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# Development of a new OpenFOAM solver for plasma cutting modelling

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## Abstract

A new OpenFOAM solver is presented, for the simulation of plasma cutting torches. The mathematical model that is introduced is based on the compressible Navier-Stokes equations coupled via source terms to the electric current conservation equation. Due to the conservative and hyperbolic nature of the model, a Godunov-type scheme is used for the first time in the context of plasma cutting simulation. The numerical method consists of a second-order Total Variation Diminishing (TVD) integration with flux Harten-Lax-van Leer-Contact (HLLC) Riemann solver for the flow conservation equations, coupled with a Laplace solver for the current conservation equation. An efficient formulation for the equation of state, accurately taking into account the plasma properties, is also presented. The solver is validated through a set of canonical test cases (shock tubes and 2D Riemann problems) and it is used to simulate a three-dimensional plasma cutting torch. Good agreement is found with the literature, with an improvement in the ability to deal with the shocks occurring during plasma cutting.

*Keywords:* Plasma cutting, HLLC, Shocks, Jets

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## 1. Introduction

Plasma cutting is an industrial process widely used for metal cutting. The principle of a plasma cutting torch is to establish an electric arc between the torch and a workpiece, producing a high temperature jet which impacts the piece and causes the metal of the plate to melt. Due to the strong velocity of this jet, the molten metal is pushed out of the impingement area and leaves a kerf in the plate. This process is usually cheaper than laser cutting, but it is less accurate. However, with the usage of new types of nozzles and other innovations in the design of plasma cutting torches, recent advances are constantly improving the efficiency and accuracy of plasma cutting. This is making the process more and more interesting in comparison with laser cutting.

Despite these new progresses, many physical phenomena related to plasma cutting still need to be investigated in order to better understand the process and design even more efficient torches. Since the first generation of plasma cutting torches dating from the late 1950's, many experimental studies have been carried out to analyse and improve this cutting process. Among these experimental works, several are based on spectroscopic and imaging determination of pressure, temperature, velocity and other quantities in the plasma column [1–8]. All these studies point out the presence of shock waves in the jets coming out of plasma cutting torches, showing their underexpanded aspect.

Along with these experimental investigations, several modellings of the process have been presented since the end of the 1990's [1, 4, 5, 9–17]. The first CFD study of a plasma cutting configuration has been proposed by González-Aguilar *et al* [10]. The method presented consists in solving the Navier-Stokes equations, coupled with the calculation of the electric and magnetic fields through joule heating and Lorentz force. For this purpose, a two-dimensional axisymmetric and steady calculation with a pressure-based solver using the SIMPLE algorithm [18] has been performed. The same numerical method has been used in

26 the investigations of Freton *et al* [4, 5], completed by a turbulence model. The influence of  
27 the radiation models has also been studied, and comparisons between 2D-axisymmetric and  
28 3D calculations have been done. The results obtained by these numerical studies have been  
29 validated with experimental data. Although there are various other modellings of plasma  
30 cutting available in the literature, none of them presents any new numerical method to  
31 solve the equations. All calculations that have been performed until now to simulate the  
32 plasma cutting process are pressure based and use SIMPLE-like algorithms. As for the  
33 experimental works, these numerical investigations agree on the underexpanded aspect of  
34 the high temperature jets with the observations of shock waves. In a realistic cutting  
35 configuration, the shock waves can have a significant influence on the quality of the cuts.  
36 Consequently, a thorough investigation of these shock waves is required in the context of  
37 improving plasma cutting.

38 For this study, we use OpenFOAM (OF) [19] which is an open source CFD code. The  
39 novel approach introduced in this paper has been performed with the development of a  
40 new three-dimensional transient solver. It is based on a Godunov-type numerical method  
41 with the HLLC approximate Riemann solver [20], particularly adapted in presence of shock  
42 waves. Several other extensions of the standard single-phase gas dynamic Godunov methods  
43 are found in the literature. For example, there are various cases in which these methods  
44 have been adapted for mixture two-phase flows and applied to the simulation of different  
45 phenomena, such as wave propagation [21], cavitation [22], as well as volcanic flows [23].

46 Thus, we propose for the first time to apply this extended solver to the simulation of  
47 plasma cutting. This choice can be explained by the conservative and hyperbolic nature  
48 of the presented model. Moreover, this kind of method is exactly designed to handle  
49 compressible flows with discontinuities, such as shock waves. The mathematical model that  
50 is used is close to those of previous studies. It solves the flow conservation equations, coupled

51 with the calculation of the electric potential by solving a Laplace equation, corresponding to  
52 the current conservation equation. An efficient interpolation method for the approximation  
53 of the plasma thermodynamics properties is presented as well. The cell-centered numerical  
54 scheme is second-order TVD, with a MUSCL reconstruction and a second order Runge-  
55 Kutta time integration. The main purpose of this work is to validate the new OF solver  
56 on a plasma cutting configuration. Simultaneously, this study allows to demonstrate the  
57 ability of the Godunov-type method to simulate this kind of plasma flows and deal with  
58 strong shock waves more efficiently than the commonly used pressure-based methods. Thus,  
59 the calculations have been performed with both the present solver and the pressure-based  
60 solver used by Freton *et al* [4, 5].

61 First, the mathematical model that has been used for the simulation of a plasma cutting  
62 torch is introduced: the main assumptions leading to the set of equations are discussed  
63 and the methods used for the interpolations of the plasma thermodynamic properties are  
64 detailed and validated.

65 Secondly, the numerical method that has been employed and developed in OF is pre-  
66 sented. After the introduction of a general three-dimensional finite volume discretization,  
67 the formulation of the HLLC flux that has been chosen to solve the hyperbolic part of  
68 the equations is given, followed by the MUSCL scheme used to reconstruct variables at cell  
69 faces in order to reach a second order discretization. The procedure used to calculate  
70 the source terms is also detailed and the time integration scheme is presented.

71 Then, validations of the HLLC scheme along with the MUSCL reconstruction are in-  
72 troduced with two common test cases: a one-dimensional Sod shock tube, and a two-  
73 dimensional Riemann problem.

74 Finally, the complete solver is used on a plasma cutting torch configuration. The results  
75 are compared with the ones obtained with another solver and they are validated.

76 **2. Mathematical model**

77 *2.1. Problem description*

78 The main configuration studied for the validation of the present solver is a simplified  
 79 model of a plasma cutting torch. The geometry used is shown in Fig. 1.

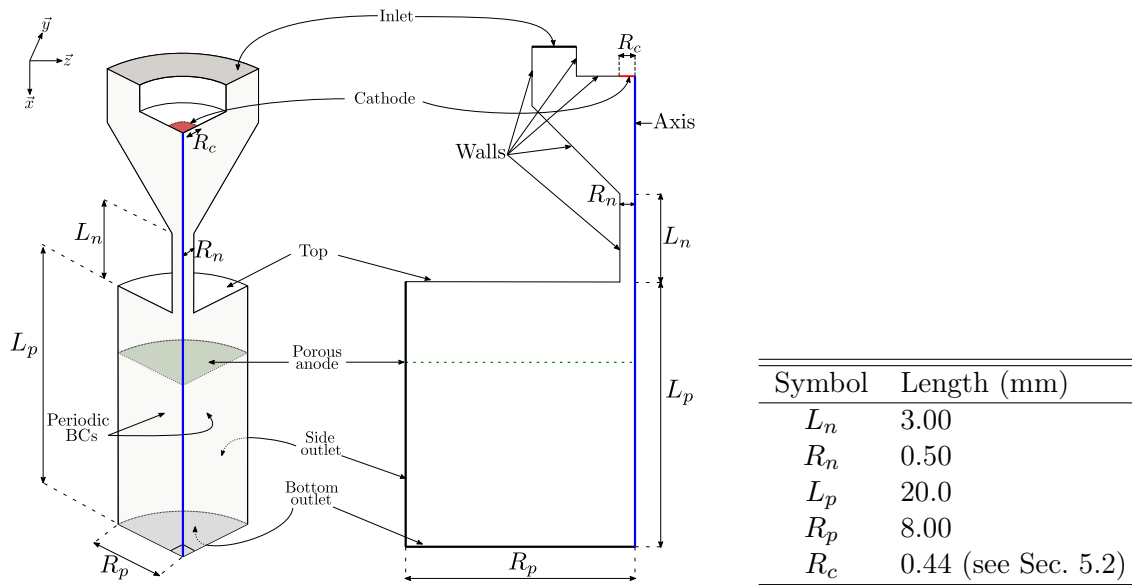


Figure 1: Details of the geometry considered for the present work.

80 The gas considered is pure air, whose properties are discussed in Sec. 2.3. The fluid  
 81 enters the plasma torch through the inlet with a high generating pressure of several atmo-  
 82 spheres. Due to a current density imposed on the cathode of radius  $R_c$ , the gas is ionized  
 83 and it becomes a plasma, conducting the current through an electric arc between the cath-  
 84 ode and the anode. Because of the joule heating and the constriction of the plasma in  
 85 the narrow nozzle of length  $L_n$  and radius  $R_n$ , the gas reaches a temperature of almost  
 86 30 000 K near the axis of the torch. Finally, the plasma is coming out of the nozzle in  
 87 the ambient air, forming a high temperature underexpanded jet in a plenum of radius  $R_p$   
 88 and length  $L_p$ . In a realistic setup, a metal plate (workpiece), which is also the anode, is

89 located a few millimeters downstream of the nozzle exit. Thus, it is impacted by the jet,  
90 whose temperature can reach several thousand kelvins. Combined with the high momen-  
91 tum of the jet, this results in the melting and the expulsion of the metal out of the piece,  
92 forming a kerf in the plate. In the present study, the interaction between the solid work  
93 piece and the jet is not taken into account. Instead, the metal plate is modeled by a porous  
94 anode [4], located 5 mm below the nozzle exit. This porous anode has no major impact  
95 on the flow, the jet can go through it without deflection, but the electric potential is fixed  
96 to 0 V downstream of it. It gives a reference value of the potential, representing the arc  
97 attachment location.

## 98 *2.2. Governing equations*

99 The plasma is considered as a compressible newtonian fluid and is assumed to be a  
100 thermal plasma at local thermodynamic equilibrium (LTE). The magnetic properties of  
101 the plasma are not considered, only the electric field is taken into account through the  
102 calculation of the electric potential, corresponding to the conservation of the current. The  
103 radiation of the plasma is modeled using a net emission coefficient assumption. The unstable  
104 areas corresponding to the arc attachment are not modeled. Since the solid workpiece is  
105 replaced by a porous anode in our model, there is no metal vapor to take into account.

106 The electric potential  $V$  is calculated by solving the following Laplace equation derived  
107 from the Maxwell-Gauss equation:

$$\nabla \cdot (\sigma \nabla V) = 0, \quad (1)$$

108 where  $\sigma$  is the electrical conductivity, depending on the local temperature and pressure.

109 The plasma flow is modeled by the compressible Navier-Stokes equations, coupled with  
110 the electric potential via source terms that take the joule heating and the radiation into

111 account:

$$\frac{\partial \mathbf{U}}{\partial t} + \frac{\partial \mathbf{F}(\mathbf{U})}{\partial x} + \frac{\partial \mathbf{G}(\mathbf{U})}{\partial y} + \frac{\partial \mathbf{H}(\mathbf{U})}{\partial z} = \mathbf{S}(\mathbf{U}). \quad (2)$$

112 In this conservative formulation of the Navier-Stokes equation,  $\mathbf{U}$  is the vector of con-  
113 served variables defined as

$$\mathbf{U} = [\rho, \rho u, \rho v, \rho w, \rho E]^\top, \quad (3)$$

114 where  $\rho$  is the density,  $u$ ,  $v$  and  $w$  are the components of the velocity vector  $\mathbf{u}$  and  $E =$   
115  $e + (u^2 + v^2 + w^2)/2$  is the total energy by mass unit, sum of internal energy  $e$  and kinetic  
116 energy.

117 The flux vectors  $\mathbf{F}(\mathbf{U})$ ,  $\mathbf{G}(\mathbf{U})$  and  $\mathbf{H}(\mathbf{U})$  are given by

$$\mathbf{F}(\mathbf{U}) = [\rho u, \rho u^2 + p, \rho uv, \rho uw, u(\rho E + p)]^\top, \quad (4)$$

$$\mathbf{G}(\mathbf{U}) = [\rho v, \rho uv, \rho v^2 + p, \rho vw, v(\rho E + p)]^\top, \quad (5)$$

$$\mathbf{H}(\mathbf{U}) = [\rho w, \rho uw, \rho vw, \rho w^2 + p, w(\rho E + p)]^\top, \quad (6)$$

118 and the source term vector  $\mathbf{S}(\mathbf{U})$  reads

$$\mathbf{S}(\mathbf{U}) = \left[ 0, (\nabla \cdot \bar{\boldsymbol{\tau}}) \cdot \vec{e}_x, (\nabla \cdot \bar{\boldsymbol{\tau}}) \cdot \vec{e}_y, (\nabla \cdot \bar{\boldsymbol{\tau}}) \cdot \vec{e}_z, \nabla \cdot (\kappa \nabla T) + \nabla \cdot (\bar{\boldsymbol{\tau}} \cdot \mathbf{u}) + \underbrace{\frac{\|\mathbf{j}\|^2}{\sigma}}_{\text{joule heating}} - \underbrace{4\pi\epsilon_N}_{\text{radiative losses}} \right]^\top. \quad (7)$$

119 In the previous expressions,  $p$  and  $T$  are respectively the pressure and the temperature,



120  $\bar{\boldsymbol{\tau}}$  is the viscous stress tensor,  $\kappa$  is the thermal conductivity,  $\varepsilon_N$  is the net emission coefficient  
 121 and  $\boldsymbol{j}$  is the current density vector. Under the Stokes hypothesis, the viscous stress tensor  
 122 for a compressible fluid can be written as

$$\bar{\boldsymbol{\tau}} = \mu \left( \nabla \boldsymbol{u} + (\nabla \boldsymbol{u})^\top \right) - \frac{2}{3} \mu (\nabla \cdot \boldsymbol{u}) \bar{\boldsymbol{I}}, \quad (8)$$

123 where  $\mu$  is the dynamic viscosity.

124

125 Finally, the current density vector is calculated as

$$\boldsymbol{j} = -\sigma \nabla V. \quad (9)$$

126 The formulation of the model (2) corresponds to the conservative form of the Euler  
 127 equations with source terms. It is conservative, and its hyperbolic nature depends on the  
 128 equation of state (EOS) discussed in the following section.

### 129 2.3. Closure models

130 For the system of equations to be fully defined, let us now detail

- 131 • the equation of state, providing the link between energy and density (related to the  
 132 conserved variables) on one hand, and pressure and temperature on the other hand,
- 133 • the procedure to compute the speed of sound  $c$  according to pressure and temperature,
- 134 • the interpolations for electrical conductivity  $\sigma$ , the dynamic viscosity  $\mu$  and the ther-  
 135 mal conductivity  $\kappa$

136 *2.3.1. Equation of state*

137 As the temperature can rise up to 30 000 K during the plasma cutting process, different  
138 reactions occurs, such as dissociation and ionisation. Consequently, the gas thermodynamic  
139 properties cannot be calculated with a simple equation of state. Due to this wide range of  
140 temperature, the transport coefficients are also pressure and temperature dependant.

141 Direct calculation of these thermodynamic properties and transport coefficient for each  
142 value of pressure and temperature is too expensive in term of computational time and is  
143 beyond the scope of the present work. Instead, two tables of calculated values have been  
144 used for each thermodynamic property and transport coefficient. One table corresponds  
145 to values calculated for temperatures going from 300 K to 30 000 K at a pressure  $p_1 =$   
146 1 bar. The second table gives the values at a pressure  $p_2 = 8$  bars for the same range of  
147 temperature. These calculated values in the tables are obtained from a plasma composition  
148 and thermodynamic properties in-house software assuming LTE [24].

149 For the computation of the thermodynamic properties at pressure values different from  
150  $p_1$  and  $p_2$ , two types of interpolations had to be performed: one allowing the calculation of  
151 the density  $\rho$  and internal energy  $e$  from pressure  $p$  and temperature  $T$ , and a second one  
152 giving  $T$  and  $p$  from  $\rho$  and  $e$ .

153 In a hyperbolic code, the first interpolation  $(p, T) \rightarrow (\rho, e)$  is only used to initialize  
154 the conservative variable vector  $\mathbf{U}$  (3); whereas the second (reverse) interpolation  $(\rho, e) \rightarrow$   
155  $(p, T)$  is required in the flux evaluation, at every grid point and time-step. The latter is  
156 therefore more critical in an efficient solver.

157 *Density and internal energy interpolations from pressure and temperature*

158 Since the numerical method used in this work is density-based,  $\rho$  and  $e$  have to be  
159 interpolated from  $p$  and  $T$  at the initial step. For this purpose, the calculated values in the  
160 tables of  $\rho$  and  $e$  at  $p = p_1$  and  $p = p_2$  have been used. We denote as  $\rho_1(T) = \rho_1(p = p_1, T)$

161 and  $\rho_2(T) = \rho_2(p = p_2, T)$  the tables of density at  $p = p_1$  and  $p = p_2$ . In the same way,  
 162  $e_1(T) = e_1(p = p_1, T)$  and  $e_2(T) = e_2(p = p_2, T)$  are the tables of internal energy.

163 For the calculation of the density, we found that a linear interpolation method is suffi-  
 164 cient:

$$\rho(p, T) = \left(1 - \frac{p - p_1}{p_2 - p_1}\right) \rho_1(T) + \left(\frac{p - p_1}{p_2 - p_1}\right) \rho_2(T). \quad (10)$$

165 Regarding the interpolation of the internal energy, a quadratic method is performed

$$e(p, T) = [1 - \xi(p)] e_1(T) + \xi(p) e_2(T), \quad (11)$$

166 with

$$\xi(p) = \begin{cases} \frac{|p - p_1|}{p - p_1} \sqrt{\frac{|p - p_1|}{p_2 - p_1}} & \text{if } p \neq p_1, \\ 0 & \text{if } p = p_1. \end{cases}$$

167 Figure 2 shows values of  $e$  and  $\rho$  interpolated at  $p = 5$  bars for different temperature  
 168 values between 300 K and 30 000 K. Note that the choice of a quadratic interpolation is not  
 169 random. The peaks in the physical properties – see, e.g. the heat capacity, corresponding  
 170 to the slope of  $e$  in Fig. 2 or thermal conductivity, Fig. 4 – are associated with dissociation  
 171 reactions (e.g.  $\text{O}_2 + \text{M} \longrightarrow 2\text{O} + \text{M}$ ) which typically have square pressure dependencies.

172 In order to assess the accuracy of these interpolated values, they are compared with the  
 173 values calculated in a third table, at  $p = 5$  bars. Relative errors between the interpolated  
 174 and calculated values are also plotted. For both variables  $\rho$  and  $e$ , the interpolation method  
 175 yields excellent accuracy: interpolation errors  $IE$  between the interpolated and calculated  
 176 values reaches a maximum of 0.03 for the density and a maximum of 0.02 for the internal  
 177 energy.

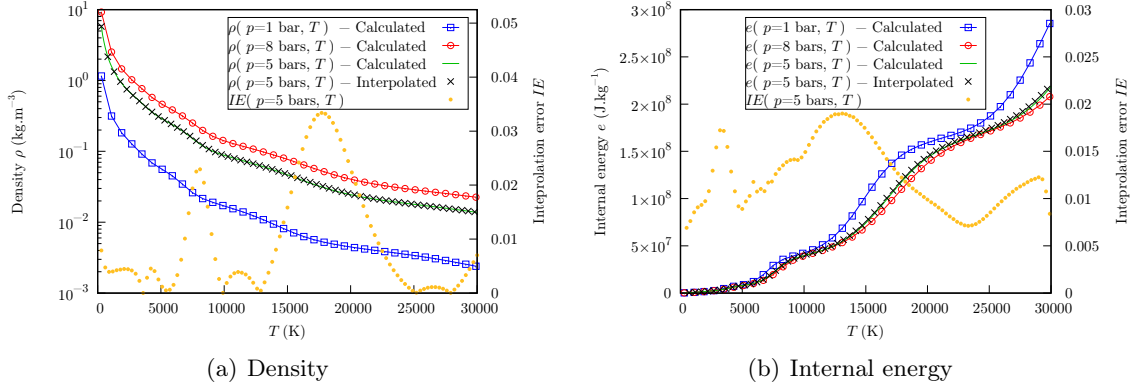


Figure 2: Interpolations of internal energy  $e$  and density  $\rho$  at  $p = 5$  bars - Error  $IE$  between interpolated and calculated values.

178 *Pressure and temperature interpolations from density and internal energy*

179 As outlined above, this interpolation is critical, as it is performed at every grid point  
 180 and time-step. The easiest and most efficient way that has been found to determine the  
 181 values of  $p$  and  $T$  according to  $\rho$  and  $e$  is achieved by using the linear interpolation of  
 182 density

$$p(\rho, e) = \left(1 - \frac{\rho - \rho_1(T_1(e))}{\rho_2(T_2(e)) - \rho_1(T_1(e))}\right) p_1 + \left(\frac{\rho - \rho_1(T_1(e))}{\rho_2(T_2(e)) - \rho_1(T_1(e))}\right) p_2, \quad (12)$$

183 where  $T_1$  and  $T_2$  are the inverse of the internal energy tables  $e_1$  and  $e_2$ :  $T_1(e) = e_1^{-1}(T)$   
 184 and  $T_2(e) = e_2^{-1}(T)$ . Once the value of the pressure is computed according to  $\rho$  and  $e$ , the  
 185 value of the temperature can easily be calculated using the expression

$$T(p, e) = [1 - \xi(p)] T_1(e) + \xi(p) T_2(e). \quad (13)$$

186 In order to assess the efficiency of this reverse interpolation, a test has been conducted  
 187 for several initial values of the pressure-temperature couple:  $(p, T)$ . For each value of  
 188  $(p, T)$ , the direct interpolation method has been performed, giving a corresponding couple

189  $(\rho, e)$ . Then, using this interpolated couple  $(\rho, e)$ , the reverse interpolation has been used  
 190 to determine a final value of the pressure-temperature couple:  $(p_f, T_f)$ . By comparing the  
 191 final couples  $(p_f, T_f)$  with the initial ones  $(p, T)$ , the following errors can be defined:

$$E_p(p, T) = \frac{|p_f - p|}{p} \quad \text{and} \quad E_T(p, T) = \frac{|T_f - T|}{T}. \quad (14)$$

192 The contour lines of the pressure error  $E_p(p, T)$  are shown in Fig. 3. A maximum error  
 193 of 0.026 is reached for a pressure of 3 bars and a temperature of 18 000 K. The same  
 194 behaviour could be observed for the temperature error  $E_T(p, T)$ , with a maximum error  
 195 below 0.016. Thus, the reverse interpolation is an accurate method that allows a low-cost  
 196 computation of the pressure and temperature according to the density and internal energy.

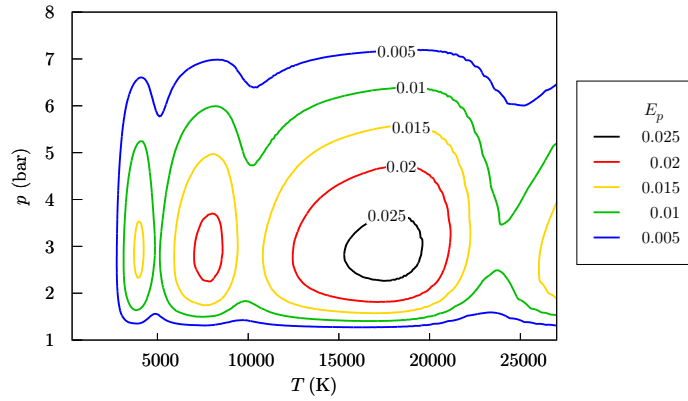


Figure 3: Contour lines for the pressure error  $E_p(T, p)$ .

### 197 2.3.2. Speed of sound calculation

198 The calculation of the speed of sound  $c$  values is required for the evaluation of the  
 199 numerical fluxes (see Sec. 3.2). In order to compute the values of  $c$  according to pressure  
 200 and temperature, the following relation has been used:

$$c(p, T) = \sqrt{\frac{\gamma(p, T) p}{\rho}}, \quad (15)$$

201 with  $\gamma$  corresponding to the pressure and temperature dependant heat capacity ratio,  
 202 which can be defined as

$$\gamma(p, T) = \frac{c_p(p, T)}{c_v(p, T)}, \quad (16)$$

203 where  $c_p$  and  $c_v$  are respectively the specific heat capacities at constant pressure and  
 204 constant volume. Regarding  $c_p$ , its values have been calculated according to temperature  
 205 at  $p = p_1$  and  $p = p_2$ :

$$\begin{cases} c_p(p = p_1, T) = c_{p,1}(T) = \left( \frac{\partial h}{\partial T} \right)_{p=p_1} = \frac{dh_1(T)}{dT}, \\ c_p(p = p_2, T) = c_{p,2}(T) = \left( \frac{\partial h}{\partial T} \right)_{p=p_2} = \frac{dh_2(T)}{dT}. \end{cases} \quad (17)$$

206 In previous relations,  $h_1(T)$  and  $h_2(T)$  are the enthalpy tables at  $p = p_1$  and  $p = p_2$ ,  
 207 which are calculated from the internal energy tables and density tables as follow:

$$\begin{cases} h_1(T) = e_1(T) + \frac{p_1}{\rho_1(T)}, \\ h_2(T) = e_2(T) + \frac{p_2}{\rho_2(T)}. \end{cases} \quad (18)$$

208 For pressure values different from  $p_1$  and  $p_2$ , the heat capacity  $c_p$  is interpolated linearly  
 209 between the two tables.

210 Concerning the calculation of the heat capacity  $c_v$ , it is done using Mayer's relation

$$c_v(p, T) = R_s(p, T) - c_p(p, T), \quad (19)$$

211 with  $R_s(p, T) = p/(\rho T)$  denoting the specific gas constant. For the present equation of

212 state, unlike the ideal gas case, the value of  $R_s$  is not constant and depends on pressure an  
 213 temperature.

214 Using Eq. (19) to express  $c_v$ , the heat capacity ratio can finally be written as

$$\gamma(p, T) = \left( 1 - \frac{R_s(p, T)}{c_p(p, T)} \right)^{-1}, \quad (20)$$

215 leading to the following expression of the speed of sound:

$$c(p, T) = \sqrt{\frac{p}{\rho \left( 1 - \frac{R_s(p, T)}{c_p(p, T)} \right)}}. \quad (21)$$

216 According to this last formulation of  $c$ , it is always positive and defined if  $R_s < c_p$ . This  
 217 condition is always true in the range of pressure and temperature considered in this study.  
 218 Consequently, as demonstrate in [25], since the speed of sound values calculated according  
 219 to the present EOS are strictly positive, they are physical and the set of equations of the  
 220 model (2) is hyperbolic.

221 To summarize, assuming that the density  $\rho$ , the pressure  $p$  and the temperature  $T$  are  
 222 known, the speed of sound  $c$  is determined in two steps:

- 223 (1) linear interpolation of  $c_p$  between the two tables at  $p = p_1$  and  $p = p_2$ ,
- 224 (2) application of Eq. (21).

225 To give an idea of the range of values in which the speed of sound lies, it has been  
 226 plotted in Fig. A.13 in the appendices, for two different pressure values:  $p = 1$  bar and  
 227  $p = 8$  bars.

228 *2.3.3. Transport coefficients*

229 The transport coefficients  $\sigma$ ,  $\kappa$  and  $\mu$  have been linearly interpolated between tabulated  
230 values, also calculated at  $p_1 = 1$  bar and  $p_2 = 8$  bars. These transport coefficients are  
231 obtained from the LTE plasma composition and the Chapman-Enskog method [24]. They  
232 take into account the reactions inside the plasma, like dissociation or ionisation. An example  
233 is given in Fig. 4 for the thermal conductivity of air at  $p = 1$  bar and  $p = 8$  bars. Some  
234 peaks can be seen, corresponding to the dissociation of nitrogen and dioxygen and to their  
235 successive ionizations. The influence of the pressure is mainly significant for temperatures  
236 above 15000 K. The temperature and pressure dependency of the viscosity  $\mu$  and electrical  
237 conductivity  $\sigma$  of air are shown in the appendices, in Fig. B.14.

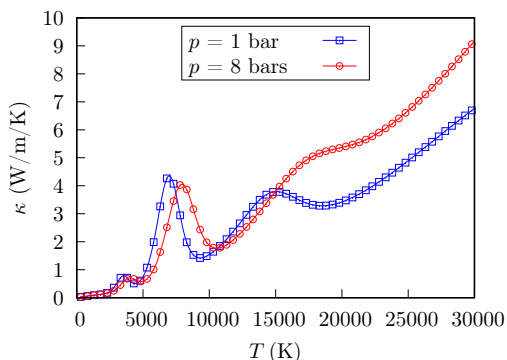


Figure 4: Calculated values of the thermal conductivity  $\kappa$  of air, at  $p = 1$  bar and  $p = 8$  bars.

238 **3. Numerical method**

239 *3.1. Spatial discretization*

240 In this work, a finite volume cell-centered method has been used, with a reconstruction  
241 of conserved variables at the centers of cells faces (see Sec. 3.3). The 3D computational  
242 domain has been discretized with an unstructured mesh, made of  $N_C$  different tetrahedral



243 or hexahedral elements. An example of two neighbouring tetrahedral cells is shown in  
 244 Fig. 5.

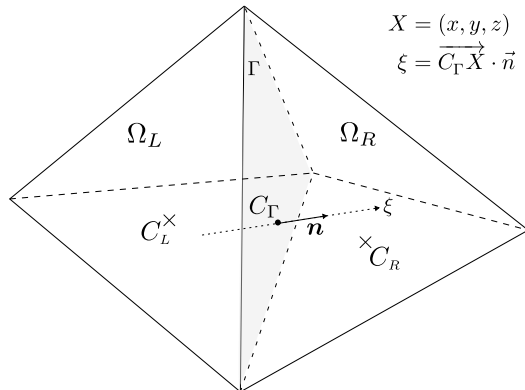


Figure 5: Example of two neighbouring tetrahedral cells  $\Omega_L$  and  $\Omega_R$ .  $C_L$  and  $C_R$  are the centroids of the cells,  $\Gamma$  denotes the face between  $\Omega_L$  and  $\Omega_R$ ,  $C_\Gamma$  is the center of  $\Gamma$  and  $\mathbf{n}$  is the unit vector normal to  $\Gamma$ .

245 In this section, we introduce the numerical scheme that has been used for the approxi-  
 246 mation of the cell-centered values  $\mathbf{U}_i(t)$  ( $i = 1, \dots, N_C$ ):

$$\mathbf{U}_i(t) \equiv \frac{1}{|\Omega_i|} \iiint_{\Omega_i} \mathbf{U}(x, y, z, t) d\Omega. \quad (22)$$

247 By integrating Eq. (2) over the volume of any cell  $\Omega_i$ , the following equation is obtained:

$$\frac{d\mathbf{U}_i}{dt} + \frac{1}{|\Omega_i|} \oint_{\partial\Omega_i} \overline{\mathbf{W}}(\mathbf{U}_i) \cdot \mathbf{n}_i d\Gamma = \mathbf{S}_i, \quad (23)$$

248 where  $\overline{\mathbf{W}}(\mathbf{U}_i) = (\mathbf{F}(\mathbf{U}_i), \mathbf{G}(\mathbf{U}_i), \mathbf{H}(\mathbf{U}_i))$  is the tensor of fluxes and  $\mathbf{n}_i$  is the outward  
 249 unit vector normal to the boundary of  $\Omega_i$ . Then, the surface integral over the boundary  
 250 may be separated in the sum of surface integrals over the  $N_{f_i}$  faces  $\Gamma_k$  ( $k = 1, \dots, N_{f_i}$ ) of  
 251 the cell  $\Omega_i$ :

$$\frac{d\mathbf{U}_i}{dt} + \frac{1}{|\Omega_i|} \sum_{k=1}^{N_{f_i}} \iint_{\Gamma_k} \overline{\mathbf{W}}(\mathbf{U}_i) \cdot \mathbf{n}_k d\Gamma = \mathbf{S}_i, \quad (24)$$

252 with  $\mathbf{n}_k = (n_k^x, n_k^y, n_k^z)^\top$  corresponding to the outward unit vector normal to the face  $\Gamma_k$ .

253 By defining  $q_k = \mathbf{u}_i \cdot \mathbf{n}_k$  the projection of the velocity on the normal vector, the projection

254  $\hat{\mathbf{F}}_k$  of the tensor of fluxes on the normal vector can be expressed as:

$$\hat{\mathbf{F}}_k = \overline{\mathbf{W}}(\mathbf{U}_i) \cdot \mathbf{n}_k = [\rho_i q_k, \rho_i u_i q_k + p_i n_k^x, \rho_i v_i q_k + p_i n_k^y, \rho_i w_i q_k + p_i n_k^z, q_k(\rho_i E_i + p_i)]^\top \quad (25)$$

255 Introducing the numerical flux  $\mathbf{F}_k$  as:

$$\mathbf{F}_k = \iint_{\Gamma_k} \hat{\mathbf{F}}_k d\Gamma, \quad (26)$$

256 Equation (24) can be simplified as

$$\frac{d\mathbf{U}_i}{dt} + \frac{1}{|\Omega_i|} \sum_{k=1}^{N_{f_i}} \mathbf{F}_k = \mathbf{S}_i. \quad (27)$$

257 Once a numerical scheme is chosen to calculate the approximation  $\hat{\mathcal{F}}_k$  of the flux value

258  $\hat{\mathbf{F}}_k$  on the face  $\Gamma_k$  (see Sec. 3.2), the numerical flux can be written as

$$\mathbf{F}_k = |\Gamma_k| \hat{\mathcal{F}}_k. \quad (28)$$

259 Thus, the following semi-discrete scheme is obtained:

$$\frac{d\mathbf{U}_i}{dt} = -\frac{1}{|\Omega_i|} \sum_{k=1}^{N_{f_i}} |\Gamma_k| \hat{\mathcal{F}}_k + \mathbf{S}_i. \quad (29)$$

260 Regarding the method used for the computation of the approximation  $\hat{\mathcal{F}}_k$ , due to the

261 conservative and hyperbolic nature of the model, a HLLC scheme (see, e.g. [20]) has been  
 262 chosen and is introduced in the next section.

263 *3.2. Numerical flux: the HLLC scheme*

264 In order to approximate the values of  $\hat{\mathbf{F}}$  on the faces of the cells, a Riemann problem  
 265 has to be solved for each face:

$$\begin{cases} \frac{\partial \mathbf{U}}{\partial t} + \frac{\partial \hat{\mathbf{F}}}{\partial \xi} = 0, \\ \mathbf{U}(\xi, 0) = \begin{cases} \mathbf{U}_L, & \text{if } \xi < 0, \\ \mathbf{U}_R, & \text{if } \xi > 0, \end{cases} \end{cases} \quad (30)$$

266 where  $\mathbf{U}_R$  and  $\mathbf{U}_L$  are the values reconstructed on the right and left side of the face, as  
 267 explained in next section. To solve these Riemann problems and compute the values of  $\hat{\mathbf{F}}$  on  
 268 the face ( $\xi = 0$ ), the approximate Riemann solver HLLC [20] has been chosen. To calculate  
 269 the HLLC fluxes, a three-wave structure is assumed. The wave speeds of the left, middle  
 270 and right waves are respectively denoted as  $S_L$ ,  $S_*$  and  $S_R$ , giving the following expression  
 271 of the HLLC flux:

$$\hat{\mathcal{F}}^{HLLC} = \begin{cases} \hat{\mathbf{F}}_L, & \text{if } 0 \leq S_L, \\ \hat{\mathbf{F}}_L^* = \hat{\mathbf{F}}_L + S_L (\hat{\mathbf{U}}_L^* - \mathbf{U}_L), & \text{if } S_L \leq 0 \leq S^*, \\ \hat{\mathbf{F}}_R^* = \hat{\mathbf{F}}_R + S_R (\hat{\mathbf{U}}_R^* - \mathbf{U}_R), & \text{if } S^* \leq 0 \leq S_R, \\ \hat{\mathbf{F}}_R, & \text{if } S_R \leq 0, \end{cases} \quad (31)$$

272 with the intermediate states  $\hat{\mathbf{U}}_K^*$  ( $K = L$  or  $K = R$ ) given by

$$\hat{\mathbf{U}}_K^* = \frac{S_K - q_K}{S_K - S^*} \begin{pmatrix} \rho_K \\ \rho_K u_K + \frac{p^* - p_K}{S_K - q_K} n_x \\ \rho_K v_K + \frac{p^* - p_K}{S_K - q_K} n_y \\ \rho_K w_K + \frac{p^* - p_K}{S_K - q_K} n_z \\ \rho_K E_k + \frac{p^* S^* - p_K q_K}{S_K - q_K} \end{pmatrix}. \quad (32)$$

273 The value of the pressure  $p^*$  is the same in both intermediate states, as it is constant  
 274 through the contact discontinuity

$$p^* = p_L + \rho_L(S_L - q_L)(S^* - q_L) = p_R + \rho_R(S_R - q_R)(S^* - q_R). \quad (33)$$

275 Regarding the wave speed  $S^*$  of the middle wave, it can be expressed as

$$S^* = \frac{p_R - p_L + \rho_L q_L (S_L - q_L) - \rho_R q_R (S_R - q_R)}{\rho_L(S_L - q_L) - \rho_R(S_R - q_R)}. \quad (34)$$

276 These previous expressions of variables in the intermediate states are obtained using  
 277 algebraic manipulations from the Rankine–Hugoniot jump conditions across the waves of  
 278 speeds  $S_L$  and  $S_R$ :

$$\begin{cases} \hat{\mathbf{F}}_L^* - \hat{\mathbf{F}}_L = S_L (\hat{\mathbf{U}}_L^* - \mathbf{U}_L), \\ \hat{\mathbf{F}}_R^* - \hat{\mathbf{F}}_R = S_R (\hat{\mathbf{U}}_R^* - \mathbf{U}_R). \end{cases} \quad (35)$$

279 More details about the derivation of the intermediate states are given in [20].

280 Regarding the left and right wave speeds  $S_L$  and  $S_R$ , they have been estimated as  
 281 proposed by Davis [26]:

$$S_L = \min(q_L - c_L, q_R - c_R) \quad \text{and} \quad S_R = \max(q_L + c_L, q_R + c_R), \quad (36)$$

282 where  $c_L$  and  $c_R$  are the left and right speed of sound values, computed as described in  
 283 Sec. 2.3.2.

### 284 3.3. Reconstruction with a flux limiter

285 In order to achieve a better precision, a second order MUSCL-type scheme [27] has been  
 286 used. For one cell  $\Omega_K$ , the idea is to reconstruct the conserved variables  $\mathbf{U}_K$  at the centers  
 287  $C_{\Gamma_k}$  of its faces. The index  $k$  takes values from 1 to  $N_{f_K}$ , where  $N_{f_K}$  denotes the total  
 288 number of faces of the cell  $\Omega_K$ . Thus, for a face  $\Gamma_k$ , the expression of the reconstructed  
 289 value is:

$$\mathbf{U}_K(C_{\Gamma_k}) = \mathbf{U}_K(C_K) + \Theta_K \overrightarrow{C_K C_{\Gamma_k}} \cdot \nabla \mathbf{U}(C_K). \quad (37)$$

290 In the previous equation,  $\Theta_K \in [0, 1]$  is a flux limiter, which can be calculated as

$$\Theta_K = \min_{k \in \llbracket 1, N_{f_K} \rrbracket} (\theta_{\Gamma_k}). \quad (38)$$

291 To prevent any overshoot in the solution, a Minmod limiter [28] has been chosen for the  
 292 calculation of  $\theta_{\Gamma_k}$ :

$$\theta_{\Gamma_k} = \max[0, \min(1, \phi_{\Gamma_k})], \quad (39)$$

293 with:

$$\phi_{\Gamma_k} = \begin{cases} \frac{\mathbf{U}^{max} - \mathbf{U}_K}{2\delta_k}, & \text{if } \delta_k > 0, \\ \frac{\mathbf{U}^{min} - \mathbf{U}_K}{2\delta_k}, & \text{if } \delta_k < 0, \\ 1, & \text{if } \delta_k = 0, \end{cases} \quad (40)$$

294 where  $\delta_k = \overrightarrow{C_K C_{\Gamma_k}} \cdot \nabla \mathbf{U}(C_K)$ , and the variables  $\mathbf{U}^{max}$  and  $\mathbf{U}^{min}$  are respectively the  
 295 maximum and minimum values between cell  $\Omega_K$  and its direct neighbors. The calculation  
 296 of the gradient at the cell center  $\nabla \mathbf{U}(C_K)$ , is performed using a Least-Squares scheme [29].

### 297 3.4. Source terms calculation

298 The components of the source term vector  $\mathbf{S}(\mathbf{U})$  defined in Eq. (7) are computed  
 299 explicitly. To calculate the current density vector  $\mathbf{j}$ , the field of electric potential  $V$  must be  
 300 determined first. This is performed at each time step by solving the Laplace Eq. (1), using  
 301 the Laplacian discretization scheme of OF. The numerical scheme that has been chosen for  
 302 this purpose is called Gauss Harmonic Corrected, using the Gauss theorem to integrate the  
 303 Laplacian. During the discretization procedure, the transport coefficient  $\sigma$  is interpolated  
 304 at the centers of the faces using an harmonic mean. In addition, a correction is performed  
 305 for non-orthogonal faces. Regarding the calculations of the gradients, it is done with a  
 306 weighted least squares method, as described in [29]. There are two main advantages of  
 307 using this method: it is made to handle strong conductivity jumps and it does not require  
 308 any correction for non-orthogonal faces. Finally, the stress tensor's divergence  $\nabla \cdot \overline{\overline{\boldsymbol{\tau}}}$  is  
 309 computed using the Gauss Linear scheme of OF [30]. With this scheme, the Gauss theorem  
 310 is used to integrate the divergence terms and the values of the fields are interpolated at the  
 311 centers of the faces thanks to a central differencing.

312 *3.5. Time integration and general procedure*

313 The time integration is performed using a second-order TVD Runge-Kutta scheme. The  
 314 timestep value  $\Delta t$  is calculated according to

$$\Delta t = CFL \min_{i \in \llbracket 1, N_C \rrbracket} \left( \frac{d_i}{\|\mathbf{u}_i\|^2 + c_i} \right), \quad (41)$$

315 where  $0 < CFL < 1$  is the Courant–Friedrichs–Lewy number and  $d_i$  is the shortest  
 316 distance between the center of the cell  $\Omega_i$  and the center of its faces:

$$d_i = \min_{k \in \llbracket 1, N_{f_i} \rrbracket} \|\overrightarrow{C_i C_{\Gamma_k}}\|. \quad (42)$$

317 Finally, the general procedure that is used at each time step of the calculation is the  
 318 following:

- 319 1. Reconstruction of conserved variables at faces centers: the gradients of conserved  
 320 variables are computed and Eq. (37) is applied.
- 321 2. Computation of HLLC fluxes through each face of the domain: the speed of sound  
 322 values are calculated according to pressure and temperature (Eq. (21)) and Eq. (31)  
 323 is used.
- 324 3. Source terms calculation (see Sec. 3.4): solving of the Laplace equation Eq. (1) to de-  
 325 termine the electric potential field  $V$  along with the current density  $\mathbf{j}$ , and calculation  
 326 of the divergence of the viscous stress tensor.
- 327 4. Time integration: the conserved variables are updated from time  $t = t^n$  to time  $t =$   
 328  $t^{n+1} = t^n + \Delta t$ . This step allows the determination of  $\rho^{n+1}$  and  $e^{n+1}$ , corresponding  
 329 respectively to the approximation of the density field and internal energy fields at

330 time  $t^{n+1}$ . At this point of the procedure, the pressure  $p^{n+1}$  and temperature  $T^{n+1}$   
 331 are not computed.

332 5. Reverse interpolation: calculation of  $p^{n+1}$  and  $T^{n+1}$  according to  $\rho^{n+1}$  and  $e^{n+1}$  using  
 333 the interpolations (12) and (13).

334 For the initialization, the temperature and pressure are imposed in the domain. Thus,  
 335 the interpolations of density (10) and internal energy (11) according to pressure and tem-  
 336 perature are performed to compute the density and internal energy at the initial time.

#### 337 4. Validations of the HLLC and MUSCL schemes

338 In order to check that the methods implemented in OF are efficient for the capture of  
 339 shocks and discontinuities, two test cases have been performed. The first one is the Sod  
 340 shock tube test case and the second one is a 2D Riemann problem. For both test cases,  
 341 only the Euler equations are solved, corresponding to Eq. (2) with the components of the  
 342 source term  $\mathbf{S}(\mathbf{U})$  set to 0. Moreover, an ideal gas law is considered for the equation of  
 343 state, with a heat capacity ratio  $\gamma$  set to 1.4.

##### 344 4.1. Sod shock tube

345 For this test case, a computational domain  $\Omega = [0, 1] \times [-0.005, 0.005] \times [-0.005, 0.005]$   
 346 has been considered and discretized with a cartesian grid of size  $200 \times 1 \times 1$ . The initial  
 347 state at  $t = 0$  s is the following:

$$[\rho, u, v, w, p] = \begin{cases} [1, 0, 0, 0, 1] & \text{if } 0 \leq x \leq 0.5, \\ [0.125, 0, 0, 0, 0.1] & \text{if } 0.5 \leq x \leq 1. \end{cases} \quad (43)$$

348 Figure 6 shows the results obtained at  $t = 0.25$  s with a  $CFL$  number of 0.9. The  
 349 numerical solution without reconstruction (first order HLLC) and the one with a MUSCL



350 reconstruction (second order) are compared with the analytical solution. Both methods are  
 351 efficient to resolve the discontinuities. As expected, a better accuracy is achieved by the  
 352 second order scheme, especially in the contact discontinuity area.

353 This case has also been performed with the second order scheme for four different grid  
 354 resolutions: 100, 200, 400 and 800 cells. The numerical results computed at  $t = 0.25$  s with  
 355 a  $CFL$  number of 0.9 on these 4 grids are compared with the analytical solution in Fig. 7.  
 356 The plotted profiles indicate that the numerical results converge to the analytical solution  
 357 when the grid is refined and the numerical scheme seems to achieve grid independence for  
 358 this test case.

#### 359 4.2. 2D Riemann problem

360 To validate our implementation, a 2D Riemann problem test case has also been carried  
 361 out. A cartesian grid of size  $400 \times 400 \times 1$  has been used with a computational domain  
 362  $\Omega = [0, 1] \times [0, 1] \times [0, 0.01]$ . At  $t = 0$  s, the initial state is divided as follow:

$$[\rho, u, v, w, p] = \begin{cases} [1, -0.75, 0.5, 0, 1] & \text{if } 0 \leq x \leq 0.5 \text{ and } 0 \leq y \leq 0.5, \\ [3, -0.75, -0.5, 0, 1] & \text{if } 0.5 \leq x \leq 1 \text{ and } 0 \leq y \leq 0.5, \\ [2, 0.75, 0.5, 0, 1] & \text{if } 0 \leq x \leq 0.5 \text{ and } 0.5 \leq y \leq 1, \\ [1, 0.75, -0.5, 0, 1] & \text{if } 0.5 \leq x \leq 1 \text{ and } 0.5 \leq y \leq 1. \end{cases} \quad (44)$$

363 Although there is no analytical solution for this test case, many references are available  
 364 in the literature. The results obtained with the present scheme and a  $CFL$  number of 0.9  
 365 are compared with the ones of [31] in Fig. 8, which shows the contour lines of the density  
 366 field at  $t = 0.3$  s.

367 These results indicate that the method used in this work is able to capture contact

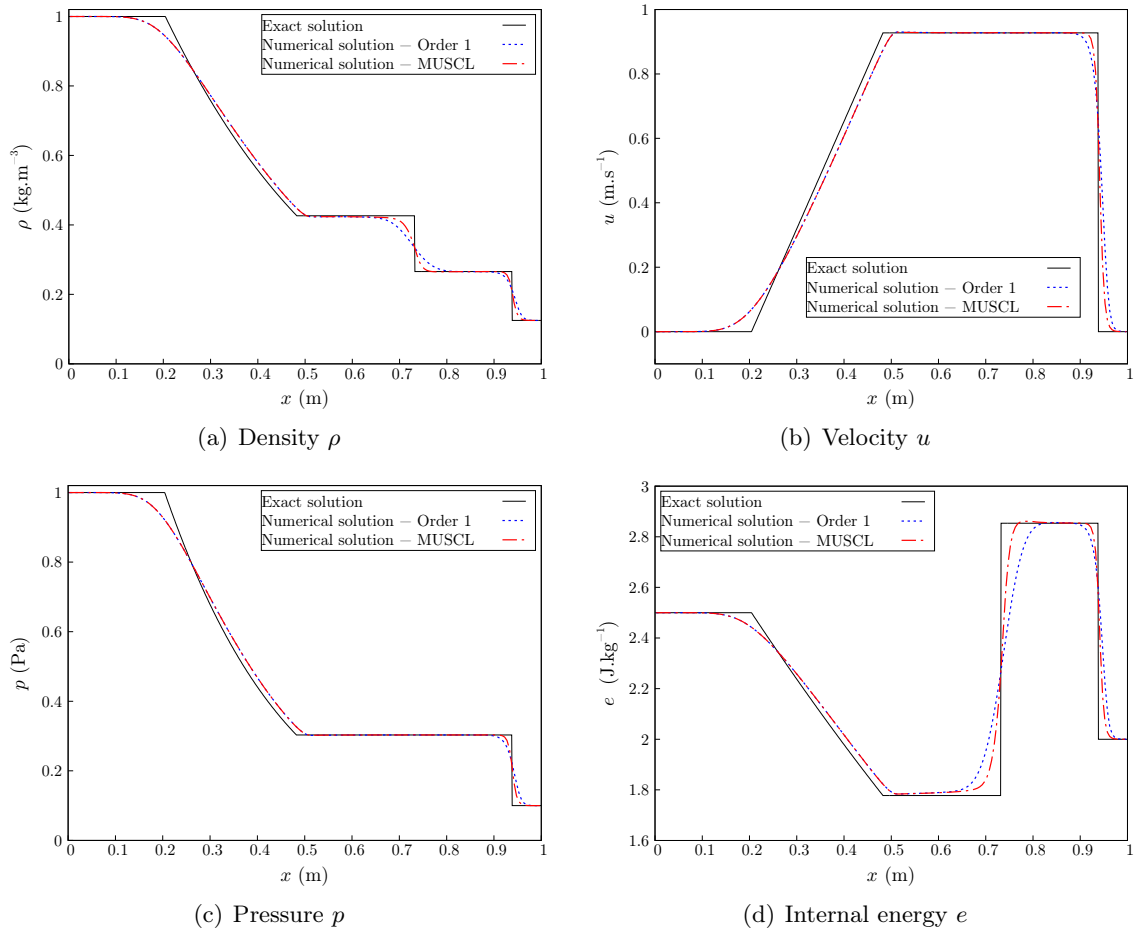


Figure 6: Profiles along the  $x$ -axis for the Sod shock tube test case at  $t = 0.25$  s ( $CFL = 0.9$ ). Comparison between the exact solution, the numerical solution without reconstruction (order 1) and the numerical solution with the MUSCL reconstruction.

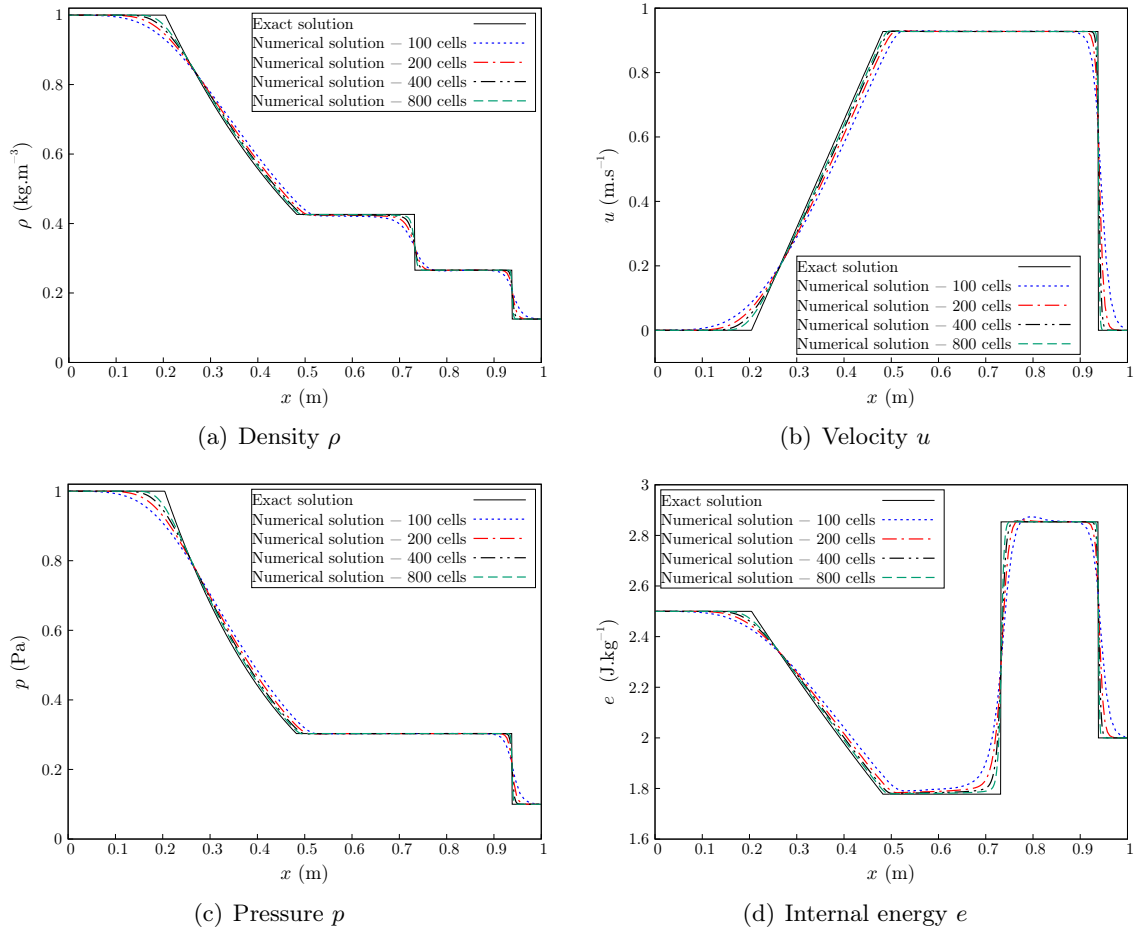


Figure 7: Profiles along the  $x$ -axis for the Sod shock tube test case at  $t = 0.25$  s ( $CFL = 0.9$ ). Comparison between the exact solution and numerical solutions (order 2) with several mesh resolutions.

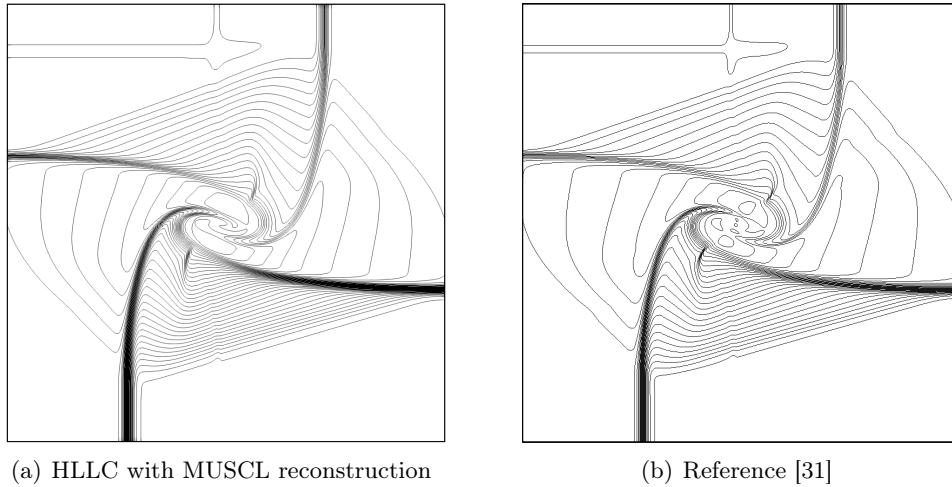


Figure 8: Density contour for the 2D Riemann problem at  $t = 0.3$  s ( $CFL = 0.9$ ).

368 discontinuities for two-dimensional cases. It has also been checked with other 2D Riemann  
 369 problems that rarefaction and shock waves are well captured with this method.

## 370 **5. Validation of the coupling with the electric potential - Simulation of a plasma** 371 **cutting torch**

### 372 *5.1. Comparison with a pressure-based method*

373 In order to validate the numerical method presented in this work, it has been compared  
 374 with a pressure-based method used on the same configuration. The geometry used for this  
 375 calculation has been presented in Fig. 1. The pressure-based method uses the SIMPLE al-  
 376 gorithm [18] of ANSYS Fluent combined with User Defined Functions. It has been assessed  
 377 and validated in various studies [4, 5]. The calculation with the current method (OF) has  
 378 been performed in 3D for a quarter of the torch, whereas a 2D-axisymmetric calculation  
 379 has been done with the pressure-based method (Fluent).

380 *5.2. Boundary conditions*

381 Boundary conditions details are given in Table 1. The cathode is considered as a slipping  
 382 wall, on which a parabolic current density profile  $J_c(r)$  is given:

$$J_c(r) = J_{\max} \left( 1 - \frac{r^2}{R_c^2} \right), \quad (45)$$

383 where  $r = \sqrt{y^2 + z^2}$  is the radial distance from the axis and  $J_{\max} = 1.6 \times 10^8 \text{ A.m}^{-2}$   
 384 is the maximum current density. This value of  $J_{\max}$  has been chosen according to an  
 385 experimental study [32]. The value of the cathode radius  $R_c$  has been calculated so that a  
 386 50 A corresponding current  $I_c$  is imposed on the cathode:

$$I_c = 2\pi \int_0^{R_c} J_c(r) r dr = 50 \text{ A}. \quad (46)$$

Boundaries	Variables			
	$p$	$T$	$\mathbf{u}$	$V$
Inlet	$p_{in} = 5 \text{ bars}$	300 K	–	$\nabla V = 0$
Cathode	$\nabla p = 0$	$\nabla T = 0$	Slip	$-(\sigma \nabla V) \cdot \mathbf{n} = J_c(r)$
Walls	$\nabla p = 0$	$\nabla T = 0$	Slip	$\nabla V = 0$
Top	$\nabla p = 0$	300 K	Slip	$\nabla V = 0$
Side outlet	1 bar	300 K	–	$\nabla V = 0$
Bottom outlet	$\nabla p = 0$	$\nabla T = 0$	$\nabla \mathbf{u} = 0$	$V = 0$

Table 1: Boundary conditions for the simulation of the plasma cutting torch.

387 *5.3. Grid independence study with the OpenFOAM solver*

388 Prior to the comparison with Fluent, the present OF solver has been used to simulate  
 389 the plasma cutting configuration with three different levels of axial grid refinement (in  
 390 the  $x$ -axis direction) in the vicinity of the shock wave. The coarsest grid has a cell size

391  $\Delta x = 1.1 \times 10^{-4}$  m in the axial direction in the shock area, which is located a few millimeters  
392 downstream of the nozzle outlet. For the most refined grid,  $\Delta x = 3.2 \times 10^{-5}$  m and for the  
393 middle level of refinement,  $\Delta x = 5.7 \times 10^{-5}$  m. The three calculations have been performed  
394 with a *CFL* number of 0.9. The profiles computed on the axe ( $r = 0$ ) for the three grids  
395 are compared in Fig. 9. Other profiles and fields are presented in the following section  
396 dedicated to the comparison between OF and Fluent. The results plotted in Fig. 9 show  
397 that the shock wave is well captured with the three meshes, even with the coarsest one.  
398 Although the amplitude of the waves is slightly damped in the case of the coarsest grid,  
399 the results obtained with second level of refinement and the maximum level of refinement  
400 are close enough to consider that the second level is sufficient in this case. Consequently,  
401 the second level of refinement ( $\Delta x = 5.7 \times 10^{-5}$  m) is the one that has been used for the  
402 comparison between OF and Fluent results.

#### 403 *5.4. Results comparison*

404 The computation with OF using the present method has been performed on a mesh  
405 made of  $4 \times 10^5$  hexahedral cells, whereas  $2 \times 10^4$  quadrilateral cells have been used for the  
406 2D-axisymmetric calculation with Fluent. Both meshes had the same refinements in the  
407 axial (see previous section) and radial directions. The transient simulation with OF has  
408 been run with a *CFL* number of 0.9 until stabilization of the mass flow rates through the  
409 inlet and the nozzle, at  $t = 3$  ms. Regarding the calculation with Fluent, it has been done  
410 with a steady solver until convergence.

##### 411 *5.4.1. Fields analysis*

412 Both methods give the same results in terms of mass flow rate and voltage: an inlet  
413 mass flow rate of  $1.47 \times 10^{-4}$  kg.s<sup>-1</sup> is obtained, along with a voltage of 132 V between the  
414 cathode and the porous anode. As the arc attachment and the sheath are not described,

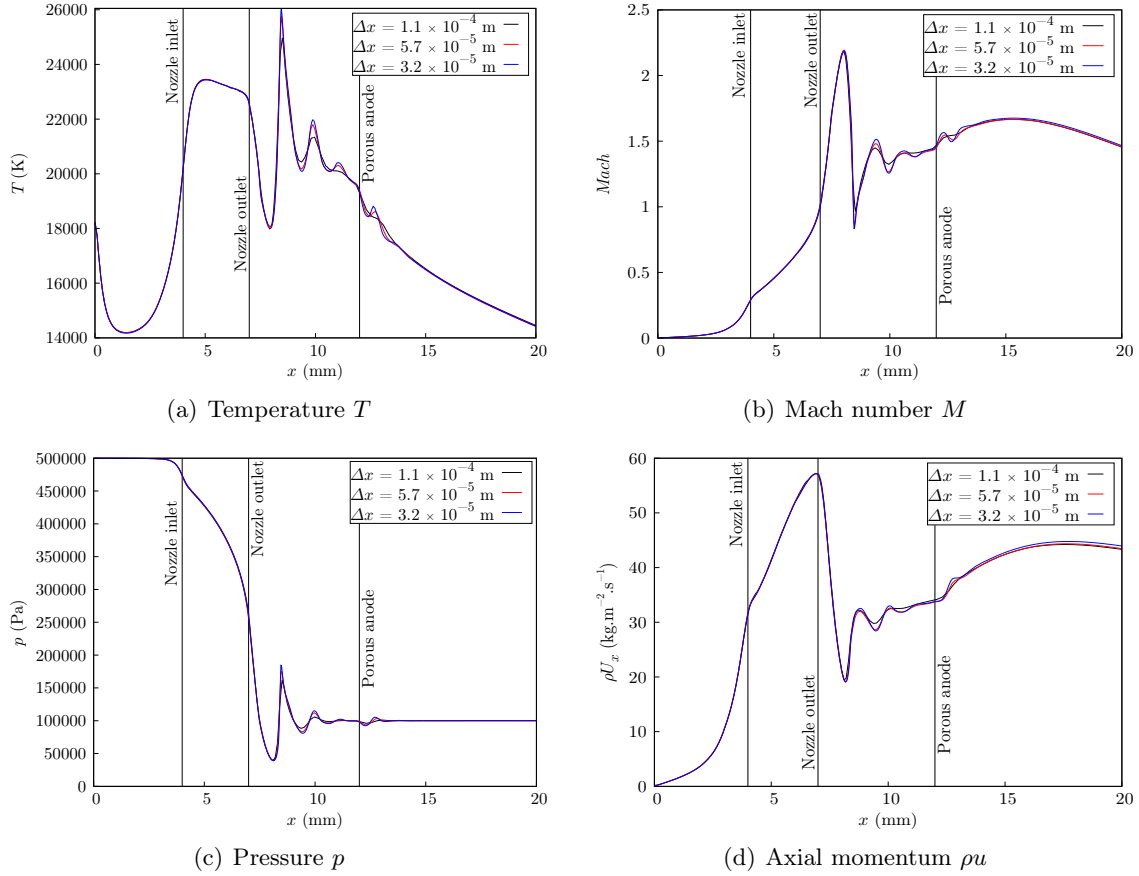


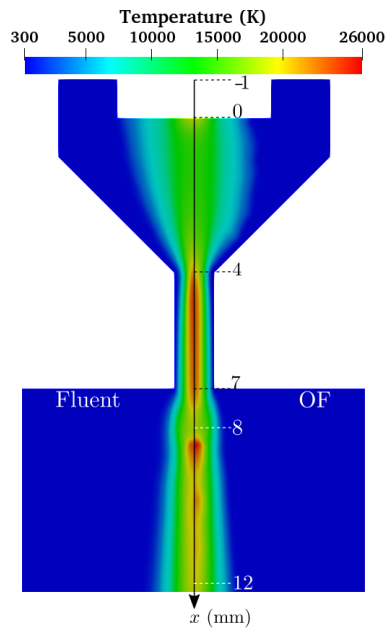
Figure 9: Profiles along the  $x$ -axis - Comparison between the different levels of grid refinement.

415 this voltage only corresponds to the plasma column. Figure 10 compares different fields  
416 computed with OF and Fluent: the temperature  $T$ , mach number  $M$ , electric potential  $V$   
417 and axial current density  $j_x$  fields are shown. Since the calculation with OF is 3D, a slice  
418 in the middle of the domain at  $z = 0$  mm is plotted. Although only one slice is shown  
419 here, it has been checked that the results were isotropic by plotting other slices for different  
420 values of  $z$ . For the four fields, a good agreement is observed overall between OF and  
421 Fluent results. The main difference appears in the top part of torch, where the plasma is  
422 more constricted towards the center of the cathode in the OF results. The behaviour is  
423 opposite at the inlet of the nozzle ( $x = 4$  cm), where Fluent results show a more constricted  
424 plasma, with higher temperature and current density than the ones calculated with OF.  
425 This can also be observed on the axial profiles shown in Fig. 11. The reasons of these  
426 differences have not been investigated precisely in the present work, since this area was  
427 not the main focus of the study. Farther downstream, as depicted in Fig. 10(b), the mach  
428 number reaches a maximum of 2.3 in the core of the underexpanded jet, upstream of a shock  
429 wave. Downstream of it, the temperature computed with OF rises up to 26 000 K, whereas  
430 it increases up to 24 000 K regarding the Fluent results. This difference of magnitude  
431 downstream of the shock wave is studied more thoroughly by analysing the axial profiles.

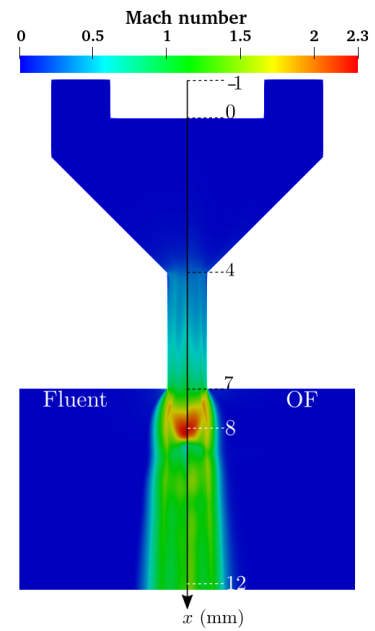
#### 432 5.4.2. Axial profiles analysis

433 Profiles along the axis ( $r = 0$ ) of  $T$ ,  $M$ ,  $p$  and  $x$ -momentum  $\rho u$  are plotted in Fig. 11.  
434 From the cathode to the shock ( $0 < x < 8$  mm), the profiles obtained with OF and Fluent  
435 are close, except for the temperature at the nozzle inlet, as mentioned previously. After the  
436 shock ( $x > 8$  mm), the values of temperature, mach number and axial momentum computed  
437 with Fluent are much lower than the values calculated with OF. As a small discontinuity  
438 is observed after the shock on the  $\rho u$  profile obtained with Fluent, the OF results seem to  
439 be more reliable. Moreover, the numerical method used with OF is suited for the capture

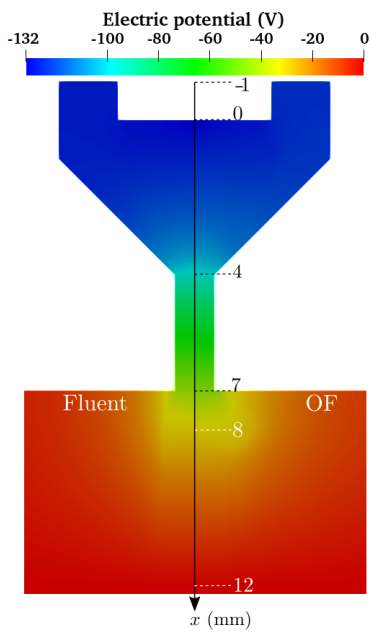




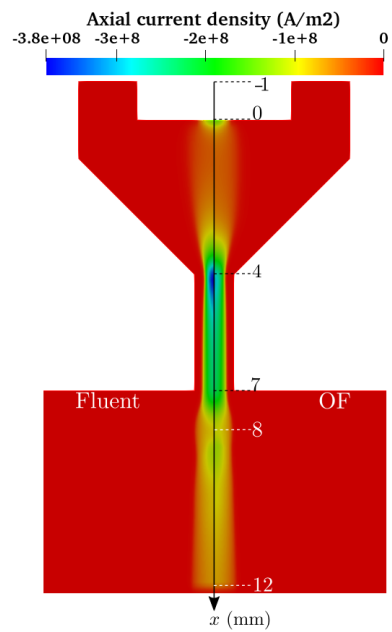
(a) Temperature field  $T$



(b) Mach number field  $M$



(c) Electric potential field  $V$



(d) Axial current density field  $j_x$

Figure 10: Comparison of fields computed with OF and Fluent -  $p_{in} = 5$  bars.

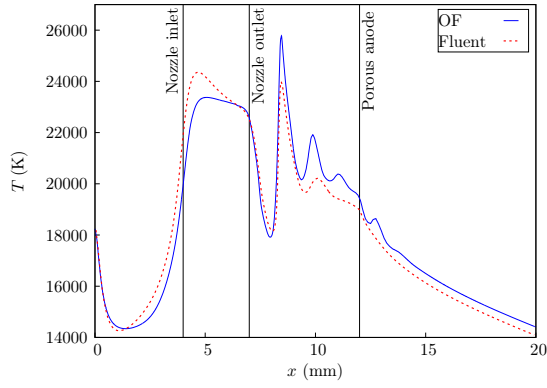
440 of shocks, as shown in the previous section.

441 In order to determine more carefully which of the two methods gives the best results  
442 downstream of the shock, more calculations have been performed and the results have been  
443 compared with the ones of other CFD softwares. These simulations have been carried out  
444 on the same geometry as the one shown in Fig. 1, without any coupling with the electric po-  
445 tential and without the diffusive source terms ( $\mathbf{S}(\mathbf{U}) = \vec{0}$ ). Thus, only the Euler equations  
446 have been solved and an ideal gas law with  $\gamma = 1.4$  has been used for the equation of state.  
447 Four different solvers have been compared: the present solver (HLLC with MUSCL) in OF,  
448 the pressure-based method of Fluent (SIMPLE algorithm), the density-based method of  
449 Fluent (Roe scheme) and the HLLC scheme of SU2 [33]. Several calculations have been  
450 done with different values of the pressure imposed at the inlet  $p_{in}$ . For  $p_{in} = 4$  bars, the  
451 results obtained with the four solvers are in good agreement with a maximal mach number  
452 of 2.0. With  $p_{in} = 5$  bars ( $\max(M) = 2.5$ ) and  $p_{in} = 6$  bars ( $\max(M) = 3.0$ ), all solvers  
453 give close results, except the pressure-based solver of Fluent. In fact, the higher  $p_{in}$ , the  
454 more Fluent's pressure-based solver results deviate from the others. This deviation only  
455 appears downstream of the shock, with still a good match with the other solvers upstream.

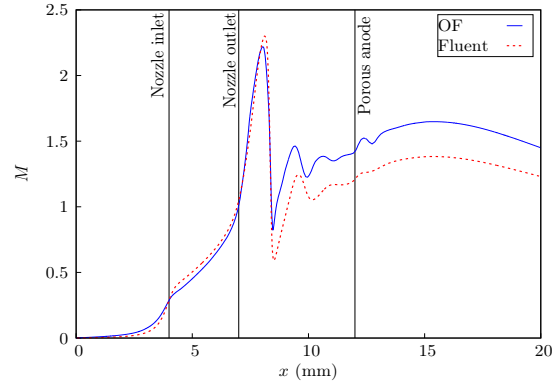
456 This complementary study leads to the conclusion that OF results shown in Fig. 11 are  
457 truly better downstream of the shocks than those of Fluent. In addition, it shows that the  
458 pressure-based solver of Fluent does not seem to be well suited for the capture of shock  
459 waves for mach number above 2.

#### 460 5.4.3. Nozzle radial profiles analysis

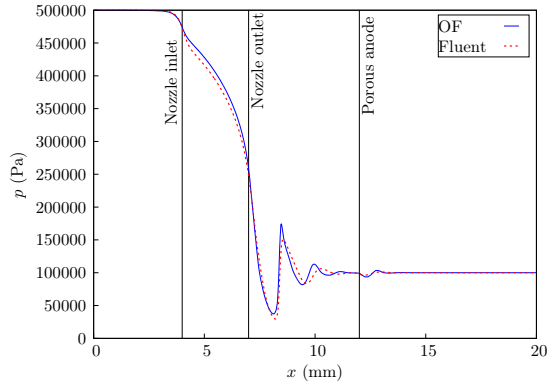
461 Figure 12 shows the radial profiles of  $T$  and  $\rho u$  computed with OF and Fluent in the  
462 middle section of the nozzle ( $x = 5.5$  mm). From the nozzle center to half its radius  
463 ( $0 < r < 0.25$  mm), OF and Fluent give very close results for both temperature and axial  
464 momentum. In the second half of the nozzle ( $0.25 < r < 0.5$  mm), close to the wall,



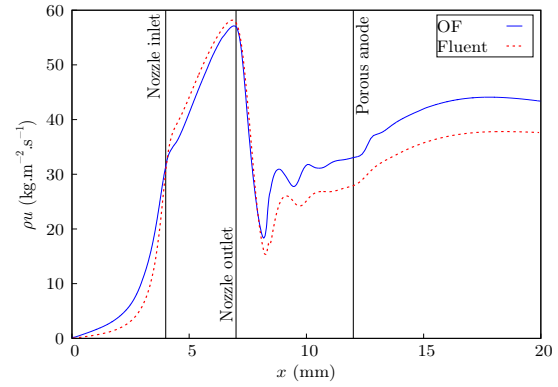
(a) Temperature  $T$



(b) Mach number  $M$



(c) Pressure  $p$



(d) Axial momentum  $\rho u$

Figure 11: Profiles along the  $x$ -axis - Comparison between OF (present method) and Fluent (pressure-based) results -  $p_{in} = 5$  bars.

465 more differences appear in the profiles. The values of temperature computed with OF in  
 466 the vicinity of the wall are lower than the ones calculated with Fluent, leading to higher  
 467 values of momentum. This might be due to a distinction in the treatment of the boundary  
 468 condition, and also to the difference of constriction near the inlet of the nozzle, as discussed  
 469 previously. The thickness of this “cold” gas layer is influenced by several parameters, such  
 470 as the pressure inlet and the intensity of the current imposed on the cathode. It could be  
 471 the focus of future investigations.

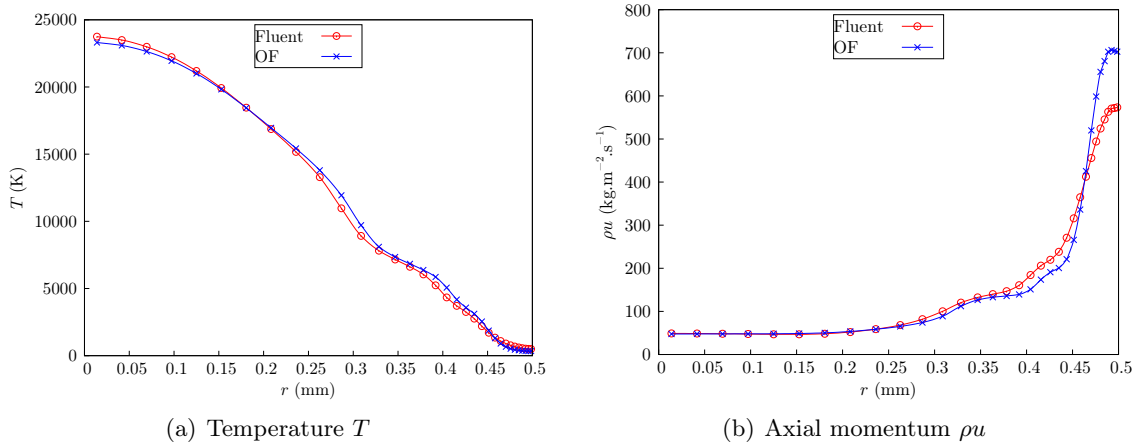


Figure 12: Radial profiles in the middle section of the nozzle ( $x = 5.5$  mm) - Comparison between OF (present method) and Fluent (pressure-based) results -  $p_{in} = 5$  bars.

## 472 6. Conclusion

473 A new OpenFOAM solver for the simulation of plasma cutting torches has been pre-  
 474 sented. It includes a new efficient formulation for the equilibrium plasma equation of state.  
 475 Through introduction of a  $2^{nd}$  order TVD scheme with HLLC Riemann solvers, the numer-  
 476 ical method is able to accurately solve discontinuities, as those present in plasma cutting  
 477 torches. This introduction of a Godunov-type scheme is novel in the scope of plasma cutting  
 478 simulation. The choice of this numerical method has been justified by the conservative and

479 hyperbolic nature of the mathematical model.

480 Following the presentation of the model and the corresponding EOS, the numerical  
481 scheme is detailed. The new OF solver is then validated through a one-dimensional shock  
482 tube and a two-dimensional Riemann problem. Finally, a three-dimensional simulation of  
483 a plasma cutting torch is presented, showing (i) results consistent with the literature [4, 5]  
484 and (ii) a much higher robustness to increasing the generating pressure.

485 Future works include the introduction of magnetic field equation, and a thorough inves-  
486 tigation of the high temperature jet downstream of the nozzle.

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585 **Appendix A. Speed of sound**

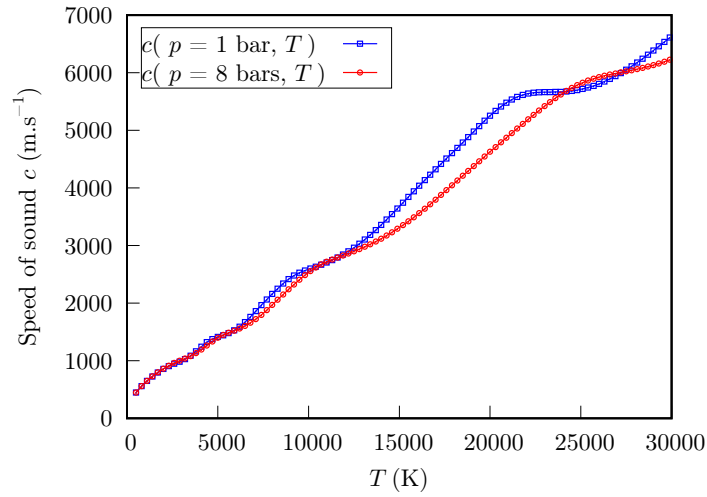


Figure A.13: Speed of sound  $c$  values at  $p = 1$  bar and  $p = 8$  bars.

586 **Appendix B. Transport coefficients**

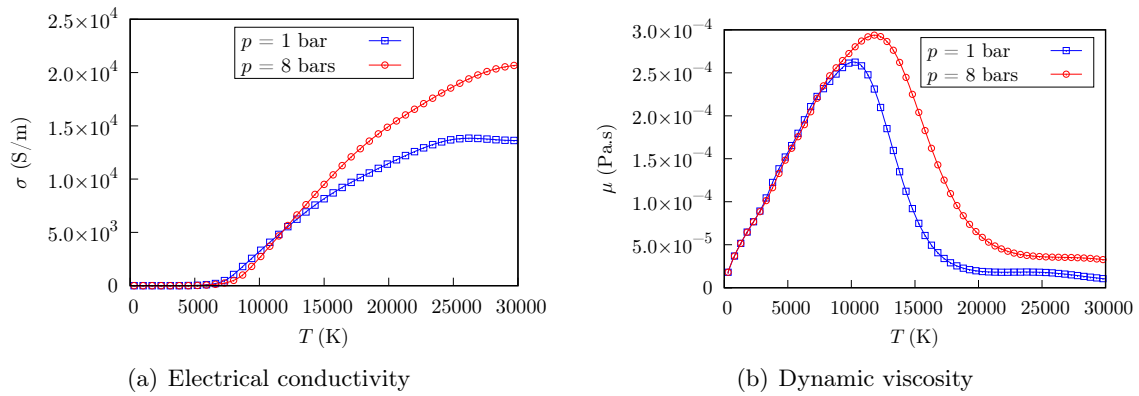


Figure B.14: Calculated values of the electrical conductivity  $\sigma$  and dynamic viscosity  $\mu$  of air, at  $p = 1$  bar and  $p = 8$  bars.