of $M_w \sim 1.6 \times 10^{-5}\, \text{M}_\odot\, \text{yr}^{-1}$ (Vaytet et al. 2011). The RG wind has a velocity of $\sim 15\, \text{km}\, \text{s}^{-1}$ and $M_{RG} = 2 \times 10^{-7}\, \text{M}_\odot\, \text{yr}^{-1}$. We assume a timescale of $t_{\text{d}} \sim 2$ weeks as in Sect. 2.2, and an orbital separation of $a = 1.48\, \text{AU}$ (e.g. Booth et al. 2016). For estimating the wind density $n_w$ at a distance $a$, we use Eq. (5) properly normalised; we calculate the fast wind density as $n_w = 3M_w t_{\text{d}}/(4\pi m_p(V_w t_{\text{d}})^3)$. Using the expressions from Sect. 2, we obtain $n_w \sim 130\, \text{km}\, \text{s}^{-1}$, $V_w \sim 5900\, \text{km}\, \text{s}^{-1}$ for $\chi \sim 5 \times 10^{-4}$. We estimate $t_{\text{cool}}$ (see Eq. 2) for the reverse shock, propagating through the fast wind, and the forward shock that develops in the RG wind in this case. We conclude that the reverse shock is highly adiabatic whereas the forward shock is highly radiative during $t_{\text{d}}$.

The magnetic field strength near the reverse shock, $B = 2 \times 10^{-2}\, \text{G}$, is estimated assuming that the magnetic pressure is a fraction $\epsilon_B = 10^{-4}$ of the thermal pressure of the post-shock gas (see Metzger et al. 2015). The target radiation fields for IC in the vicinity of the reverse shock are the RG photon field $U_{RG}$ at a distance of $\sim a$ and the optical $U_{\text{opt}}$ emission from reprocessed X-rays (Metzger et al. 2014). The companion star of the system is a M2III giant star (Zamanov et al. 2018), and we estimate a luminosity of $L_{RG} \sim 2.5 \times 10^{38}\, \text{erg}\, \text{s}^{-1}$ (see the adopted stellar parameter in Table 2). For estimating $L_{\text{opt}}$, we use $L_{\text{sh}}$, which is a small fraction, $10^{-2}$, of the shock power $L_{\text{sh}}$, and was radiated and reprocessed into optical emission.

In Figure 11 we show the relevant timescales involved in particle acceleration, diffusion, and radiation losses in the reverse adiabatic shock. Inverse Compton losses are very efficient, giving electrons maximum energies of $\sim 0.3\, \text{TeV}$. In the case of protons, the losses for $pp$ do not affect the acceleration, which is limited only by the timescale of the event giving $E_{p,\text{max}} \sim 30\, \text{TeV}$. The diffusion of particles into the region of the contact discontinuity is fast, allowing the particles to further radiate there.

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1. http://ooruri.kusastro.kyoto-u.ac.jp/mailarchive/vsnet-alert/26131
2. We do not consider any effect from the neutrals that might affect particle acceleration (see e.g. Metzger et al. 2016).
The maximum energies of electrons and protons are compatible with high and very-high-energy gamma emission.

RS Oph was detected for the first time at gamma rays on August 2021 by Fermi Large Area Telescope (LAT) operating from 20 MeV to 300 GeV; the satellite detected a transient gamma-ray source positionally consistent with the nova. For its part, H.E.S.S., which operates in the energy range 10 GeV to 10 TeV, also detected a very-high-energy gamma-ray excess compatible with the direction of RS Oph. Following the detection, H.E.S.S. observations continued during the nova outburst. Here we show with a simple estimation that an adiabatic reverse shock is expected in RS Oph, and that it is capable of accelerating electrons and protons to high energies. The expected maximum energies are 0.3 TeV for electrons and 30 TeV for protons. In such dense environments, we expect significant pp emission up to 3 TeV. IC and Bremsstrahlung would be important at hundreds of GeV. Additional high-energy radiation, especially Bremsstrahlung and pp, is expected when particles diffuse to the contact discontinuity and compressed RG wind regions. The interaction of these high-energy particles can explain the detected gamma-ray emission. A detailed model of the system will be possible when observations become available.

6. MHD scaling of YSO jets and novae to laboratory experiments

In order to extend the scope of this work, we present preliminary estimates of the feasibility to scale the shocks from YSO jets and novae outflows to laboratory experiments. Laboratory experiments with dense, magnetised plasmas provide a novel approach to the study of astrophysical jets and outflows. The experiments are typically conducted on high-power laser and pulsed-power facilities (Remington et al. 2006), with each experimental approach allowing for different ranges of plasma parameters which can be chosen to match different regimes of interest. Typically, the experiments are characterised by temperatures of approximately hundreds to thousands of eV, flow velocities of approximately hundreds to thousands of km s⁻¹, plasma volumes of ~mm-cm³, timescales of approximately 1 to 100 ns, and electron densities of ≥ 10¹⁸ cm⁻³. The effect of magnetic fields can be controlled in the experiments depending on the way the plasma is driven in the experiment. In the case of pulsed-power generators, the magnetic field is produced from the strong electrical currents that drive the plasma, whereas in the case of laser experiments they can be generated from strong gradients of electron density and temperature in the plasma due to the Biermann battery effect, or introduced externally using capacitor-coil targets (Santos et al. 2018) or pulsed-power systems like MIFEDS (Fiksel et al. 2015).

MHD scaling arguments (e.g. Ryutov et al. 1999, 2000; Fäl-ize et al. 2011; Cross et al. 2014) make it possible to study astrophysical processes through laboratory experiments, for instance the launching and propagation of YSO jets (Lebedev et al. 2019) and accretion shocks (Van Box Som et al. 2017). Recently, the self-similar dynamics of the collision between adiabatic and radiative supersonic flows was studied by Gintrand et al. (2021), including an analysis of their scaling to laboratory experiments. Following the details given by Ryutov et al. (1999), the scaling is based on five characteristic physical parameters for the astrophysical and laboratory systems: length scale (R), density (n), pressure (P), velocity (v), and magnetic field (B). The ratio of length scales, density, and pressure result in three arbitrary scaling factors a, b, and c, which are further combined to constrain the timescale (t) and magnetic field.

The two systems will evolve identically if the initial conditions for the physical parameters are geometrically similar and the two dimensionless parameters, the Euler number $E_t = v \sqrt{P/\rho}$ and thermal plasma beta $\beta = 8 \pi P/B^2$, are the same. The scaling will be valid if both systems have a fluid-like behaviour, i.e. a localisation parameter $\delta < 1$, and negligible dissipation processes, i.e. Reynolds ($Re$), Peclet ($Pe$), and magnetic Reynolds numbers ($Re_M$) ≳ 1.

Table 2 summarises the MHD scaling for YSO jets and novae to a possible laboratory experiment. We fix the input astrophysical parameters (e.g. based on Table 1) and propose sensible laboratory parameters that match the scaling, resulting in scaling factors $a = 1$ and $b = c = 10^{18}$. For both YSO jets and novae, the laboratory parameters result in lengths scales of $R_{lab} \sim 1$ mm, densities of $n_{lab} \sim 10^{19}$ cm⁻³, pressures of $P_{lab} \sim 10^3$ bar, velocities of $v_{lab} \sim 100 - 700$ km s⁻¹, magnetic fields of $B_{lab} \sim 1 - 10$ T, and timescales of $t_{lab} \sim 1$ ns. The temperature $T_{lab} \sim 1$ keV was obtained under the assumption of astrophysical temperatures $T_{astro} = 50$ eV (5.8×10¹⁵ K) which is in line with those obtained in the simulations in Fig. 3. The dimensionless parameters in Table 2 fulfil the MHD scaling, with the exception of the Peclet number which for the laboratory case is ~1.

The parameters for scaled laboratory experiments are in line with plasma conditions achievable on current high-energy density facilities, such as the OMEGA laser at the University of Rochester, and future energetic, high-repetition lasers such as ELI-Beamlines in the Czech Republic (Jourdain et al. 2021).

7. Summary and conclusions

We study the interaction regions of YSO jets with an ambient medium and of classical novae outflows with previous ejected material. We show in Section 2 that for certain values of the jet or outflow and ambient densities, the working surface in both sources is composed of an adiabatic and a radiative shock. This particular system, in which the bow shock is radiative and the reverse shock is adiabatic, is very common in other astrophysical sources such as stellar bow shocks (e.g. del Valle & Pohl 2018). This shock combination is of interest for particle acceleration and subsequent non-thermal radiation. Particles are expected to be efficiently accelerated in strong adiabatic shocks, while the radiative shock produces a strong compression of the plasma (and magnetic field).

High-energy particles accelerated in the reverse adiabatic shock diffuse up to the dense layer downstream of the radiative shock where they can undergo further re-energisation by compression (e.g. Enßlin et al. 2011) and also an enhancement of radiative losses in the denser layer in the form of relativistic Bremsstrahlung for leptons and proton-proton inelastic collis-
ions in the case of hadrons. We make order-of-magnitude estimates for the timescales and gamma-ray luminosity for the case of the YSO jet. In the case of novae, we model the source RS Oph recently detected for the first time in gamma rays. Our estimations indicate that the reverse shock is adiabatic and might accelerate particles up to high energies that could be responsible for the observed emission.

We find that the parameters for scaled laboratory experiments for YSO jets and nova outflows are in line with plasma conditions achievable in current high-power laser facilities. This provides new laboratory astrophysics working scenarios, especially in the case of novae outflows, which had never been explored with this approach before now.

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References

Ackermann, M., Ajello, M., Albert, A., et al. 2014a, Science, 345, 554
Ackermann, M., Ajello, M., Albert, A., et al. 2014b, Science, 345, 554
Anupama, G. C. 2008, in Astronomical Society of the Pacific Conference Series, Vol. 401, RS Ophiuchi (2006) and the Recurrent Nova Phenomenon, ed. A. Evans, M. F. Bode, T. J. O'Brien, & M. J. Darnley, 31
Araudo, A. T., Padovani, M., & Margon, A. 2021, MNRAS, 504, 2405
Aydi, E., Sokolovsky, K. V., Chomiuk, L., et al. 2020, Nature Astronomy, 4, 776
Blandford, R. D. & Payne, D. G. 1982, MNRAS, 199, 883
Blandford, R. D. & Cowie, L. L. 1982, ApJ, 260, 625
Blandford, R. D. & Payne, D. G. 1982, MNRAS, 199, 883
Booth, R. A., Mohamed, S., & Podsiadlowski, P. 2016, MNRAS, 457, 822
Bode, M. F. 1987, in RS Ophiuchi (1985) and the Recurrent Nova Phenomenon, ed. M. F. Bode, 241
Booth, R. A., Mohamed, S., & Podsiadlowski, P. 2016, MNRAS, 457, 822
Booth, R. A., Mohamed, S., & Podsiadlowski, P. 2016, Monthly Notices of the Royal Astronomical Society, 457, 822
Calderón, D., Cuadra, J., Schartmann, M., et al. 2020, MNRAS, 504, 2405
Caprioli, D. & Spitkovsky, A. 2014, ApJ, 783, 91
Carrasco-González, C., Rodríguez, L. F., Anglada, G., et al. 2010, Science, 330, 1209
Chomiuk, L., Metzger, B. D., & Shen, K. J. 2020, arXiv e-prints, arXiv:2011.08751
Crosse, J. E., Reville, B., & Gregori, G. 2014, ApJ, 795, 59
de Gouveia Dal Pino, E. M. 2005, Advances in Space Research, 35, 908, fundamentals of Space Environment Science
del Valle, M. V. & Pohl, M. 2018, ApJ, 864, 19
Dobrzycka, D. & Kenyon, S. J. 1994, AJ, 108, 2259
Doss, F. W., Kline, J. L., Flippo, K. A., et al. 2015, Physics of Plasmas, 22, 056303
Erdl, S., Pfummer, C., Miniati, F., & Subramanian, K. 2011, A&A, 527, A99
Falize, E., Michaut, C., & Bouquet, S. 2011, ApJ, 730, 96
Feeney-Johansson, A., Purser, S. J. D., Ray, T. P., et al. 2019, ApJ, 885, L7
Fiksel, G., Agliata, A., Barnak, D., et al. 2015, Review of Scientific Instruments, 86, 016105
Franckowiak, A., Jean, P., Wood, M., Cheung, C. C., & Busson, S. 2018, A&A, 609, A120
Fujikawa, S., Yamaoka, H., & Nakano, S. 2012, Central Bureau Electronic Telegrams, 3202, 1
Gurian, S. M., DeRosa, R. J., & Vurm, I., et al. 2014, MNRAS, 442, 713
Hartigan, P. 1989, ApJ, 339, 987
Heathcote, S., Reipurth, B., & Raga, A. C. 1998, AJ, 116, 1940
Jourdain, N., Chaufrignon, U., & Mavlyk, H., et al. 2021, Matter and Radiation at Extremes, 6, 015401
Kuranz, C. C., Park, H. S., Huntington, C. M., et al. 2018, Nature Communications, 9, 1564
Lebedev, S. V., Frank, A., & Ryutov, D. D. 2019, Rev. Mod. Phys., 91, 025002
Li, K.-L., Metzger, B. D., Chomiuk, L., et al. 2017, Nature Astronomy, 1, 697
Livio, M. 1999, Phys. Rep., 311, 225
Martin, P., Dufresne, J. L., Gillet, S., & Point, V. 2014, A&A, 564, A128
Metsger, B. D., Caprioli, D., Vurm, I., et al. 2016, MNRAS, 457, 1786
Metsger, B. D., Finzell, T., Vurm, I., et al. 2015, MNRAS, 450, 2739
Metsger, B. D., Hascoët, R., Vurm, I., et al. 2014, MNRAS, 442, 713
Mignone, A., Bodo, G., Massaglia, S., et al. 2007, ApJS, 170, 228
Nelson, T., Mihajlovic, M., & Fisk, P. 2016, ApJ, 827, 121
O'C Drury, L., Duffy, R., & Kirk, J. G. 1996, A&A, 309, 1002
Raga, A. C., Canto, J., & Cabrit, S. 1998, A&A, 332, 714
Remington, B. A., Drake, R. P., & Ryutov, D. D. 2006, Rev. Mod. Phys., 78, 755
Rodríguez-Kamenetzky, A., Carrasco-González, C., Araudo, A., et al. 2017, The Astrophysical Journal, 851, 16
Ryu, T., Drave, R. P., Kane, J., et al. 1999, ApJ, 518, 821
Ryu, T. D., Drake, R. P., & Remington, B. A. 2000, ApJS, 127, 465
Santos, J. J., Bailly-Grandvaux, M., Ehret, M., et al. 2018, Physics of Plasmas, 25, 056705
Schure, K. M., Kosenko, D., Kastra, J. S., Keppens, R., & Vink, J. 2009, A&A, 508, 751
Shu, F., Najita, J., Ostriker, E., et al. 1994, ApJ, 429, 781
Steinberg, E. & Metzger, B. D. 2018, MNRAS, 479, 687
Sutherland, R. S. & Dopita, M. A. 1993, ApJS, 88, 253
Suzuki-Vidal, F., Lebedev, S. V., Ciardi, A., et al. 2015, ApJ, 815, 96
Walder, R., Folini, D., & Shore, S. N. 2008, A&A, 484, L9
Zamanov, R. K., Boeva, S., Latev, G. Y., et al. 2018, MNRAS, 480, 1363

Table 3. MHD scaling of YSO jets and novae to laboratory experiments. The first seven rows are physical (dimensional) parameters, whereas the last six rows are dimensionless parameters. Please refer to the text for further details on the different parameters.

| Parameter                        | YSO jet Scaled experiment | Novae Scaled experiment |
|----------------------------------|---------------------------|-------------------------|
| Length scale [$\mathcal{R}$] = cm | 10$^{16}$                 | 6$\times$10$^{13}$     |
| Density [$\rho$] = cm$^{-3}$     | 10$^{3}$                  | 0.1                     |
| Pressure [$P$] = bar              | 10$^{-3}$                 | 10$^{9}$                |
| Velocity [$v$] = km s$^{-1}$     | 10$^{4}$                  | 8$\times$10$^{-8}$      |
| Magnetic field [$B$] = G         | 10$^{-4}$                 | 10$^{2}$                |
| Timescale [$\tau$] = s           | 10$^{8}$                  | 1.2$\times$10$^{6}$    |
| Temperature [$T$] = eV            | 50$^{10}$                 | 2$\times$10$^{9}$     |

Localisation parameter $\delta$

Reynolds number $Re$

Peclet number $Pe$

Magnetic Reynolds number $Re_M$

Euler number $Eu$

Thermal plasma beta $\beta$