An unfortunate mistake has been identified in figure 5 of the article ‘Interaction of an atmospheric pressure plasma jet with grounded and floating metallic targets: simulations and experiments’ [1]. The correct spatial and temporal distributions of the electron impact ionization source term ($S_e$) are given in the corrected figure 5 below. The mistake results in a three order of magnitude deviation in the figure, but not in the plasma simulation. Indeed, in the referenced work the magnitude of $S_e$ has maxima between $10^{20}$ cm$^{-3}$ s$^{-1}$ and $10^{21}$ cm$^{-3}$ s$^{-1}$. These values have the same order of magnitude as those in other simulation works on He plasma jets [2–7].

This correction does not change the main conclusions of the article, namely, the electric potential of the floating target, on the presence of the return stroke with both targets and the event at the fall of the pulse and on the effects these events have on discharge parameters. However, the conclusion that the observed increases of electron density ($n_e$) (with the return stroke and with the electric field redistribution at the fall of the pulse) are mostly attributed to electron emission from metallic surfaces should be changed, as the corrected $S_e$ also contributes significantly to electron production. Thus, both processes (ionization in volume and electron emission from the target), together with electron transport in the plasma, can constitute the main sources of the experimentally and numerically observed increases of $n_e$. This correction affects the text...
Figure 5. Cross section of the spatial distribution of the corrected electron-impact ionization source term, from the simulations, with \( V_P = 5 \) kV and \( t_f = 1 \) μs at different times, for the case with a metallic target at floating potential (on top) and with a grounded metallic target (on bottom).
Interaction of an atmospheric pressure plasma jet with grounded and floating metallic targets: simulations and experiments

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Abstract
The interaction of kHz μs-pulsed atmospheric pressure He jets with metallic targets is studied through simulations and experiments, focusing on the differences between floating and grounded targets. It is shown that the electric potential of the floating target is close to grounded in the instants after the impact of the discharge, but rises to a high voltage, potentially more than half of the applied voltage, at the end of the 1 μs pulse. As a result, a return stroke takes place after the discharge impact with both grounded and floating targets, as a redistribution between the high voltage electrode and the low voltage target. Electric field, electron temperature and electron density in the plasma plume are higher during the pulse with grounded target than with floating target, as gradients of electric potential progressively dissipate in the latter case. Finally, at the fall of the pulse, another electrical redistribution takes place, with higher intensity with the highly-charged floating target than with the grounded target. It is shown that this phenomenon can lead to an increase in electric field, electron temperature and electron density in the plume with floating target.

Keywords: plasma jet, plasma-surface, metallic surfaces, floating, grounded, benchmarking

(Some figures may appear in colour only in the online journal)

1. Introduction
The interactions between non-thermal plasmas at intermediate to high pressures and surfaces is of great interest due to an increasing number of applications. These include biomedical treatment (Fridman et al 2008, Kong et al 2009, Graves 2014, Weltmann and von Woedtke 2017), surface modification (Noeske et al 2004, Cheng et al 2006, Fanelli and Fracassi 2017), catalysis (Guaitella et al 2006, Neyts 2016), and nitrification of liquids (Lindsay et al 2014). Plasma jets are very useful tools for the study of those interactions, as they are able to repetitively deliver in remote locations a wide range of reactive and charged species, high electric fields and UV photons, at atmospheric pressure and while keeping low gas temperature.
Thus, they are very promising devices for applications in the areas of material science, biomedicine (Collet et al 2014, Laroussi 2015) and agriculture (Ito et al 2018). Plasma jets have attracted a lot of attention in the past 20 years and are generally described in several reviews (Schantz et al 1998, Laime and Stori 2007, Laroussi and Akan 2007, Lu et al 2012, 2014, Winter et al 2015, Reuter et al 2018).

Several works have studied interactions of jets with targets, finding not only that the plasma affects the surface, but also that discharge dynamics can vary dramatically when interacting with surfaces with different electrical properties (Sakiyama et al 2008, Bornholdt et al 2010, Urabe et al 2010). As such, in the last years there have been several investigations of jet interactions with targets of different electrical character: dielectric at floating electric potential (Sobota et al 2013, Wild et al 2014, Guaitella and Sobota 2015, Slikboer et al 2016, 2017, Liu et al 2017, Koné et al 2017, Slikboer et al 2018a, 2018b, Ji et al 2018, Klarenaar et al 2018a, Lazarou et al 2018) (Viegas et al 2018b, Hofmans and Sobota 2019, Sobota et al 2019, Slikboer et al 2019, Viegas and Bourdon 2020, Babaeva et al 2019, Wang et al 2019); dielectric attached to a grounded plate (Hofmann et al 2014, Breden and Raia 2014, Boselli et al 2014, Guaitella and Sobota 2015, Norberg et al 2015, Wang et al 2016, Yan and Economou 2016) (Ji et al 2018, Yue et al 2018, Schweigert et al 2019b, Simoncelli et al 2019, Viegas and Bourdon 2020); conductive at floating potential (Ito et al 2010, Robert et al 2015, Koné et al 2017, Ji et al 2018, Klarenaar et al 2018a, Hofmans and Sobota 2019, Schweigert et al 2019a, Babaeva et al 2019, Martinez et al 2019); conductive grounded (Hofmann et al 2014, Boselli et al 2014, Robert et al 2015, Norberg et al 2015, Yan and Economou 2016, Darney et al 2017) (Viegas et al 2018a, Ji et al 2018, Yue et al 2018, Schweigert et al 2019b, 2019a, Babaeva et al 2019, Wang et al 2019, Simoncelli et al 2019, Martinez et al 2019).

These studies have distinguished the effect of dielectric targets of different relative permittivities $\epsilon_r$ and conductivities and of metallic targets on discharge dynamics. With low values of $\epsilon_r$ (approximately $\epsilon_r \approx 20$) and conductivity and low capacitance, after the impact of the discharge on the target, the surface of the target is charged locally in a short time, which quickly leads to the depletion of the axial component of electric field and the rise of the radial component that sustains the propagation of the discharge on the surface (Norberg et al 2015, Yan and Economou 2016, Wang et al 2016, Yue et al 2018, Klarenaar et al 2018a, Viegas and Bourdon 2020). With high values of $\epsilon_r$, or conductivity (liquid water-based surfaces) and with metallic targets, the charging of the surface is slower or inexistent, there is no radial component of electric field and no discharge propagation on the surface. Instead, a higher voltage drop remains in the gap, which promotes a return stroke and the formation of a conductive channel (Loeb 1965, Sigmond 1984, Raizer 1991, Norberg et al 2015, Yan and Economou 2016, Darney et al 2017, Viegas et al 2018a, Yue et al 2018, Klarenaar et al 2018a, Hofmans and Sobota 2019, Babaeva et al 2019). The return stroke is an ionization wave that propagates with reverse polarity with respect to the first ionization front (Loeb 1965, Sigmond 1984, Darney et al 2017). It starts at the target where electron emission takes place and propagates in an already ionized channel towards the powered electrode, provoking charge separation of opposite sign to that generated by the first wave, thus partially neutralizing the plasma channel (Raizer 1991, Viegas 2018). The return stroke is driven by the gradient of electric potential between the target and the powered electrode.

Despite some recent studies, the difference between grounded and floating conductive targets is less characterized. It has been reported by experiments that the He flow channeling in jets is stronger over grounded metallic targets than over those at floating potential (Robert et al 2015). Moreover, in Ji et al (2018), experiments and simulations of He jets have been used to compare the discharge dynamics with floating and grounded dielectric and metallic targets, obtaining higher velocities of propagation, maximum electric fields, higher species production and higher $E. coli$ cell inactivation with the grounded targets. The faster discharge propagation towards grounded targets has also been observed in Ji et al (2018), Babaeva et al (2019), Viegas and Bourdon (2020), Martinez et al (2019). The simulation results in Babaeva et al (2019) have also shown higher ionization source term during propagation and higher electric field and electron density after impact on a conductive grounded target than on a conductive floating target. The experiments in Yue et al (2018), Schweigert et al (2019a) have also observed higher production of reactive species in jets interacting with grounded conductive targets than in jets interacting with dielectric or floating conductive targets. Schweigert et al (2019a) have reported lower viability of cancer cells when using a grounded substrate under the cells during plasma jet treatment. Furthermore, it has been shown through simulations in Viegas and Bourdon (2020) that a dielectric target attached to a ground is significantly more charged by a He jet than the same target at floating potential. These differences suggest the importance of grounding or not the target for applications.

Another subject not fully understood in jet-target interaction is the discharge dynamics at the fall of the applied voltage. In jet experiments with dielectric targets using AC voltages a faint back discharge has been reported at the reversion of applied voltage polarity (Sobota et al 2013, Slikboer et al 2016). In pulsed jets, similar phenomena have been reported at the fall of the pulse. In Norberg et al (2015) the dynamics of charges at the fall of the voltage pulse has been described through simulations, not as a new discharge but as a balance between remaining positive and negative charges in the plasma and on the target surface. In Yan and Economou (2016) simulations of jets with grounded dielectric and conductive targets have observed an electric field reversal and a brief heating of electrons at the fall of the applied voltage. In the free jet experiments in Lu et al (2017) a secondary discharge has been observed at the end of the pulse and it has been attributed to the residual charges left from the first discharge. It has been found to have opposite polarity with respect to the first discharge and to be associated to an electric field below 6 kV cm$^{-1}$. Moreover, in Yue et al (2018), Klarenaar et al (2018a) a faint discharge at the fall of the voltage pulse has been reported, observed in the whole plasma channel, more pronounced with metallic target at floating potential than
with dielectric or grounded metallic targets. The faint glow has been attributed to the neutralization of the space charge in the plasma channel in Yue et al (2018). In fact, in Kim et al (2018) the electric field reversal at the falling edge of a positive voltage pulse and consequent secondary ionization have been investigated with a full kinetic treatment in argon discharges between planar electrodes on nanosecond time scales. It is claimed that the secondary ionization is induced by charge transport in the bulk plasma region. In our previous works on pulsed jets with dielectric targets (Viegas et al 2018b, Slikboer et al 2019), we have shown through comparisons between experiments and simulations on the electric field in the target induced by surface charges that the electrical redistribution at the fall of the pulse neutralizes the positive charge on the target surface and in some cases charges the target negatively.

In our last work (Hofmans et al 2020) we have characterized a kHz atmospheric pressure He plasma jet without target powered by pulses of positive applied voltage through quantitative comparisons on several key parameters of plasma jet dynamics (i.e. length and velocity of discharge propagation, gas mixture composition, electron temperature and density and peak electric field) between experimental measurements combining different diagnostic techniques and two-dimensional numerical results. Excellent agreement has been obtained between experiments and simulations on the length and velocity of discharge propagation, the gas mixture distribution and the peak electric field in the discharge front, as well as a qualitative agreement on the electron density and temperature measured behind the high field front. Moreover, we have shown how the fall of the pulse of applied voltage leads to lowering the electric potential in the plasma and, in the case of short pulses, to stopping discharge propagation. In this paper we combine different experimental diagnostic techniques (imaging, Stark polarization spectroscopy peak electric field measurements, Thomson scattering measurements of electron properties in the plasma plume and high-voltage probe measurements of the temporal evolution of electric potential of the floating target) and 2D fluid simulations to study the interaction of a positive pulsed He plasma jet with metallic targets and the influence of grounding or not the target on discharge dynamics before the impact, after the impact and after the end of the pulse. Moreover, we compare discharge parameters with metallic targets with those from the jets in Hofmans et al (2020) without target.

Firstly, both the experimental and numerical setups for free jet, jet with metallic target at floating potential and jet with grounded metallic target are described in section 2. A set of assumptions to describe the floating metallic target in the model is proposed. Then, section 3.1 describes the general discharge dynamics with the three jet configurations in both experiments and simulations, focusing on discharge propagation and the associated peak electric field. Finally, section 3.2 describes in detail the charging of the floating metallic target and shows the influence it has on plasma parameters after the impact of the discharge on the target and after the fall of the pulse. As a result, the electrical redistribution associated to the fall of the pulse and its dependence on the target are characterized.

2. Setup

2.1. Experimental setup

The plasma jet used in the experiments consists of a pyrex tube with a stainless steel tube inside as powered electrode and a grounded copper ring on the outside. Helium flows through the inner electrode with a flux of 1.5 slm. All piping in the setup is made of stainless steel to limit impurities in the gas flow. In figure 1, the dimensions of the jet are shown. This jet has the same geometry as the jet used in Sobota et al (2016), Viegas et al (2018b), Slikboer et al (2019), Hofmans et al (2020) and is operated vertically downwards. The jet is powered by positive square high voltage pulses at a repetition rate of 5 kHz. The width $t_f$ and amplitude $V_p$ of the pulses are varied to be 1 or 10 $\mu$s and 4, 5 or 6 kV, respectively. A function generator (Agilent 33220A) and high voltage power supply (Spellman UHR10P60/CL/220) couple, respectively, the pulse shape and the DC high voltage to a high voltage pulse generator (DEI PVX-4110) that supplies the high voltage pulses to the inner electrode of the jet. Three jet configurations are used: free jet, jet with metallic target at floating potential and jet with grounded metallic target.

As target, a copper plate of 8 mm $\times$ 8 mm with a thickness of 1 mm is used. For the measurements with target, the target is placed on a plastic, insulated plate that is connected to the holder of the jet itself. The distance between the target and the nozzle of the jet is set to 1 cm. In general, the distance between the target and the closest grounded plane, which is a table, is around 30 cm. The target can be grounded by connecting a cable to the ground on one end and to the target on the other end.

As in Hofmans et al (2020), the voltage and current that are applied to the jet are measured at the inner ring by a high voltage probe (LeCroy PHV4-3432) and a Rogowski coil (Pearson Current Monitor 6585), respectively. The current measured
with both the helium flow and the applied voltage, in the presence of the plasma, corresponds to the total current $I_{\text{tot}}$, while the current measured only with the applied voltage, without plasma and helium flow, corresponds to the capacitive current $I_{\text{cap}}$. This capacitive current arises from the circuit of the plasma equipment, which behaves as a capacitor. Subtracting the capacitive current from the total current yields the conductive current $I_{\text{con}}$, which is the current that flows through the plasma acting as a conductor. Figure 2 shows the applied voltage as function of time, as well as the capacitive current peak. The implications of these differences will be analysed in section 3.2.

The behavior of the plasma jet in the different cases (free jet, floating and grounded target) is studied by imaging with an ICCD camera (Stanford Computer Optics 4Picos S20Q). From these images, the position and velocity of the ionization front along the axis of the jet are determined.

To determine the electric field in the ionization front of the jet, the Stark polarization spectroscopy setup of Hofmans and Sobota (2019) is used as in Hofmans et al (2020). The He I 492.2 nm line is studied to determine the electric field from the peak-to-peak wavelength difference between the allowed and the forbidden component of the line, according to the calibration in Hofmans and Sobota (2019). Scanning the jet inside the tube and outside in the effluent yields the axial profile of the electric field at the center of the jet. Measurements of electric field are performed for the free jet and for the jet interacting with a floating as well as a grounded metallic target.

The same Thomson scattering setup as in Klarennaar et al (2015), Hofmans et al (2020) is used to measure the electron density and temperature in the plasma jet when interacting with a floating metallic target. The jet is positioned in such a way that the focus of the 532 nm laser beam lies in the center of the jet at 8.7 mm from the tube exit, thus at 1.3 mm above the target, which is the closest distance to the target possible without the laser interacting with the target. Then, the electron density and temperature are probed as a function of time. The laser operates at 100 Hz with 140 mJ per pulse and at the focal point each laser pulse has a width of 100 μm and a duration of 10 ns. This means that all scattered light within a cylindrical volume of approximately $7.9 \times 10^{-4}$ mm$^3$ and during 10 ns is captured. A volume Bragg grating is used to filter out the Rayleigh stray light from the Thomson signal (Klarennaar et al 2015). More details about this setup can be found in Klarennaar et al (2015), (2018a), (2018b). As in our previous work (Hofmans et al 2020), we take the statistical spread as error instead of the error that results from the fit. By performing measurements on different days at the same settings and at a position of 8.7 mm from the tube exit, a spread of 20% in the values of $n_e$ was found. This error is larger than the error resulting from the fit, which is only a few percent. Considering the uncertainties in the experiments as discussed in Hofmans et al (2020), among which the influence of the laser pulse on the plasma, we consider 20% as error for $n_e$ and $T_e$ more realistic.

To measure the potential on the floating metallic target as function of time, a high voltage probe (Tektronix P6015A 100 MΩ 3.0 pF) is connected to the target and the results are shown on an oscilloscope (LeCroy waveRunner 6100A 1 GHz [dual 10 GS/s, quad 5 GS/s]). This setup is shown schematically in figure 1.

2.2. Numerical setup

We use a two-dimensional axisymmetric fluid model described in our previous works (Slikboer et al 2019, Viegas and Bourdon 2020, Hofmans et al 2020). The numerical setup is shown in figure 3. The model assumes in the whole domain atmospheric pressure and room-temperature $T = 300$ K. The geometries taken are as close as possible as in the experiments. A dielectric pyrex tube with a relative permittivity of $\epsilon_r = 4$, length 3.3 cm (between $z = 0.0$ cm and $z = -3.3$ cm), internal radius $r_{in} = 1.25$ mm and outer radius $r_{out} = 2.0$ mm is used. Helium flows through the tube with a 1.5 slm flux as in the experimental conditions. A ring electrode is set inside the tube between $z = -2.8$ cm and $z = -3.3$ cm with inner radius 0.4 mm and outer radius 1.25 mm and a grounded ring is wrapped around the tube between $z = -2.0$ cm and $z = -2.3$ cm. The inner ring is powered by a positive applied voltage that
increases from zero at $t_0 = 0$ ns during 50 ns until it reaches a plateau voltage $V_P$. It is then constant until $t = t_f$ and decreases until $t_f + 50$ ns, when it reaches zero, as in Hofmans et al. (2020).

Three geometries are studied: (1) a free jet with no target present, with a grounded plane set far from the tube at $z = 20$ cm. This is the same jet configuration as in Hofmans et al. (2020); (2) a jet with conductive metallic target at floating potential placed at $z = 1$ cm with 1 mm thickness and $\sim 64$ cm$^2$ surface ($\sim 4.5$ cm radius) as in the experiments, with a grounded plane set at $z = 20$ cm. This is the case represented in figure 3; (3) a jet with conductive metallic target at grounded potential placed at $z = 1$ cm.

The metallic targets are modelled as in Viegas (2018), with infinite conductivity. To model the conductive target at a floating potential, we assume a very high relative permittivity $\varepsilon_r = 1000$, which guarantees that the target is isopotential. Unlike the grounded metal, the floating target charges anduncharges through the interaction with the plasma. However, unlike the case with dielectric surfaces, we consider that charges are conducted instantaneously inside the metallic material. Thus, we integrate the fluxes of inwards- and outwards-directed charged particles in time, to obtain the total net charge on the surface of the target $Q$ and we distribute this charge instantaneously and homogeneously in the target as net volume charge density $\rho$. This approach is similar to the one recently used in Babaeva et al. (2019), describing a floating metal as a material with $\varepsilon_r = 80$ and high conductivity. Finally, we consider the target as an ideal metal, i.e. a perfect absorber and perfect emitter, where an infinite number of free conducting charges can be exchanged with the plasma by mediation of the electric field. Therefore, we consider as for the grounded metallic target and the inner ring electrode, that electrons are emitted and absorbed and that ions are neutralized following a Neumann boundary condition for their fluxes through electric drift.

Figure 3 shows that in the model, the discharge setup is placed inside a grounded cylinder with a radius of 10 cm, to clearly define boundary conditions. The discharge dynamics is simulated through drift–diffusion-reaction equations for mean electron energy, electrons, positive ions and negative ions, and reaction equations for neutral species, coupled with Poisson’s equation in cylindrical coordinates $(z, r)$:

\[
\frac{\partial}{\partial t}(n_e \varepsilon_m) + \nabla \cdot \mathbf{j}_e = -|q_e| \mathbf{E} \cdot \mathbf{j}_e - \Theta_e \tag{1}
\]

\[
\mathbf{j}_e = -n_e \varepsilon_m \mu_e \mathbf{E} - D_i \nabla (n_e \varepsilon_m) \tag{2}
\]

\[
\frac{\partial n_i}{\partial t} + \nabla \cdot \mathbf{j}_i = S_i \tag{3}
\]

\[
\mathbf{j}_i = (q_i/|e|) n_i \mu_i \mathbf{E} - D_i \nabla n_i \tag{4}
\]

\[
\varepsilon_0 \nabla \cdot (\varepsilon_r \nabla V) = -\rho - \delta \sigma \tag{5}
\]

\[
\mathbf{E} = -\nabla V; \rho = \sum q_in_i \tag{6}
\]

where $n_i = n_e \varepsilon_m$ is the electron energy density, defined as the product of the electron density with the mean electron energy, $\Theta_e$ represents the power lost by electrons in collisions and $j_e$ is the flux of $n_e$ by drift and diffusion. $n_i$, $q_i$, $j_i$, $\mu_i$, and $D_i$ are the number density, charge, flux, mobility and diffusion coefficient of each species $i$, respectively. $S_i$ is the rate of production and destruction of species $i$ by kinetic processes and by photoionization. $V$ is the electric potential, $\mathbf{E}$ the electric field, $e$ the electron charge, $\varepsilon_0$ the vacuum permittivity, $\varepsilon_r$ the relative permittivity and $\delta$ the Kronecker delta (1 on the dielectric/gas interface). At the surface of the tube, secondary emission of electrons by ion bombardment ($\gamma = 0.1$ for all ions) is taken into account. The surface charge density $\sigma$ on the surface of the dielectric is obtained by integrating in time charged particle fluxes through electric drift to the surface. We consider that these charges then remain immobile on the surface of the dielectric.

In the experiments there is a high repetition rate ($f = 5$ kHz). However, there is uncertainty on what the exact initial conditions should be to reproduce the repetitive discharges (Naidis 2011). To take this into account, we consider, as in our previous works (Viegas et al 2018a, 2018b, Slikboer et al 2019, Viegas and Bourdon 2020, Hofmans et al 2020), a standard uniform initial preionization density $n_{\text{init}} = 10^9$ cm$^{-3}$ of electrons and $O_2^-$. However, no initial surface charges are considered on the surfaces. As in our previous works, the static flow is precalculated using COMSOL (2016) (Arjunan et al 2016, Viegas 2018). In Hofmans et al (2020), the flow calculation from Sobota et al (2016) with 1.5 slm of helium with 1000 ppm of air impurities flowing through the tube into air has been used and compared with radially-resolved Raman scattering measurements of air density ($N_2 + O_2$) in free jet configuration, yielding a good agreement. In this work, we use the same flow calculation for the free jet. For the cases with target, we use the same model to recalculate the flow for the geometry in figures 1 and 3 with a flow rate of 1.5 slm. Then, to use
the local gas mixture compositions in the plasma model, we consider that helium contains O₂ impurities and flows downstream into an O₂ environment, as an approximation to air, as in Hofmans et al. (2020). The spatial distribution of O₂ in the He–O₂ mixture obtained from the flow calculation with target is presented in figure 3.

The reaction scheme proposed in Viegas and Bourdon (2020) is used to describe the kinetics in the He–O₂ plasma, including a total of 55 reactions with 10 species. All the parameters related to electron kinetics are calculated with the electron Boltzmann equation solver BOLSIG+ (Hagelaar and Pitchford 2005), using the IST-Lisbon database of cross sections in LXCat Pancheshnyi et al. (2012), IST (2018), as functions of both the local gas mixture and the local mean electron energy $\epsilon_m$. We describe photoionization using the approach described in Bourdon et al. (2016), Slikboer et al. (2019). The ionizing radiation is assumed to be proportional to the excitation rate of helium atoms by electron impact and the photoionization source term is taken as proportional to the amount of O₂ ($X_{O_2}$) and thus we use as photoionization proportionality coefficient $A_{ph} = 100 \times X_{O_2}$ (Bourdon et al. 2016).

A finite volume approach and a Cartesian mesh are used in the model. The mesh size is 10 $\mu$m, axially between $z = -3.3$ cm and $z = 5.0$ cm (free jet case) or $z = 1.1$ cm (floating target case) or $z = 1.0$ cm (grounded target case) and radially between $r = 0$ and $r = 3.0$ mm. Then, in the rest of the domain the mesh size is expanded using a geometric progression. The average computational time required for a 2 $\mu$s simulation run to obtain the results presented in this paper was of four days with 64 MPI processes on a multicore cluster ‘Hopper’ (32 nodes DELL C6200 bi-pro with two 8-core processors, 64 GB of memory and 2.6 GHz frequency per node). Further details on the numerical schemes and other characteristics of the simulations are given in Viegas (2018).

3. Results and discussion

3.1. Characterization of discharge propagation and peak electric field

In this section, $V_F = 5$ kV and $t_f = 1 \mu s$ are used. Firstly, we compare discharge propagation with metallic targets at floating potential and at grounded potential. Figure 4 presents the experimental imaging from light emission in the two cases at different instants: during discharge propagation, at discharge impact on the target, before the fall of the pulse at $t_f = 1000$ ns and after the fall of the pulse. These emission images are wavelength integrated and show mostly emission in the range of 200–600 nm, since the sensitivity of the camera drops exponentially outside this wavelength range. From the emission spectra (not shown here) it is visible that the main sources are atomic helium (He I), the second positive system of N₂ and the first negative system of N₂⁺.

Figure 4 shows a similar propagation towards the floating and grounded targets. However, the discharge propagates faster towards the grounded target, as impact takes place at around 300 ns after the start of the pulse, which is about 60 ns earlier than in the case with the target at floating potential. Then, both cases in figure 4 show a return stroke shortly after the impact on the target. In the grounded case, light emission from the plasma persists until the end of the pulse, and is particularly high in the plume region between the tube and the target, which suggests that reactivity in the plasma persists during that time. Conversely, with the target at floating potential, the emission intensity severely decreases until the end of the pulse in the whole plasma but especially in the plume. As the applied voltage falls to zero from 1000 to 1050 ns, the emission intensity progressively decreases in the whole plasma in the grounded case, while in the floating case it increases near the inner ring electrode from $t = 900$ to 1060 ns, which suggests an electric redistribution in that region. A light emission event at the end of the pulse has also been observed experimentally in Yue et al. (2018), Klarenaar et al. (2018a), Slikboer et al. (2019) for jets with different targets. The results of
of the electron-impact ionization rate shown in figure 5. This figure presents the spatial distribution of the electron-impact ionization source term, from the simulations, with values of \( S_e \) around \( 10^{16} \text{ cm}^{-3} \text{s}^{-1} \) in the floating target case at \( t = 410 \text{ ns} \) and up to \( 10^{17} \text{ cm}^{-3} \text{s}^{-1} \) with grounded target at \( t = 350 \text{ ns} \). The return stroke is driven by the gradient of electric potential between the target and the powered electrode that transports electrons emitted from the metallic targets, as will be shown in section 3.2. Its presence with the target at floating potential suggests that the target has a low potential immediately after the impact of the discharge. With the floating target, it is visible that \( S_e \) in the plasma in the simulations decreases from the time of impact to the end of the pulse. Indeed, \( S_e \) is no longer visible at \( t = 900 \text{ ns} \) in the floating target case. Conversely, with the grounded target, \( S_e \) remains visible until the end of the pulse, with higher intensity in the plume than in the tube, with values up to \( 10^{17} \text{ cm}^{-3} \text{s}^{-1} \). Both results agree with the experimental observations in figure 4 and the simulation results in Babaeva et al (2019). These results imply that an electric field remains in the plasma between the powered electrode and the grounded target, while in the case of the target at floating potential the potential gradients in the plasma dissipate as the target is charged, as will be shown in section 3.2.

After the end of the pulse, as the voltage of the inner ring electrode falls to zero, a new dynamics takes place. Indeed, as happens with light emission from experiments, \( S_e \) increases close to the inner electrode, which requires the presence of an electric field in that region, between the grounded inner ring and the plasma. \( S_e \) after the fall of the pulse is more intense with the charged target at floating potential, reaching \( 10^{17} \text{ cm}^{-3} \text{s}^{-1} \), than with the grounded target. During the pulse, the floating target is charged, its electric potential can rise to values of the same order of the applied voltage, which is not the case with the grounded target, as will be shown in section 3.2. Then, when the applied voltage falls to zero, the gradient of potential between the new grounded electrode and the plasma is higher in the case with floating target than with grounded target.

In order to deepen the understanding on discharge propagation and quantitatively compare simulations and experiments, in figure 6 we follow the position of the discharge front in time in experiments and simulations. Besides the two cases presented in figures 4 and 5, the results with free jet (no target) and \( V_P = 5 \text{ kV} \) are also shown and compared. In the experiments, the position of the discharge front is obtained from the maximum of the light emission intensity, with an errorbar of 0.07 cm, while in the simulations it is obtained from the maximum of the axial component of electric field \( E_z \). That approach allows to follow the propagation of the first ionization wave in both experiments and simulations and of the

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**Figure 5.** Cross section of the spatial distribution of the electron-impact ionization source term, from the simulations, with \( V_P = 5 \text{ kV} \) and \( t_f = 1 \mu\text{s} \) at different times, for the case with a metallic target at floating potential (on top) and with a grounded metallic target (on bottom).

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**Table:**

| Floating Metal | Grounded Metal |
|----------------|---------------|
| 200 ns | 370 ns | 410 ns | 900 ns | 1070 ns | 1100 ns |
| \( S_e [\text{cm}^{-3} \text{s}^{-1}] \) | \( S_e [\text{cm}^{-3} \text{s}^{-1}] \) |
| 10^{16} | 10^{17} |
| 10^{15} | 10^{16} |
| 10^{14} | 10^{15} |

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**Equations:**

\[ P = V \cdot I \]

**Symbols:**

- \( P \): Power
- \( V \): Voltage
- \( I \): Current

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return stroke in simulations but not in experiments. The numerical result is shifted by 30 ns to account for the difference in time of ignition with respect to the experiments.

Figure 6 shows, as in Hofmans et al. (2020), a small difference between the ignition time in experiments and simulations of about 30 ns and an excellent agreement in discharge propagation for every case. In both experiments and simulations, the discharge propagating towards the grounded target is faster than in the other cases, all along the propagation, due to the proximity of the ground, as has also been shown in experiments and simulations in Ji et al. (2018), Babaeva et al. (2019), Viegas and Bourdon (2020), Martínez et al. (2019). Conversely, the discharge propagates faster towards the floating target than in the free jet case only when the discharge is close to the target. The ground is placed at \( z = 20 \) cm in both cases and thus does not justify the difference. Indeed, the difference in velocity is due to the influence of \( \epsilon_r \) on propagation \( (\epsilon_r = 1000 \) in the floating target case and \( \epsilon_r = 1 \) in air in the free jet case). It has been shown in Viegas and Bourdon (2020) that \( \epsilon_r \) does not significantly change the velocity of discharge propagation, except when the discharge front is very close to the target surface, mostly in the last 5 mm of propagation. There, the velocity of propagation increases with \( \epsilon_r \). A difference between different targets on electron density and peak electric field only at a few mm from the surface has also been measured in Klarenaar et al. (2018a), Sobota et al. (2019) and simulated in Ji et al. (2018), Schweigert et al. (2019b). Finally, the simulation results in figure 6 show the return stroke propagating from the target towards the inner electrode, with both grounded and floating targets. The return stroke propagates faster than the first ionization wave, in agreement with the cases in Darny et al. (2017), Viegas et al. (2018a) for grounded target and both positive and negative polarities of applied voltage.

In figure 7, the value of \( E_{\text{MAX}} \) along the propagation is presented for the three previously described cases. Both experimental Stark shift measurements and simulation results are shown. In both cases, \( E_{\text{MAX}} \) is the peak electric field in the center of the front (Hofmans and Sobota 2019), with a radial uncertainty of the size of the slit width of 100 \( \mu \)m. As explained in Hofmans and Sobota (2019), \( E_{\text{MAX}} \) in the experiments comes from the distance between the allowed and forbidden lines of the studied helium band, where the position of the forbidden line changes the most due to the high electric field. The error in the electric field values is taken as the uncertainty of the fit in determining the wavelength position of both lines, yielding values of around \( \pm 1 \) \( \text{kV cm}^{-1} \), as can be seen in figure 7. Hence, the measured \( E_{\text{MAX}} \) might not be the highest at the measured position, but actually an average value within a range of \( \pm 1 \) \( \text{kV cm}^{-1} \). This range of \( \pm 1 \) \( \text{kV cm}^{-1} \) around the maximum of \( E_z \) corresponds to a distance of around \( \pm 0.1 \) mm around its position, according to the simulations. Therefore, we take the average \( E_z \) within a distance of \( \pm 0.1 \) mm around the maximum of \( E_z \) found in the center for \( r < 0.1 \) mm, accounting for the slit width. The axial averaging of \( E_z \) has been shown in Hofmans et al. (2020) to be a more accurate way to compare simulation results with Stark shift electric field measurements than taking the local maximum of \( E_z \). Moreover, we have verified that the difference between the maximum of \( E_z \) for \( r < 0.1 \) mm and its radial average within \( r < 0.1 \) mm is negligible.

Between the three different jet configurations, in both experiments and simulations, there are only small differences...
ments in discharge structure and in figure 6, the differences between simulations and experiments are sharp in the model, while they are rounded in the experiments. As the experimental discharge is tendentiously wider close to the nozzle, the electric field is more off-centered in the simulations and thus its value at the center is lower in the experiments. Likewise, close to the electrodes, as the discharge is more centered in the experiments, \( E_{\text{MAX}} \) is higher at the center than in the simulations. Moreover, the increase of \( E_{\text{MAX}} \) just outside the tube in the simulations may be due to the change of permittivity between the tube with \( \varepsilon_r = 4 \) and the ambient air with \( \varepsilon_r = 1 \) (Jánský and Bourdon 2011). The fact that the tube edges are sharp in the model, while they are rounded in the experiments, might contribute to the difference. However, as shown in figure 6, the differences between simulations and experiments in discharge structure and \( E_{\text{MAX}} \) do not lead to a significant difference in discharge propagation velocity. Indeed, we have shown in Hofmans et al. (2020) through comparisons of discharge dynamics with different applied voltages, that \( E_{\text{MAX}} \) profiles are not directly related to discharge propagation velocities, in agreement with studies in air streamer discharges (Babaeva and Naidis 1996). Velocities are dependent on geometry and on the magnitude of applied voltage, while the electric field is related to the local charge separation.

### 3.2. Jet-target interaction

In this section, we study the dynamics taking place after the impact of the discharge on the target. Firstly, figure 8 presents the temporal evolution of the electric potential in the conductive metallic target at floating potential. Experimental and numerical results are shown for three cases of \( V_P \): 4, 5 and 6 kV. In the experiments, two different lengths of pulses are used for each case: \( t_f = 1 \) \( \mu \)s and \( t_f = 10 \) \( \mu \)s. In the simulations, only \( t_f = 1 \) \( \mu \)s is used.

Figure 8 shows that the target potential in the experiments for pulses with \( t_f = 10 \) \( \mu \)s starts increasing at the impact of the discharge and slowly rises due to electron emission and ion neutralization until saturation is reached after some \( \mu \)s at a potential slightly below \( V_P \). It is visible that both the time of impact of the discharge and the time of saturation are inversely proportional to \( V_P \). Thus, the pulse width and the applied voltage allow to control the charging of the floating target, as shown also in Slikboer et al. (2019) for a dielectric target. With short pulses of \( t_f = 1 \) \( \mu \)s, the charging of the target is interrupted. As the applied voltage in the inner ring electrode is dropped, the target changes from cathode to anode and the electric potential slowly decreases by electron absorption, reaching almost zero at \( t = 2 \) \( \mu \)s. Negative charge deposition after the fall of the pulse due to reversal of electric field direction has also been observed in experiments and in simulations with different dielectric targets (Viegas et al. 2018b, Slikboer et al. 2019, Viegas and Bourdon 2020). However, even with short pulses of \( t_f = 1 \) \( \mu \)s, the target potential reaches non-negligible values at the end of the pulse, of almost 2/3 of \( V_P \) when \( V_P = 6 \) kV, 1/2 of \( V_P \) when \( V_P = 5 \) kV and 1/4 of \( V_P \) when \( V_P = 4 \) kV. Conversely, in the simulations the target potential has a first increase with the approach of the discharge, mostly due to electron emission by effect of the electric field. Then, the potential increases faster after the discharge impact. Indeed, it rises faster than in experiments and saturates at about 2 kV below \( V_P \). In the cases with \( V_P = 5 \) and 6 kV, the saturation takes place during the 1 \( \mu \)s pulse, approximately 400 ns after the impact with \( V_P = 6 \) kV and 600 ns after the impact in the \( V_P = 5 \) kV case. After the end of the pulse, the potential also decreases faster in simulations than in experiments, decreasing to half its value in about 150 ns, instead of 400 to 700 ns registered in experiments. In both experiments and simulations the rate of charging and uncharging grows with \( V_P \). The difference between the experimental and numerical results of charging of the target will be discussed in section 3.3.

The total charge in the target can also be obtained from the simulations. Although not shown here, it follows approximately the same temporal profile as the electric potential in the target, reaching values of around 2.0, 3.6 and 4.0 nC at the end of the pulse, respectively, for \( V_P = 4 \), 5 and 6 kV. These values agree with those presented in Viegas and Bourdon (2020) for \( V_P = 6 \) kV, where it has been shown through simulations that a floating dielectric target of \( \varepsilon_r = 80 \) charges up to 2 nC in about 400 ns and a grounded dielectric target of \( \varepsilon_r = 56 \) charges up to 10 nC in the same timescale. As expected for a floating conductive target, the value obtained here for \( V_P = 6 \) kV stands between those two cases. However, as the charge is distributed in the large metallic target, 4.0 nC corresponds to only ~ 0.06 nC cm\(^{-2}\) of surface charge density. This value is much lower than those

![Electric potential in the metallic target at floating potential](image-url)
potential and

in dielectric targets that charge locally up to 70 nC cm−2 (Slikboer et al 2019, Viegas and Bourdon 2020). As the target is charged in the model through ion neutralization and electron emission, both driven by the electric field, and considering that electrons are approximately 100 times more mobile than ions, we can conclude that the number of electrons emitted during the charging is of the order of 1010 (1 nC corresponding to approximately 6 × 109 elementary charges). Considering the case with \( V_p = 6 \text{ kV} \), where the target charges approximately 1 nC per 100 ns, we can calculate a flux of electron emission through the discharge cross section of \( \sim 0.05 \text{ cm}^2 \) of approximately 0.2 nC ns−1 cm−2 or 109 ns−1 cm−2 electrons. Likewise, electrons are absorbed after the fall of the pulse with a flux \( \sim 0.3 \text{ nC ns}^{-1} \text{ cm}^{-2} \). Finally, we calculate the self-capacitance of the target as \( C = Q/V \) to be between 1.0 and 1.2 pF, where \( Q \) is the total charge in the target and \( V \) is its potential. In Itô et al (2010), with a jet powered by a voltage with peak of 6–7 kV impacting on a copper target with a flux \( \sim 0.2 \text{ nC ns}^{-1} \text{ cm}^{-2} \) or 109 ns−1 cm−2 electrons. The experimental values of \( n_e \) after the fall of the pulse show that the electrical redistribution taking place between the inner electrode (now cathode) and the plasma, limited by a target charged at 4 kV (figure 8), can effectively transport or produce a significant amount of electrons. The experimental values of \( T_e \) and \( n_e \) in this work agree with already published results (Klarennaar et al 2018a), where the same jet has been used, but add the increase in \( n_e \) after the fall of the pulse.

Despite the differences, both experimental and numerical results in figure 8 support the conclusion that the metallic target at floating potential has a voltage close to zero at the time of discharge impact, which allows it to behave approximately like a grounded target in the instants after the impact. However, at the end of the pulse the target is charged and thus its interaction with the plasma is expected to be different from that of a grounded target. Then, we analyze the consequences of jet-target interaction on the plasma. In the experiments, the temporal evolution of electron temperature \( T_e \) and electron density \( n_e \) has been measured through Thomson scattering in a jet with \( V_p = 6 \text{ kV} \) and floating copper target, in the center at \( r = 0 \) and at \( z = 8.7 \text{ mm} \), at only 1.3 mm from the target. This is represented in figures 9 and 10.

In the experiments, \( T_e \) first increases to 4 eV as the discharge arrives. As the position assessed is very close to the target, the propagation of the return stroke is not distinguishable from this first peak of \( T_e \). Then, as the return stroke propagates further into the tube, \( T_e \) decreases close to the target. However, \( T_e \) returns to 4 eV and stays with that value until the end of the pulse, which suggests a continuous reactivity in the plume. This would not be the case if the target would charge up to \( V_p = 6 \text{ kV} \), in which case the plasma would tend to be a quasineutral channel. Then, as the applied voltage falls, there is an increase in \( T_e \) to almost 5 eV, to which follows a slow decrease. The experimentally-measured \( n_e \) follows the same evolution, remaining close to \( 10^{14} \text{ cm}^{-3} \) during the pulse and then increasing after the fall of the pulse. The increase of \( n_e \) after the fall of the pulse shows that the electrical redistribution taking place between the inner electrode (now cathode) and the plasma, limited by a target charged at 4 kV (figure 8), can effectively transport or produce a significant amount of electrons. The experimental values of \( T_e \) and \( n_e \) in this work agree with already published results (Klarennaar et al 2018a), where the same jet has been used, but add the increase in \( n_e \) after the fall of the pulse.

The experimental \( T_e \) and \( n_e \) are compared with the simulation results in the same figures (figures 9 and 10) for the same case. The simulation results of \( T_e \) and \( n_e \) have been retrieved every 10 ns without temporal averaging, as 10 ns is also the duration of each measurement. The numerical results are presented both locally at \( r = 0 \) and \( z = 8.7 \text{ mm} \) and averaged within the volume of the laser beam in the Thomson scattering measurements, i.e. within a cylinder of 50 μm radius and 100 μm length centered at \( r = 0 \) and \( z = 8.7 \text{ mm} \). The temporal evolution of the axial component of electric field \( E_z \) is also presented, in figure 11, at the same position and in the middle of the plume, at \( r = 0 \) and \( z = 5 \text{ mm} \). This quantity is not accessible in experiments and therefore is represented exclusively as a simulation result with a resolution of 1 ns.
The numerical result for the case with floating target and $V_p = 6$ kV (black and grey curves) qualitatively agrees with the experimental measurements. Firstly, the comparison of the two numerical curves of $T_e$ and $n_e$ allows to conclude that although the averaging affects $n_e$, by about 20%, its effect is invisible on $T_e$ and it is not a fundamental factor when analysing the data in this case, due to the relatively small volume of the laser beam. $T_e$ has a peak as the discharge front impacts the target, within the errorbar of the experimental one. The peak takes place 40 ns later than in the experiments due to the difference in time of impact. Then, $T_e$ falls to 2 eV and increases slowly during the pulse until 3 eV. At the end of the pulse, the numerical $T_e$ has a sudden drop and a sudden increase. Although the drop is not obtained by the measurements, the increase is in agreement between simulations and experiments. The numerical values of $T_e$ are generally lower than the experimental ones. We should notice that $T_e$ is obtained from Thomson scattering measurements assuming that the lowest energy electrons (most of the population) follow a Boltzmann EEDF. However, the EEDFs calculated from Bolsig$+$ present deviations from the Boltzmann EEDFs. Indeed, the EEDFs have more populated bulk and less populated tail than the equilibrium solution. Thus, the Boltzmann assumption potentially leads to an overestimation of $T_e$ from Thomson scattering measurements. This effect has been quantified for argon microwave discharges in Ridenti et al (2018), leading to differences in $T_e$ up to a factor 4.

For the same case with floating target and $V_p = 6$ kV (black and grey curves), $n_e$ follows approximately the same evolution in simulations and experiments, but is more than one order of magnitude lower in simulations than in experiments. This difference has already been observed in Hofmans et al. (2020) in a free jet case and is discussed in that paper, along with the values of $n_e$ and $T_e$ in experiments and simulations. In that work we have verified that the difference between simulations and experiments is in agreement with literature and that it is not expected to be due to any perturbation of the studied discharge by the laser used in the experiments. Then, as in Hofmans et al. (2020), we assume that the difference may be related to the assumption of oxygen instead of air in the model and to the unknown memory effects of discharge repetition.

$E_z$ presented in figure 11 for the floating target case at $z = 8.7$ mm and $z = 5.0$ mm with $V_p = 6$ kV (black curve) confirms the conclusions taken from the analysis of the temporal evolution of $T_e$. In addition, figure 11 allows to observe the direction of the electric field at each stage. Firstly, we observe the peak of $E_z$ in the direction of propagation with amplitude 14 kV cm$^{-1}$ associated to the arrival of the discharge front. Then, the return stroke propagates from the target towards the powered electrode as a wave of opposite polarity, with a second peak of $E_z$ that is also positive, as shown in Darny et al. (2017), Viegas et al. (2018a). At $z = 8.7$ mm, at only 1.3 mm from the target, the second peak of $E_z$ is very close in time to the first peak and thus is not identifiable. At $z = 5$ mm, the return stroke is associated with a peak of $E_z$ of $\sim 4$ kV cm$^{-1}$. During the rest of the pulse, the electric field in the plasma remains directed downwards with a much lower amplitude close to 1 kV cm$^{-1}$. After the fall of the applied voltage, $E_z$ reverses sign and then is directed from the charged target towards the inner grounded electrode and tends to neutralize the net charge in the plasma. The reversal of direction causes $E_z$ to pass by zero, which explains the drop of $T_e$ around $t = 1050$ ns in figure 9. Then, $E_z$ has a peak at $t = 1100$ ns in both axial positions of around $-2$ kV cm$^{-1}$, which results in a small increase in $n_e$. The similar peak of $E_z$ in both positions shows that the electrical redistribution at the end of the pulse has a diffusive character and not that of a wave, as suggested by the observations of faint emission in Yue et al (2018), Klarenhaar et al. (2018a). Its direction and value below 6 kV cm$^{-1}$ agree with the findings in Lu et al (2017).

Besides the cases already discussed, the numerical temporal evolutions of $T_e$, $n_e$ and $E_z$, are represented in figures 9–11, respectively, for three different jet configurations with $V_p = 5$ kV: free jet, jet with floating target and jet with grounded target. The results for these cases are represented only locally, and not averaged. The simulation results for floating target and $V_p = 5$ kV are very similar to those with $V_p = 6$ kV. Nevertheless, these are significantly different from the results with free jet and with grounded target. For all the cases, the peak $T_e$ and $E_z$ at the arrival of the discharge front at $z = 8.7$ mm stand between 5 and 7 eV and between 12 and 16 kV cm$^{-1}$, respectively. Then, with free jet, as there is

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**Figure 11.** Temporal evolution of the electric field at $t = 0$. On the left, at $z = 8.7$ mm. On the right, at $z = 5.0$ mm. From simulations, with $V_p = 5$ kV and three different configurations and with $V_p = 6$ kV and metallic target at floating potential. The horizontal dashed line signals $E_z = 0$. 

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Besides the cases already discussed, the numerical temporal evolutions of $T_e$, $n_e$ and $E_z$, are represented in figures 9–11, respectively, for three different jet configurations with $V_p = 5$ kV: free jet, jet with floating target and jet with grounded target. The results for these cases are represented only locally, and not averaged. The simulation results for floating target and $V_p = 5$ kV are very similar to those with $V_p = 6$ kV. Nevertheless, these are significantly different from the results with free jet and with grounded target. For all the cases, the peak $T_e$ and $E_z$ at the arrival of the discharge front at $z = 8.7$ mm stand between 5 and 7 eV and between 12 and 16 kV cm$^{-1}$, respectively. Then, with free jet, as there is
no return stroke and a quasineutral plasma is formed behind the discharge front, $T_e$ and $E_z$ decrease to low values close to zero after the propagation, in agreement with the experimental observation in Klarenaar et al. (2018a). $n_e$ in the channel is lower in free jet than with metallic targets, also in agreement with Klarenaar et al. (2018a). As figures 7 and 11 show that the peak electric field during discharge propagation is not lower in the case without target, we attribute the lower $n_e$ to the absence of the target and of the return stroke. The influence of the return stroke on the increase of $n_e$ has been demonstrated experimentally in Klarenaar et al. (2018a) and numerically in Viegas (2018), supported by the measurements of helium metastable density in Darny et al. (2017), and in Yan and Economou (2016).

Here we assess the origin of the increase of $n_e$ during the return stroke. Figure 10 shows that $n_e$ in the plume at $z = 8.7$ mm obtained from the simulations for the case with $V_p = 5$ kV and floating target reaches a value around $2 \times 10^{12}$ cm$^{-3}$ higher than in the case without target and thus without return stroke. This increase takes place mostly in the first 100 ns after the impact of the discharge on the target at around $t = 370$ ns. On the one hand, figure 5 shows that the electron-impact ionization source term $S_e$ during those 100 ns does not exceed $10^{17}$ cm$^{-3}$ s$^{-1}$ and thus cannot produce more $n_e$ than $10^{10}$ cm$^{-3}$. Even though other ionization processes take place in volume (Penning ionization, photoionization, associative ionization (Viegas 2018)), they cannot justify an increase in $n_e$ of the order of $10^{12}$ cm$^{-3}$. On the other hand, figure 8 shows that the potential of the target rises up to 1.5 kV until $t = 470$ ns, corresponding to a charging of approximately 2 nC and thus the emission of $1.2 \times 10^{10}$ electrons. If these were homogeneously distributed through diffusive and convective transport in the plasma with 3.8 cm length and 0.125 cm radius (volume 0.19 cm$^3$), an increase in $n_e$ of $6.4 \times 10^{10}$ cm$^{-3}$ (approximately one third of $2 \times 10^{12}$ cm$^{-3}$) could be expected. This distribution is not homogeneous in the whole volume and thus the electron emission has a larger impact close to the target, where $n_e$ has been assessed. These results allow to conclude that electron emission from the target and the subsequent transport of electrons in the plasma are the main source of the experimentally and numerically observed increase of $n_e$ and are important aspects of the return stroke.

With grounded target, as a sharp potential gradient remains in the plasma during the pulse, a conductive channel is formed between the electrodes (with the possibility of a transition to an arc phase on longer timescales). $T_e$ and $E_z$ remain relatively high during the pulse, close to 3.5 eV and to 2 kV cm$^{-1}$, respectively, and $n_e$ increases with time during the pulse up to $8 \times 10^{15}$ cm$^{-3}$. This increase in $n_e$ is also associated with electron emission from the target, since $S_e$ presented in figure 5 could only account for an increase in $n_e$ of the order of $10^{11}$ cm$^{-3}$ in a few hundreds ns. In agreement, Babaeva et al. (2019) have also reported higher $E_z$ and $n_e$ after discharge impact with grounded target than with floating target. This can explain the higher species production, higher $E.\ coli$ cell inactivation and lower cancer cell viability with grounded targets reported in Ji et al. (2018), Yue et al. (2018), Schweigert et al. (2019a). The difference between targets during the pulse is also visible in the experimental results of conductive current at the inner ring with $V_p = 5$ kV, presented in figure 2. Indeed, by integrating $I_{con}$ in time, we have measured 6.0 nC in the case with grounded target during the pulse, excluding the positive and negative peaks. Conversely, only 1.1 nC and 1.9 nC have been measured with free jet and with floating target, respectively. These values agree in order of magnitude with the 3.6 nC simulated at the floating target (figure 8).

Then, as the applied voltage falls, the electrical redistribution between the inner ring and the plasma affects the plasma differently in each case. As in the case with floating target with $V_p = 5$ kV, the fall of the pulse in the free jet brings a rise in $T_e$ up to 2 eV and in negative $E_z$ up to $-1$ kV cm$^{-1}$, which results in approximately constant $n_e$. The electric field in these cases is directed from the plasma at positive potential towards the inner grounded ring. Conversely, with grounded target, the electrical redistribution takes place between a grounded inner electrode and a grounded plane. For that reason, its effects are weaker than in the case of a floating target charged at 3 kV in the $V_p = 5$ kV case. Indeed, with grounded target, $T_e$ and $E_z$ decrease to very low values after the pulse and $n_e$ decreases in time. With grounded target, as $E_z \approx 0$, transport through electric drift is excluded and thus the decrease in $n_e$ after the pulse is attributed to diffusive and chemical losses. As such, we can conclude that in the other two jet configurations $n_e$ is kept constant or increases (with floating target and $V_p = 6$ kV) after the fall of the pulse due to electron emission from the inner electrode and electron transport towards the target. This analysis is reinforced by the negligible values of $S_e$ in the plasma plume after the fall of the pulse observed in figure 5 and by the decrease in electric potential of the target due to electron absorption observed in figure 8. In Kim et al. (2018) the secondary ionization at the falling edge of a pulse of applied voltage is also claimed to be induced by charge transport. These results concerning the fall of the pulse constitute a major difference between grounding and not grounding the target.

3.3. Discussion on the discrepancy of charging and uncharging the floating metallic target

The faster charging and uncharging of the floating metallic target in simulations than in experiments (figure 8) leads to questioning the conditions for comparison and the assumptions taken in the model. Firstly, both experiments and simulations have verified that changing the position of the grounded plate behind the target between $z = 15$ and $z = 31$ cm has no influence on the results. Then, we should consider that in Slikboer et al. (2019) we have used the same model with secondary electron emission ($\gamma = 0.1$) instead of a perfect electron emitter assumption to describe the interaction between the discharge and a dielectric BSO target. We have found an excellent agreement with experiments on the electric field evolution inside the target, which is closely related to surface charge, for both charging (ion neutralization and electron emission) and uncharging (electron absorption) of the target. However, taking the same secondary electron emission assumption as in Slikboer et al. (2019) for the metallic targets has a negligible effect ($\sim 0.2$ kV) on the results of figure 8, although it removes the
potential increase before the discharge impact. Furthermore, it decreases the agreement with experiments in figures 9 and 10.

The model describes metallic surfaces as perfect absorbers and perfect emitters of electrons, and thus ignores the cathodic sheath between the plasma and the metallic surface and simplifies the dynamics of charges between the plasma and the surface. In Naidis (1999) a voltage drop of 0.2–0.3 kV in the sheath between air streamers at atmospheric pressure and cathodes has been suggested. This value is too low to justify the different rate of charging in figure 8. However, a recent work (Cernak et al 2020) that highlights the importance of streamer-cathode sheaths has shown through numerical simulations a voltage drop of 1.5 kV over 50 μm when a nitrogen streamer at 26.7 kPa approaches a grounded cathode. Cernak et al (2020) also suggest that the description of streamer-cathode sheaths lies outside the conditions for validity of both the drift–diffusion approximation used in fluid models and the two-term approximation for solution of the electron Boltzmann equation.

Furthermore, the works of Bronold and co-authors (Bronold et al 2018) have initiated a microscopic description of charge transfer across plasma walls leading to the calculation of electron absorption, backscattering and secondary emission coefficients. Other works (Rees and Paillon 1997, Pechereau et al 2016, Babaeva et al 2019) simulate the fluxes of electron emission from metallic surfaces from ion bombardment, thermionic emission, field emission or photo-emission processes. These factors point to potential improvements of the model that could lead to better agreement between numerical and experimental results in figure 8.

Finally, in the experiment a thin oxide layer is very likely to be formed on the copper surface interacting with the plasma and could be responsible for diminishing the conductivity of the target (Altieri et al 2017). As the model supposes the conductivity to be infinite, this could justify the slower rise and fall of electric potential in the experiments.

4. Conclusions

This work has addressed the interaction of kHz μs-pulsed atmospheric pressure He jets with metallic targets through simulations and experiments, focusing on the differences between floating and grounded targets. Three jet configurations have been studied with positive polarity of applied voltage: free jet, jet with metallic target at floating potential and jet with grounded metallic target. The same conditions have been taken in experiments and simulations. Experimentally, the jets have been studied through imaging and Stark polarization spectroscopy peak electric field measurements. In the case with floating copper target, electron properties in the plasma plume have been assessed through Thomson scattering measurements and the temporal evolution of electric potential of the target under plasma exposure has been measured with a high voltage probe. Numerically, an axisymmetric two-dimensional plasma fluid model has been used. A description of the floating metallic plate as an isopotential infinitely conductive surface where ions are neutralized and electrons are emitted and absorbed through effect of the electric field has been proposed.

Experiments and simulations have observed the same discharge dynamics. The discharge propagates faster towards the grounded target than in the other two configurations. With floating target, the discharge only propagates faster with respect to the free jet case in the last 5 mm of propagation. Moreover, experimental and numerical results both show that the peak electric field at the discharge front during the propagation is approximately the same between the three different configurations. With both grounded and floating targets, a return stroke has been observed after discharge impact on the target, as an ionization wave propagating from the target towards the powered electrode in an already ionized channel. With grounded target, reactivity stays in the plasma plume during the 1 μs pulse, while with floating target it severely decreases in a few hundred ns. At the fall of the applied voltage pulse, another electrical redistribution takes place between the now grounded inner electrode and the positive plasma. This has been shown to have higher intensity with floating target than with grounded target.

The explanations for the differences between grounded and floating targets have been found in the temporal evolution of electric potential of the floating target. The discrepancy between simulations and experiments in that temporal evolution has been discussed, taking into account that the model describes metallic surfaces as perfect absorbers and emitters of electrons. A more accurate description of electron absorption, backscattering and emission from the surfaces has been pointed as a potential future improvement of the model. However, both experiments and simulations have shown that the potential of the floating target after discharge impact increases a few kV per μs, depending on the amplitude of applied voltage. Thus, the pulse width and the applied voltage allow to control the charging of the target. After the pulse, the potential decreases at approximately those rates, until approaching zero. As such, during dozens of ns after the impact, the target is close to grounded but, at the end of the 1 μs pulse, it is at a high voltage, potentially more than half of the applied voltage. That explains the similar return stroke with floating and grounded targets. Furthermore, it justifies the decay in reactivity during the pulse with floating target as the target charges and potential gradients in the plasma dissipate. As a result, simulations have shown that the electron temperature and electric field remain high in the plasma with grounded target and the electron density in the plasma plume increases during the pulse with grounded target but not with floating target or without target.

Finally, the charging of the floating target has shown that the redistribution at the end of the pulse takes place between a grounded inner electrode and a plasma limited by a charged surface in the floating target case, while with grounded target it takes place between two grounded electrodes. That justifies the stronger intensity of that redistribution with floating target. Experiments and simulations have shown an increase in electron temperature, magnitude of electric field and electron density in the plume with floating target after the pulse, which is not the case with grounded target. The increases in electron density in the plume after the pulse with floating target, during the return stroke with both floating and grounded targets and
during the whole pulse with grounded target, have been shown to be mostly due to electron emission from metallic surfaces and charge transport in the plasma.

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