DROPLET-BASED LASER-PRODUCED PLASMA SOURCES – PHOTON EMISSION, MATTER EXPANSION AND MITIGATION

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Non fuerat nasi nisi ad has scientiae artes – harum palmarum feritis

Impetu, Exemplo, Constantia, Inventione, Experimento, Acumine, Meditatione, Labore,
Intuendo, Disciplina, Conoscendo, Libertate, Indole, Audacia, Numine, Cura

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Abstract

The field of this work is related to laser-produced plasmas generated with regenerative targets in form of droplets. Applications of laser-produced plasmas are from radiation sources to fusion with inertial confinement. In the experiments the energy flux delivered to the target is compatible with radiation sources in the extreme ultraviolet (EUV) range. Commercially, the most promising application is in extreme ultraviolet lithography. In laser-produced plasma sources, together with radiation, matter expands eventually up to functional components of the source, such as collector mirror in an EUV source. Matter (debris) coming up to collector mirrors has detrimental effects on its reflectivity and on the collector mirror lifetime. EUV sources require continuous operation at an affordable cost of ownership. Therefore, the study of the EUV emission, matter expansion and debris mitigation with droplet targets is necessary, whereas the use of droplets is fundamental to achieve continuous operation.

The emission of EUV radiation in a EUV +/-2% bandwidth centered at 13.5 nm (in-band) is measured around the droplet targets at angles from the laser axis up to 120°. The effects of irradiance, droplet diameter and ambient pressures on the in-band angular emission are addressed experimentally. Analytical models representing the physical processes are developed and calibrated with experimental data. The results of the experiments are explained with the calibrated models and physical characteristics of plasma dynamics are inferred. EUV emission up to 150° from the laser axis is extrapolated from the model. Larger droplets lead to a faster decay of the EUV emission at larger angles from the laser axis. A decrease in the irradiance has an opposite effect. And, the models are used to estimate EUV emission distribution for a desired droplet diameter and laser irradiance.

Matter expansion from droplet-based laser-produced plasmas is measured at distances from the irradiation site, which are characteristic for the location of collector mirrors in an EUV source. Diagnostics are developed to measure the overall ion current and to distinguish between the ion charge states in the presence of ambient gas and at different angles from the laser axis. A model representing the dominant physical phenomena in the three-
dimensional expansion is developed, calibrated with experiments, and successfully used for mean charge state angular and spatial distribution. The population of doubly charged ions is peaked towards the laser; instead the singly charged ions are slower and more isotropic. During irradiation with offset between droplet and laser focal spot, the main direction of propagation of ions correlates with the one of EUV emission due to opacity effects. The distribution of fragments ejected by the irradiated droplet targets is measured at different angular positions, ambient pressures, and droplet diameters. The fragments load strongly increases after 105° from the laser axis mainly due to recoil momentum, it increases with droplet diameter and ambient pressure.

A system to mitigate matter expanding up to the collector mirror is developed. The mitigation system bases on deflecting the particles from the laser-produced plasma with a high-momentum gas flow, which is accelerated through an aerospike nozzle up to rarefied conditions. Experimental, computational and analytical studies are performed to optimize the aerospike design in rarefied operating condition. Kinetic boundary layer formation on the surface of the spike lifts the flow during acceleration. The optimization of the geometry leads to an order of magnitude reduction of the jet-lifting angle. Finally, the mitigation system is tested with droplet-based laser-produced plasma, and deflection capabilities are demonstrated. Fragments are deflected by the high-momentum gas flow, and increase of stagnation pressure and stagnation temperature leads to a larger deflection.
Sommario

Il campo di questo lavoro è legato a plasmi prodotti da laser generati colpendo bersagli rigenerativi in forma di gocce. Applicazioni di plasmi prodotti da laser spaziano da sorgenti di radiazione a fusione a confinamento inerziale. Negli esperimenti il flusso di energia che arriva al bersaglio è compatibile con sorgenti di radiazione nell’ultravioletto estremo (EUV). Commercialmente, l’applicazione più promettente è nella litografia EUV. In sorgenti basate su plasmi prodotti da laser, materia del bersaglio si espande fino a elementi funzionali della sorgente, come lo specchio collettore in una sorgente EUV. La materia (debris) che arriva fino allo specchio collettore va a deteriorarne la riflettività, e quindi compromette anche la vita dello specchio collettore nella sorgente. Sorgenti EUV devono operare in continuo e avere un costo di proprietà fattibile. Per questi motivi, lo studio dell’emissione EUV, dell’espansione di materia, e dell’attenuazione del debris sono necessari usando gocce come bersagli, il cui utilizzo è fondamentale per garantire l’operazione in continuo.

L’emissione di radiazione EUV in una larghezza di banda di EUV +/-2% centrata a 13.5 nm (in banda) è misurata attorno alle gocce bersaglio ad angoli rispetto all’asse del laser fino a 120°. L’effetto dell’irradianza, del diametro delle gocce e della pressione ambiente sull’emissione in banda sono indirizzate sperimentalmente. Modelli analitici che rappresentano i processi fisici sono sviluppati e calibrati con dati sperimentali. I risultati degli esperimenti sono spiegati tramite i modelli calibrati, e sono dedotte caratteristiche fisiche della dinamica del plasma. Emissione EUV è estrapolata dai modelli fino a 150° dall’asse del laser. Gocce più grandi portano a una decrescita più importante della radiazione EUV con l’angolo rispetto all’asse del laser. Una decrescita dell’irradianza ha un effetto opposto. Inoltre, i modelli sono usati per stimare l’emissione EUV per un dato diametro della goccia e una data irradianza.

L’espansione di materia dal plasma prodotto da laser basato su gocce come bersaglio è misurata a distanze dal punto di irradiazione che sono caratteristiche della posizione di specchi collettori in una sorgente di radiazione EUV. Diagnostiche sono sviluppate per misurare la corrente ionica totale e per distinguere tra i diversi stati di carica in presenza di
gas ambiente ed a diversi angoli rispetto all’asse del laser. Un modello è sviluppato che rappresenta i fenomeni fisici dominanti nell’espansione tridimensionale. Il modello è calibrato con esperimenti, ed è usato per stimare la distribuzione angolare e spaziale dello stato di carica del plasma. La popolazione di ioni con carica doppia è preponderante verso il laser, invece ioni con carica singola sono più lenti e più isotropici.

Durante l’irraggiamento con disallineamento tra goccia e fuoco del laser, la direzione di propagazione principale degli ioni correla con quella dell’emissione EUV dovuto a effetti di opacità. La distribuzione di frammenti espulsì dalle gocce irradiate è misurata a diverse posizioni angolari, diverse pressioni ambiente, e diversi diametri della goccia. Il carico di frammenti aumenta in modo importante dopo 105° dall’asse del laser dovuto principalmente alla pressione di rinculo del plasma sulla goccia, e lo stesso carico aumenta con il diametro della goccia e con la pressione ambiente.

Un sistema è sviluppato per attenuare la materia in espansione fino allo specchio collettore. Il sistema di attenuazione è basato sulla deflessione delle particelle espulse dal plasma prodotto dal laser con un getto di gas ad alto momento. Il getto di gas è accelerato tramite un ugello aerospike fino a condizioni rarefatte. Studi sperimentali, computazionali, e analitici sono eseguiti per optimizzare il design dell’aerospike in condizioni operative rarefatte. La formazione di uno strato limite cinetico sulla superficie dello spike alza il getto durante l’accelerazione. L’ottimizzazione della geometria porta a un ordine di grandezza di riduzione dell’angolo d’innalzamento del getto. Infine, il sistema di attenuazione è testato con il plasma prodotto da laser basato su gocce come bersaglio, e la capacità di deflessione è dimostrata. I frammenti sono deflessi dal getto di gas ad alto momento, e un aumento della pressione di stagnazione e della temperatura di stagnazione porta a un aumento della deflessione.
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At the beginning of my PhD the facility consisted of a vacuum chamber and a 10 Hz laser that was shooting on metallic plates. At the end of the PhD I am leaving behind me 4 facilities, in particular a laser-produced plasma facility based on droplet targets and working at 6 kHz, and a facility to study and enhance discharge-produced plasma based on a magneto-plasma compressor. The opportunity to develop these working facilities has been an invaluable experience, and I am deeply grateful to Prof. Reza S. Abhari for entrusting me these projects and for leaving freedom to my creativity, which translated in investing over a million of Swiss francs. Furthermore, I would like to acknowledge Prof. Abhari for the time that he dedicated to me. Thanks to the many meetings I could greatly improve, especially in the method to approach problems, technically, and in numerous soft skills as the way to socially interface with colleagues and associates.

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Chapter 1

INTRODUCTION

The field of this work is the study of laser-produced plasmas (LPPs) generated with droplet targets. Depending on the laser system, droplet-based LPPs find applications from radiation sources to inertial confinement fusion. The use of droplets is fundamental in processes in order to ensure continuous operation.

1.1 Laser-produced Plasmas and Main Applications

1.1.1 Laser-produced Plasmas

Laser plasmas are experimentally studied from the early 1960s when the first high-power lasers were created. With the first lasers, plasmas could be obtained with peak laser intensity in the order of $10^9 \text{ W/cm}^2$.\(^1\) The progress in laser technology in the last decades led to an increase of the peak intensity up to $10^{22} \text{ W/cm}^2$.\(^2\) A number of phenomena are possible during laser irradiation including heating, melting, evaporation of condensed matter, thermoionic emission, change of target reflectivity during irradiation, thermal and hydrodynamic instabilities, transfer of recoil momentum to the target, multiphoton ionization, high-harmonic generation, collisional absorption of laser radiation, emission of radiation up to $\gamma$-ray with the highest laser intensities, electron and ion acceleration, generation of strong magnetic fields, and fusion. Depending on the intensity of the laser different phenomena become dominant.
In a moderate intensity regime ($10^9$-$10^{13}$ W/cm$^2$) with laser photon energies lower than the first ionization potential of the target and with durations in the order of nanoseconds, the main processes that take place in few hundreds of micrometers are the following. In the first tens of picoseconds of irradiation fast electrons are emitted by the target and ablation and ionization begins. In the first phase ionization can take place only with multiple photons as one is not enough to ionize an atom of the target, and the process is defined as multiphoton ionization. The target supplies mass to the plasma that forms on its surface. At the same time, as mass leaves the surface of the material, a high-pressure region forms on the surface and a recoil pressure begins to act on the target surface. Then, as the flux of laser energy increases further, plasma expands, and shields the surface of the target from laser photons becoming the main absorber of laser energy through collisional absorption. The energy that the laser electromagnetic wave pumps into the plasma is reemitted in form of radiation, increases the internal energy, leads to further evaporation and ionization of the target material, and is used in the expansion. Until the source of mass (target) and/or the source of energy (laser pulse) are not depleted, the process continues in a quasi-stationary manner. The density of the plasma formed above the surface of the target is close to that of a typical solid and its temperature of some millions of kelvin. Condition favorable for emission of radiation down to soft-x-ray. Typical expansion velocities of ions from high-Z materials are in the order of tens of kilometers per seconds. Plasma expands from the region where the electromagnetic wave is absorbed and cools down to temperature where recombinations with electrons become dominant over ionizations and the plasma loses part or all of its charge. Then, as density in the three-dimensional expansion rapidly decreases with the third power of the distance from the irradiation site, plasma becomes rarefied and ions, atoms and electrons travel further until they reach a surface. In the meantime on the surface of the target the important pressure (order of GPa) and the high temperature (that might reach the thermodynamic critical point) cause the ejection of fragments, which expands at velocities orders of magnitude smaller than ions.
1.1.2 Radiation Sources
Depending on wavelength, power and brightness of the radiation that is generated, the application of the source differs. Plasma-based radiation sources are employed mainly at wavelengths between ultraviolet and x-ray, where radiation cannot anymore be generated by common lamps, but rather by lasers, plasma sources or synchrotrons. Main applications of LPP as radiation source follows.

Lithography
Nowadays, transistors are made by photolithography with immersion and multiple patterning. Semiconductor manufacture strives to reduce the size of the features on integrated circuits (IC), which requires a reduction of wavelength of the radiation used during projection down to extreme ultraviolet and soft x-rays.\(^6\) The minimum achievable line width on a wafer \(L_w\) is proportional to the wavelength \(\lambda_p\) of the radiation during projection:\(^7\)

\[
L_w = k_1 \frac{\lambda_p}{NA}.
\] \(\text{(1)}\)

\(k_1\) is a constant that depends on the optical system and on the lithographic process, and \(NA\) is the numerical aperture seen at the wafer. As today 193 nm photons from ArF excimer lasers are used, the transition to extreme ultraviolet (EUV), would greatly reduce the minimum line width writable on a wafer.

Photons with the shortest wavelength produced by a commercial laser have 157 nm, and come from \(F_2\) excimer lasers.\(^8\) However, \(F_2\) lasers application as radiation source in lithography was dropped, the main reason was for the too large difficulties in the mirror manufacture compared to the gain of using the shorter wavelength photons.\(^9\) The technology used for excimer lasers cannot be used to produce more energetic photons in the EUV range; this requires plasmas as lasing medium, which leads to a x-ray laser.\(^10,11\) The solution is suitable for processes were a low repetition rate is needed (in the order of Hz). But for industrial processes as lithography, at the present the use of plasma as a lasing medium in the source does not allow achieving the power required.\(^12,13\) Another alternative that was investigated is the use of synchrotron radiation,\(^14\) the disadvantage of the latter is
the cost-of-ownership, which is at the present too high for the industrial process. Therefore, spectrally-tailored plasma sources are the main viable solution for the near future, in order to reduce $\lambda_p$. And, EUV lithography (EUVL) based on a plasma source would ensure the further miniaturization of integrated circuits needed by the semiconductor industry (see Section 1.2).

Alternative technologies for $L_w$ reduction in lithography are nanoimprint, direct self-assembly and maskless lithography, which includes e-beam and focused ion beam lithography.

**Microscopy**

Microscopy is a technical field that in order to resolve smaller features needs illumination of the sample with radiation at shorter wavelengths. Of particular interest is soft x-ray microscopy in the water window (between 2.34 nm and 4.37 nm), which complement the capabilities of visible light and electron microscopy.\(^7\,15\) The advantage of imaging in the water window is that organic materials show a strong absorption at these wavelengths, as opposed to water, which is relatively non-absorbing. Therefore, this allows the imaging of carbon-based materials in aqueous environment. Most of the soft x-ray microscopes are based on synchrotron as radiation source, but effort has been made to develop a compact soft x-ray source based on laser-produced plasma from a jet target.\(^16\,17\) The source was developed to work with liquid nitrogen, which is dispensed from a cryogenically cooled fused silica capillary nozzle with a 15 $\mu$m – 20 $\mu$m orifice.\(^15\) The jet of liquid nitrogen is generated with a pressure of around 60 bar and exits from the orifice with a velocity of around 100 m/s.\(^18\) The jet is irradiated with a Nd:YAG laser working at 100 Hz at laser intensities around $4 \times 10^{13}$ W/cm\(^2\).\(^15\)

**1.1.3 Inertial Confinement Fusion**

Inertial confinement fusion (ICF) is one of the two main approaches that are currently investigated for the development of a clean source of energy (the other approach is based on magnetic confinement). The main challenge in fusion is to maintain high-density and
high-temperature plasma for a long enough time to allow fusion to occur (Lawson criterion). In the sun the condition are satisfied as the gravitational compression is huge, and indeed the largest fusion power density in the sun is in the center, within 30% of its radius. But on earth, this is not the case. ICF fights against the natural mechanism of release of energy through expansion in the three-dimension with an inertial implosion were a multitude of lasers irradiate a target generating pressure load on its surface.

It was recently (year 2014) demonstrated that the ICF is able to generate more energy during the deuterium-tritium implosion than the amount of energy deposited into the target. In the experiment 192 lasers delivered up to 1.9 MJ into a gold hohlraum of 9.425 mm in length and 5.75 mm in diameter where the deuterium-tritium capsule was contained. The main progress that enabled the last achievement was a “high-foot” implosion scheme. This was one of the last improvement of a large series from 1972 when the first work on ICF was officially published (before ICF was under military classification).

1.1.4 Other Applications of Laser-produced Plasmas

LPP were proposed as source of ions for particle accelerators in 1969 for the first time. From then laser ion sources found applications in different fields including proton imaging/radiography/deflectometry, radioactive ion beam facilities, and possibly future cancer therapy machines.

Spectroscopy

Radiation and ions emitted by the target irradiated with the laser can be analyzed with spectrometers and mass spectrometer to gather information on the material. Laser-induced breakdown spectroscopy has been under development almost since the advent of the laser, and portable LIBS became a reality since the late 1990s. The maturity of the measurement techniques allowed the use of a LIBS system on Mars; indeed it is an element of the ChemCam on board of the Mars science laboratory rover. And, laser induced mass
spectroscopy has made possible space-resolved elemental and isotopic analysis of sample matrix.\textsuperscript{29}

**Pulsed Laser Deposition**

Pulsed laser deposition (PLD) is based on the ablation of a target material that has to be uniformly deposited on a substrate in form of thin film. Normally, in PLD the laser system irradiates a slab target, which ensures flexibility to different target materials with melting temperature that would be challenging for a droplet dispenser. PLD is mainly applied to the growth of crystalline oxides, in the synthesis of epitaxial heterostructures and superlattices, and in the growth of high-temperature superconducting thin films.\textsuperscript{30} These processes enable the production of sensors from photon to gas detection, protective and biocompatible coatings, and nanomaterials as nanocrystals, nanowires and nanocomposites.

**Laser Propulsion**

The large pressure force that the laser plasma exerts on the ablated surface is considered as possible propulsion mechanism for space flight, and to deorbit debris in the space immediately around earth.\textsuperscript{31} Different forms of laser-plasma-based propulsion are being investigated. The main advantage of these systems is that the energy source (laser) is decoupled from the momentum source on the spacecraft. This reduces the overall mass of the spacecraft or allows a larger payload. The disadvantage is the need of line-of-sight between the remote laser system and the spacecraft. This becomes particularly a problem if the laser system is earth-based, as ablation of airliner and other possible flying objects/living creatures should be avoided.

**1.1.5 Impact**

Commercially, the most important application for droplet-based LPP is as possible radiation source in the soft x-rays to ultraviolet range for semiconductor manufacture. In 2013, the semiconductor industry exceeded $300 billion in sales. The largest semiconductor category by sales was logic with $85.9 billion followed by memory and
MOS micro-ICs with $67 billion and $58.7 billion, respectively.\textsuperscript{32} Within these categories, the majority (62\% estimated in 2012) of the chips produced at a 45 nm or lower node width are used in mobile-devices.\textsuperscript{33} Semiconductor industry has been driven by what is known as the Moore’s law, which states that the number of transistors placed on integrated circuits double approximately every two years.\textsuperscript{34} Moore’s law has led to continuously decreasing prices of semiconductors in the last four decades, and it was possible because of continuous investments in the research and development of new manufacturing technologies. As in a company the investment in research scales with the volume of production, larger companies have been able to develop further than smaller ones, and occupy and remain in leading-edge position. To continue on this path, EUV lithography is one of the main possible technology to grant the further feature size reduction, however it is related with significant investments. In order to maintain the leading-edge position and to generate demand for products with an always better performance, smaller size and lower cost, the main industrial players are interested in continuing with the technological progress. But, the rate could slow down due to large necessary financial investments. A possible scenario is a further performance improvement through multiple patterning until EUV lithography demonstrates solid performance and enters the market, or continuing the hybrid use of both EUV and conventional scanners. Seen by the producers of semiconductors, the pressure to decrease further the feature size could lead to a continuous increase in production costs and reduction of margins. This could push more and more companies to become fabless. Leading to few companies as semiconductor producers with leading-edge technology that would have the advantage of economy of scale. Eventually, the next generation lithographic tools will be used, but its initial costs and cost of ownership could be too high to continue the trend of cost improvements predicted by the Moore’s law. The consequence could be a deceleration in the demand for leading-edge semiconductor devices.
1.2 EUV Lithography

A schematic of an EUVL scanner for high volume manufacturing (HVM) is shown in Fig. 1. EUV radiation is generated and collected by the source and conditioned in the illuminator up to the reticle stage where the EUV mask is located. The latter is made of a reflective (to EUV) part and of an absorber, which defines the features to be written on the wafer. After reflection onto the EUV mask, radiation enters the reduction optics, which bring the beam up to the wafer. A 4x reduction and a $NA = 0.33$ are used in the state-of-art scanner. After reflection onto the EUV mask, radiation enters the reduction optics, which bring the beam up to the wafer. A $4x$ reduction and a $NA = 0.33$ are used in the state-of-art scanner.\(^\text{35}\) Once the EUV radiation arrive on the wafer, the resist that covers the surface absorbs the radiation and emits photoelectrons and secondary electrons, which initiate locally chemical reactions printing the feature that were on the EUV mask.

![Overview of ASML’s TWINSCAN NXE:3300B.](image)

Fig. 1: **Overview of ASML’s TWINSCAN NXE:3300B.** Radiation is generated by the EUV source, is conditioned by the illuminator up to the reticle stage and reflected by the projection optics onto the wafer.

In the lithographic process the EUV mask is fundamental, since it contains the information that has to be replicated on hundreds of wafers per hour. Therefore, it is critical to inspect the mask with EUV radiation (actinic) in order to see possible defects as wafers would. Tools for mask inspection and metrology are actinic blank inspection to check quality of multilayer stack (A-BI), actinic pattern inspection to detect minimum printing particles (A-PI), and aerial image measurement system to determine defect printability (AIMS).
1.3 EUV Sources for Lithography

EUV sources are needed in both high volume manufacturing (HVM), metrology and inspection. Droplet-based LPP is used in both, but the EUV sources are optimized differently as the requirements are not the same.

1.3.1 General Characteristics of an EUV Source

In order to bring the radiation generated by the EUV source up to the wafer, Bragg’s reflectors have to be used for normal reflection, and the whole machine has to be in vacuum. Otherwise, EUV photons would be completely absorbed by conventional mirrors and by the ambient gas. Bragg’s reflectors are made of multiple layers of materials, which are deposited on a super polished surface one after the other with a thickness per layer in the order of nanometers. And, in the case of aspherical optics the thickness of the layers depends on the radial position on the surface of the mirror, as the angle of incidence (AOI) on the illuminated surface changes.

In the EUV range the most efficient reflector is the molybdenum/silicon multilayer (ML), which has the highest reflection of around 70% at a wavelength of 13.5 nm. For a normal incidence ML mirror the optimum thickness for the Si/Mo layers ranges between 2.8 nm and 4.5 nm. Considering that more than ten reflective surfaces are needed in an EUV scanner, the highest reflection of the ML mirrors sets the wavelength at which the EUV source has to operate, i.e. 13.5 nm.

Tin, xenon, and lithium are among the target materials to be excited to plasma and generate the largest amount of EUV energy centered at 13.5 nm per consumed laser energy, which is defined as conversion efficiency (CE). The main factors for the choice of the target material are the reliability in the regenerative target generation, the amount of debris produced with EUV emission, and the conversion efficiency (CE). Tin is the best material for target generation and CE, and is the target material used in this work.

The concept of debris describes matter expanding in the vacuum chamber of the EUV source. Debris is composed by different kind of particles: ions, neutrals, and fragments of the target. The challenge related to debris is that it expands up to the collector mirror surface degrading its reflectivity with time, which has a direct impact on the cost-of-
ownership of the scanner as ML collector mirrors are one of its most expensive components.

The main components of a source are laser system, droplet dispenser, collector module, and debris mitigation system. The laser system is highly dependent on the application, as output power and wavelength varies with the type of laser. Different strategies have been proposed to increase the CE through modification of the laser system, as multiple laser beams, a hollow laser beam and pre-pulsing. An increase of around 10% was reported for a three beams configuration using Nd:YAG lasers working at the fundamental frequency and at the optimum angle between the beams of 30°. The hollow laser beam configuration affects the plasma expansion and increases the CE up to around 17%. And, pre-pulsing was found to be the most promising way to increase the CE from CO$_2$ lasers, and decrease the ion kinetic energy.

The droplet dispenser has stringent requirements in both the applications, because the stability of the source directly depends on the accuracy of delivering the droplet targets at the irradiation site when the laser fires.

The collector module design depends on the application –HVM, metrology or inspection– and in particular on the amount of power that the EUV source is expected to generate. Because of the cost of ML mirrors the size of the collector mirror should be minimized to fulfill the required EUV power, and size and position in the vacuum vessel should match the extraction solid angle.

Debris mitigation is fundamental in a source with viable cost of ownership. Different strategies were developed based on magnetic field, electric field, and inertial collisions. The main difference is that electric and magnetic field acts only on electrically charged particles, instead inertial collisions acts on all the components of debris. Another way to cope with the problem of debris on collector mirrors is cleaning, which demonstrated to be successful on HVM sources. The latter is based on the reaction between tin and the hydrogen ambient gas, which produces stannane (SnH$_4$) that is volatile and detach from the collector surface leaving the coating undamaged. However, as the compound is highly toxic and tin is not only present on the collector surface, the complexity of the pumping system and of the ambient gas delivery increases.
1.3.2 Sources for HVM

HVM sources are expected to deliver 250 W of EUV for insertion in the market of EUVL in 2015.\textsuperscript{52} In order to achieve the required power with typical conversion efficiencies in the order of 1%, the laser system has to be able to deliver tens of kW. This sets the choice of the laser system: CO\textsubscript{2} lasers.\textsuperscript{53} The droplet dispenser should deliver at least a droplet per laser pulse. Typical operating frequency of the droplet dispenser is around 50 kHz.\textsuperscript{51} Another aspect to maximize in order to collect as many EUV photons as possible is the collector numerical aperture. For this reason in HVM the collector module covers up to 5 sr,\textsuperscript{51} and the debris mitigation strategy has to be chosen accordingly.

Two types of collector mirrors were originally used in demonstration tools, i.e. normal incidence and grazing incidence collectors. The first are ML mirrors as the ones used in the illuminator and as reduction optics, and the second reflects the EUV photons at large angle of incidence and are made of multiple shells.\textsuperscript{54} Grazing incidence collectors were mainly used in combination with discharge-produced plasma (DPP), and a rotating foil trap as debris mitigation mechanism.\textsuperscript{55} However, DPP has the disadvantages of having a large source size and of difficult scalability to higher power, and LPP with normal incidence collectors remained as the mostly investigated source concept for HVM.

Spectral purity of the radiation exiting the source is important for the use of the source in the lithographic process, as radiation from other wavelength arriving at the wafer has detrimental effects on the process. The reason is that the resist is sensitive also to out-of-band radiation (not within the 2% bandwidth centered at 13.5 nm),\textsuperscript{56} and a further thermal load is added on the wafer leading to thermal expansion, whereas the position of the wafer has to be controlled within nanometers. In order to limit out-of-band radiation, modification of the collector substrate and filters have been developed.\textsuperscript{57,58}
1.3 EUV Sources for Lithography

![Diagram of EUV source](image)

**Fig. 2:** *Schematic of the main components of a high volume manufacturing EUV source. EUV is generated by the interaction between droplets and laser, and is focused by the ML collector mirror to the intermediate focus.*

A general schematic of a LPP-based HVM source is represented in Fig. 2. The dispenser delivers droplets up to the irradiation site, where the CO\textsubscript{2} laser is fired. Plasma forms and EUV radiation is emitted. The latter is collected by a ML mirror with ellipsoidal geometry and focused at the intermediate focus. Position and shape of the ML collector mirror is such that one focus corresponds to the irradiation site, and the second to the intermediate focus.

### 1.3.3 Sources for Metrology and Inspection

For metrology applications brightness as opposed to power is the challenge. Brightness in excess of 100 W/(mm\textsuperscript{2}sr) is required.\textsuperscript{59} As brightness is inversely proportional to the source size squared, a small laser focal spot is advantageous. For this reason and in order to have a compact source Nd:YAG lasers are normally used.\textsuperscript{60}
Fig. 3: Schematic of the configuration of an EUV source for metrology and inspection. Droplet targets are delivered from the top (perpendicular to the drawing).

Fig. 3 shows a schematic of a source for metrology and inspection. The collector system is not anymore axisymmetric with respect to the laser axis, and EUV is captured with a smaller solid angle. Droplets are dispensed at the irradiation site, and irradiated by the Nd:YAG laser. The generated EUV is then collected by a first planar mirror and a second spherical mirror, which focus the radiation up to the intermediate focus. Another variant is to have just one collector mirror, which is aspherical and off-axis. This solution would minimize the aberration of the spot at the intermediate focus, but increase the cost of ownership. State-of-art LPP-based sources for metrology and inspection are capable of brightness up to 259 W/(mm²sr).

1.4 Motivation

Even though the use of a regenerative target is unavoidable for a commercial EUV source, LPPs were mainly studied with slab targets. The synergistic study of EUV emission and matter (debris) expansion is needed for the design of a source with the minimum cost of ownership. Furthermore, the development of a mitigation system from a droplet-based LPP is necessary for application of the EUV source to usable machines.

The influence of the target geometry on the EUV emission in the three-dimensions around the droplet target is essential for the location of the ML collector mirror. And, the dependence of the three-dimensional EUV emission on operating parameters as irradiance...
1.5 Literature Review

of the laser system, droplet diameter, and ambient pressure in the vacuum vessel allows to set of the optimum source operating condition.

On the other hand, also debris has to be estimated around the droplet target, because it has detrimental effects on the reflectivity of the ML collector mirror, which is one of the most expensive items in an EUV source (more than 100 kEUR for a state-of-art ML mirror with a diameter over 200 mm). Indeed to find the optimum design for a EUV source, compromises have to be made between highest possible EUV emission and lowest possible debris load.

Finally, as no location in the vacuum vessel around the droplet target is debris-free, a mitigation strategy that applies to all the components of debris has to be developed and tested. The mitigation system has, therefore to act on electrically charged as well as on neutral particles, it has to be compatible with the other components of the EUV source, and it has to operate in vacuum.

The knowledge of these three aspects – EUV emission, matter expansion, and mitigation – with droplet targets is fundamental for a successful EUV source design, and it is a novel contribution to the scientific community.

1.5 Literature Review

Three main research areas are explored in this work: EUV emission from droplet-based LPP, matter expansion from droplet-based LPP, and flow expansion to rarefied conditions. The latter topic is studied because it is necessary for the optimization of the debris mitigation system, where a nozzle accelerates a gas flow to high velocity, in order to deflect expanding matter away from functional surfaces as collector optics.

EUV Emission

EUV emission from LPP was studied extensively, in order to select the best target material to emit in the EUV +/-2% bandwidth centered at 13.5 nm. Tin, xenon, and lithium were studied in detail, and tin produces the highest conversion efficiency. In a tin plasma EUV is primarily emitted by the unresolved transition array of Sn$^8$ to Sn$^{13+}$ at an
optimum temperature of the laser-produced plasma of 30 eV – 40 eV and an electron number density around $10^{19}$ cm$^{-3}$.\cite{48,71}

The influence of laser wavelength and irradiance on the EUV emission from tin targets was also explored. A laser wavelength of 10.6 μm (CO$_2$ laser) generates a maximum conversion efficiency of around 3% – 4% at laser irradiances from $10^9$ to $10^{10}$ W/cm$^2$, instead at a shorter wavelength of 1.064 μm (Nd:YAG laser) the conversion efficiency decreases to around 2% at laser irradiances of few $10^{11}$ W/cm$^2$.\cite{70,72-78}

The angular emission of EUV radiation from tin-based LPPs was extensively studied following two approaches, for both of which slab targets were used. In the first approach, the laser and the EUV energy monitor remained stationary and the slab target was rotated.\cite{79,80} In the second approach, the laser and slab target remained stationary and the EUV energy monitor moved.\cite{72,81,82} The EUV emission around the target was assessed using different experimental and/or theoretical approaches, including $\cos^9(\theta)$ fit on the measured distribution\cite{72,80,81} (where $\theta$ is the angle between the laser axis and the plasma observation direction and $\alpha$ is a constant), a model with the absorption of EUV from a 1D plasma layer,\cite{82} and a model with the absorption of EUV from a plasma layer around a rectangular region that is assumed to emit EUV.\cite{79}

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**Matter Expansion**

The angular and kinetic energy distribution of ions from LPP generated in high vacuum was studied extensively with slab targets of tin,\cite{83} and of other materials.\cite{84-88} The effects of the wavelength on the plasma dynamics was studied with tin slab targets mainly comparing Nd:YAG and CO$_2$ lasers.\cite{73,89} The influence of the background gas was considered with slab targets out of silver,\cite{90} carbon,\cite{91} silicon,\cite{92} and tin.\cite{93,94} LPP from droplet targets was studied experimentally mainly in the firsts nanoseconds, during EUV emission. The influence of the ratio between droplet diameter and laser focal spot size (only for droplets larger than the focal spot size) on the plasma density and on the location of EUV emission was addressed by Yuspeh *et al.*\cite{95}. The spatial and temporal characteristic of the plasma during the laser pulse was characterized with 280 μm diameter
droplets, and down to 50 μm diameter droplets with a laser spot size of 26 μm. And, the ionic population was measured in high vacuum irradiating tin coated micro-balloons of 500 μm in diameter at 180 mm from the irradiation site. The expansion of atoms from LPP was studied experimentally with slab targets in terms of angular distribution, kinetic energy, and of ambient pressure influence. Fragments generated by laser plasma were observed and measured with droplet targets at the irradiation point, and at distances representative for collector mirrors only with slab targets.

**Debris Mitigation System**

A debris mitigation system based on inertial collisions was developed in this work, and consisted in the supply of a high-momentum layer of particles on the surface/aperture to be protected. An aerospike nozzle was chosen to deliver the two dimensional high-momentum layer. Measurements for linear aerospikes are available mainly for continuum or close to continuum cases, for application to engines of reusable space shuttles, with maximum Knudsen numbers smaller than 0.05. However, for small sized nozzles operating at low pressures, as in the applications treated in this work, the Knudsen number is larger, reaching the transition regime (0.1 < Kn < 10) in the divergent part of the nozzle. The simulations performed on linear aerospikes in the transition regime are focused on effects on the thrust and on the specific impulse, and not directly on the characterization of the non-ideal phenomena developing downstream of the throat. Analytically, literature is available on the solution to the Kramers problem, which is similar to the flow expanding along the surface of the spike. The Kramers problem consists of a gas filling the space above a physical wall, which flows with a direction parallel to the plane and with a non-zero velocity gradient along the component perpendicular to the wall. Experimentally, this scenario was measured by Reynolds et al. The difference between the Kramers problem and the scenario of this thesis is that macroscopically the spike is not a flat wall, but rather a convex one with a variable curvature. The curvature is largest close to the throat, and particles close to the wall are pushed against the wall by the free stream.
The further downstream, the more the geometry of the spike is flat approaching
the scenario of the Kramers problem. The implementation of a convex surface with a
variable curvature in the analytical procedure is non-trivial and is beyond the scope of this
work. Nevertheless, the solution to the Kramers problem is used to approach the problem
at a phenomenological level.

1.6 Research Objectives

The key in applications of a radiation source is continuous operation at an affordable cost
of ownership.

In LPP, the first fundamental point – continuous operation – is ensured by the use of a
regenerative target in form of droplets. For each laser pulse a new target is delivered at the
focus of the laser. In this work, the emission of EUV and the expansion of plasma from a
droplet target are studied experimentally, analytically and computationally. And, the
studies are extended to non-ideal alignment between the droplet target and the laser spot
size. The research is novel as previous works on LPP focused on slab targets (see Section
1.5) regarding both EUV emission and matter expansion. In addition, the understanding of
debris field following repetitive LPP source operation in terms of electrons, ions, neutrals,
and fragments is gained. It is an objective of this work to characterize the droplet-based
LPP in terms of EUV emission and matter expansion to tailor the design of future sources,
and exploit the whole potential of droplet-based LPP.

The other fundamental point – affordable cost of ownership – depends on the frequency of
replacement of items damaged by matter (debris) expanding in the vacuum chamber from
the LPP. For this reason in the second part of this work a mitigation system is developed,
based on inertial collisions with the incoming debris. The system is characterized at rarefied
condition experimentally and computationally. And it is applied to LPP to detect the
effects on debris expansion. It is an objective of this work to maximize the lifetime of
components of the source by optimizing the constellation of the components and by
developing and characterizing a debris mitigation system.
1.7 Thesis Outline

The first chapter is dedicated to the definition of laser-produced plasma and its role in the market. The main applications are presented, and EUV lithography is treated in details. The role of droplet-based EUV sources is shown, and is followed by motivation, literature review, and research objectives of this work.

Chapter 2 describes the apparatuses that were developed and used for the experiments.

Chapter 3 gives an overview of the different instruments and measurements techniques, in particular for the detection of: radiation from EUV to infrared; ion total charge; contribution of different ion species; correlation between main propagation direction of EUV emission and ions/electrons expansion in the space; and fragments load.

Chapter 4 exposes the experimental results of the three-dimensional EUV emission distribution. The developed analytical models for density distribution estimation are presented and applied. And, further experimental studies on the effect of irradiance, droplet diameter, and ambient pressure are reported and analyzed.

Chapter 5 presents the measurements on matter expansion in the case of droplets aligned with the laser. The ion measurements are complemented with a model, which is developed and calibrated with experiments. The model contains the main processes happening at the droplet and in the plasma from ablation to rarefied expansion. The influence of various parameters is studied experimentally and explained through the model, as angle from the laser axis, ambient pressure and ambient gas.

Chapter 6 shows the influence of offset between droplet position and laser focal spot. EUV emission is analyzed together with the ion load. Angular distribution, dependence with ambient pressure and droplet diameter of fragment ejection is presented. And, the role of offset in the droplet target ablation is considered together with the mechanism of fragment formation.

Chapter 7 describes the debris mitigation concept and its application with the droplet-based laser-produced plasma. The debris mitigation system is first analyzed in detail with experiments, computations and models to optimize its design, and successively the mitigation capability is presented.
Finally, Chapter 8 reports the main conclusions of this work, and contains some suggestions for future works.
Chapter 2

EXPERIMENTAL APPARATUS

The experiments presented in this thesis were carried out in two different laboratories. The first one is known as ALPS 1, and the second one, which was designed and commissioned as part of this work, is called ALPS 2. The four main hardware components in the laboratories are the vacuum system, the irradiation system, the droplet generation system, and the control system. The description of each component follows in the next Sections.

2.1 Vacuum System

The vacuum system consists primarily of a vacuum vessel and a vacuum pump. In order to achieve the desired vacuum level, the choice of the two components has to match. The operating ambient pressure and the ambient gas were set to 0.1 mbar and argon, respectively. Thus, the vacuum inside the vessel enters the category medium vacuum.\textsuperscript{119} The initial choice of operating pressure and gas comes from the trade-off between the EUV absorption through the gas and the debris mitigation effectiveness.\textsuperscript{49} According to the vacuum level, the flanges of the vacuum vessels were selected of the type ISO-K, which is suitable for the vacuum level and is advantageous for easy replacement compared to ISO-CF.

Both the laboratories were equipped with the same vacuum pump. The pump is made of two stages in series, a screw type main pump, and a roots type booster pump. The pump
was chosen mainly because the volumetric flow rate was sufficient to perform the debris mitigation experiments, and because it is a dry pump, hence oil contamination in the vacuum vessel could be avoided.

More components were added to the vacuum system for the instrument used for ion detection presented in Section 3.2.2. The instrument needed high vacuum, therefore a differential pumping system with a turbomolecular pump was designed and built.

### 2.1.1 ALPS 1

The vacuum vessel in the laboratory ALPS 1 has a spherical geometry with 37 flanges, which are all oriented towards the center of the chamber. The inner diameter of the chamber is 800 mm. All the flanges for instrumentations are DN63 ISO-K. The inside of the chamber is accessed through the door of 625 mm in diameter. The flange for the vacuum pump is at the bottom of the chamber with a standard flange size DN160 ISO-K. The laboratory infrastructure was changed in the frame of this work to host the vacuum pump in the room situated below ALPS 1. The vacuum tubing was optimized to have the lowest possible pressure loss by minimizing the length and by designing the conduct as a straight tube. In order to decouple the vacuum chamber and the vacuum pump for safety and operation reasons, a gate valve was introduced between the two components. Fig. 4 shows the pressure inside the vacuum vessel as a function of the introduced mass flow of argon. The curve is used to estimate the maximum mass flow that can be injected by the debris mitigation system for a desired pressure in the vessel. The trend of the curve is approximately linear and at an ambient pressure of 0.12 mbar, 0.1 g/s of argon can be injected in the vessel.

In the measurements presented in this thesis two flanges stayed occupied by the same components: the laser focusing optics (part of the irradiation system described in Section 2.2) were inserted from the flange opposite to the door and the droplet target dispenser was held by the flange at 45° from the equatorial plane. Inside the vessel, the beam dump (part of the irradiation system) and the tin trap (part of the droplet generation system described in Section 2.3) were kept in the same position during all the experiments.

Experiments performed in ALPS 1 are described in Chapter 4, and Section 7.1.
2.1.2 ALPS 2

The laboratory ALPS 2 has been designed and realized in the frame of this work to extend the capabilities of ALPS 1. The laboratory infrastructure is qualified for the use of Ar, N\textsubscript{2}, Xe, He, H\textsubscript{2}, which imposed constrains on the air conditioning and electrical system. The main working gases Ar, N\textsubscript{2}, and compressed air are directly supplied to the laboratory. Multiple power supplies and cooling lines were specified and placed inside and outside the laboratory. The vacuum system was designed to host the vacuum pump on the outside of the laboratory at a distance of approximately 3 m from the vacuum vessel. In order to minimize the pressure losses between the pump and the vessel, the vacuum pipe was chosen to be a DN320 ISO-K, and four 90° bend were necessary to connect the vessel to the vacuum pump. A gate valve was added between the vacuum vessel and the pump to decouple the two components in case of emergency shutdown and for operation. The optical table was custom made to support laser, beam line and vacuum vessel. Passive damping was chosen for the legs of the optical table with 50% damping at 8 Hz. The distribution of the legs was set according to the weight load on the table. The vacuum vessel is illustrated in Fig. 5. The vessel was designed to accommodate multiple design...
2.1 Vacuum System

ccepts. It has a sub-chamber to host and protect the motorized stage in charge of keeping the droplet targets at the irradiation site, i.e. in charge of moving the droplet dispenser. Four latitudinal windows were implemented and allow a laser-photodiode system to detect the droplet position. This information together with the EUV emission is then used to control the targets position. A sliding flange was added to one face of the vessel for maintenance of the ML collector mirror.

Fig. 5: ALPS 2 facility during operation. The custom-made vacuum vessel is visible on the right. The vacuum vessel is connected with the pump outside the laboratory by the pipe visible below the table.

As for the vacuum system in ALPS 1, the maximum injectable massflow of argon for debris mitigation at a desired pressure was quantified. The trend is approximately the same as in ALPS 1, showing that the increase in length and bends of the vacuum tubing was compensated as aimed by the larger diameter. Experiments performed in ALPS 2 are described in Chapter 5, 6, and Section 7.2.
2.2 Irradiation System

The photons for plasma generation are produced by a laser, reflected by a number of mirrors and focused at the irradiation site, and when no target is present, the photons defocus and travel up to the beam dump, where they are absorbed.

The laser used was derived from a Starlase AO16 from PowerLase. The Nd:YAG diode pumped solid state laser (DPSS) worked at the fundamental wavelength of 1064 nm. The laser was operated at a frequency of 6 kHz, at a range of laser power from 360 W to 1.44 kW (nominal power), depending on the current applied on the gain modules.

The beam-line in ALPS 1 consisted of multiple mirrors and a focusing lens. In order to minimize the focal spot size a large input beam diameter at the lens was set using the natural divergence of the beam with a length of the beam line extended to approximately 4 m. Both the faces of the lens were coated with anti-reflective coating to maximize the transmission of the lens to 1064 nm wavelength to approximately 99%. In ALPS 2 the beam-line was improved by adding a custom-made fused silica window after the lens, towards the plasma. This feature allowed an improved cooling of the optical elements by forced convection, and the easy replacement of the cheaper window instead of the lens in case of transmission decrease. The cooling system was designed and implemented in the irradiation system together with a debris mitigation system for the optical element facing the plasma. A patent on the beam-line of the laser in a EUV source was submitted.\(^{62}\)

The beam dump was designed to absorb the energy of the laser with a safety factor of 1.5. Since in the vacuum vessel during operation the convection is negligible and the radiation losses are too small to prevent the melting of the beam dump, a cooling system was designed. The latter was a serpentine along the cylindrical wall of the beam dump. A cone was placed in the cylindrical cooled mantel made out of stainless steel. The latter was necessary to absorb and reflect the radiation from the laser to the walls. The geometry of the beam dump components was designed to have at least two reflections of the laser photons on the cooled walls. The material of the cone was chosen to be aluminum, in order to maximize the conduction to the cooled wall. And, in order to withstand the higher temperature at the spike of the cone, an insert was added as spike out of a high-melting-
temperature material. Both the laboratories were equipped with the same beam dump design.

**Fig. 6:** Intensity of the focused laser beam at 25 µm from the focus of the lens. The contour of the levels is not perfectly circular due mainly to the pentagonal distribution of the diodes around the Nd:YAG crystal.

The focal spot size at irradiation site was quantified using a CCD. In order to have a range of power compatible with the CCD, four attenuators were added to the beam line. Each attenuator reflected only 1% of the laser beam intensity. The distance between the focusing lens and the laser output was kept the same as in operation. In order to find the focus, the CCD was mounted on a translational stage and data was stored every 12.5 µm. The data presented were taken at the base operating condition of the laser, repetition rate of 6 kHz, and power of 1.44 kW, with the irradiation system used in ALPS 2. **Fig. 6** shows the cross-section of the laser spot 25 µm from the focus. The levels of the contour are not perfectly circular mainly because of the position of the diodes inside the gain modules of the laser, where the diodes are placed at the vertexes of a pentagon around the Nd:YAG crystal. The horizontal and the vertical cut (in x and y direction, respectively) that cross the maximum of the spot are shown in **Fig. 7**. A Gaussian distribution in the space approximates the intensity profile of the laser spot. The curves are used to extract the full width half maximum (FWHM) of the intensity profile, which remains approximately constant along both vertical and horizontal direction.
Fig. 7: Cuts of the spot of the laser at 25 µm from the focus of the lens in horizontal and vertical direction. The FWHM is approximately the same. A Gaussian profile predicts approximately the spot in space.

The radius at half maximum (HM) is half of the FWHM and is represented versus the depth of focus in Fig. 8. The minimum radius at HM is around 32.5 µm, the FWHM 75 µm, which is taken as diameter of the focal spot. The measured radius at HM is not symmetric along the depth from the focus position; which is due to aberration of the focusing element and to the natural laser beam divergence. The same experimental approach was used to determine the focal spot with the irradiation system used in ALPS 1.

Fig. 8: Radius at HM of the laser spot along the depth of focus of the lens. The focal spot size is 75 µm.
Temporally, the laser pulse was characterized with a fast photodiode. The pulse at base operating condition is shown in Fig. 9 together with a Gaussian profile. The latter was constructed to follow the photodiode measurement, and it matches the signal up to around 40 ns, then the laser pulse decay is slower than the Gaussian profile. The laser pulse duration at FWHM is derived from the time profile, at base operating condition was 23.9 ns, and increased up to 39 ns at 360 W. In the modeling of the laser pulse temporal profile in the next Sections, the Gaussian profile satisfies both the FWHM, and the overall laser pulse energy (integral of the laser power) as measured experimentally.

The irradiance resulting from the measurements at the base operating condition of 1.44 kW at 6 kHz was of approximately $2 \times 10^{11}$ W/cm$^2$; and at the reduced laser output power of 360 W the irradiance was $5 \times 10^{10}$ W/cm$^2$.

![Laser pulse time profile](image)

**Fig. 9:** Laser pulse time profile. The blue points represent the photodiode measurements, and the red dashed curve is a Gaussian profile. The laser pulse follows the Gaussian profile up to around 40 ns, and afterwards the laser pulse decay is slower.
2.3 Droplet Generation System

A droplet dispenser, and a trap for the targets that were not irradiated compose the droplet generation system. The irradiated targets used in the present work were tin droplets, which were formed by piezo-electrically modulated Rayleigh break-up. The diameter of the droplets was in a range from 30 µm to 50 µm depending on the size of the orifice of the nozzle, on the operating frequency, and on the applied backpressure. The material of the targets was tin. The design and operation of the droplet dispenser is detailed in Rollinger,\textsuperscript{121} and a patent on the concept has been filed.\textsuperscript{122} The trap for the droplet targets was designed to keep the tin not used for radiation generation in liquid phase, in order to prevent the building of a tin tower. The trap was electrically heated and insulated through ceramic holders from the rest of the vacuum vessel. Additionally, special care was taken in the design of the surface where droplets eventually impact, in order to minimize the bouncing of droplet targets from the surface in the vacuum vessel.

2.4 Control System

A single system was developed to integrate the control of all the actuators of the facility together with the acquisition of the measurements. The first version of the control system was developed in ALPS 1, and successively extended in ALPS 2. The system was capable of controlling the vacuum system, the irradiation system and the droplet generation system simultaneously. The majority of the features of the control system were implemented for continuous and steady operation of the EUV source. Additional features were integrated for the measurements. In particular, the control of the irradiation system was extended with a lasing gate mode. With the latter it was possible to shoot at 6 kHz with nominal power for a fraction of a second, and to stop for a given interval during which the signals retrieved from the measurement cards would be saved to the hard disk. This mode was used in particular for measurements where coating of debris would deposit on part of instruments, leading to a decrease of the signal intensity imposing a systematic error. Specific extensions to the control system were made for each instrument. The most extensive were made for
the three-dimensional emission profile of EUV radiation and for the ion detection using electrostatic filtering. In the former, a robotic arm had to be moved around the chamber to a wanted matrix of points and the motion had to be coordinated with the signal acquisition (see Section 3.1.1). In the latter, the voltages of the electrodes for ion filtering and focusing had to be scanned to measure the ion populations at a given ambient pressure and position (see Section 3.2.2).
Chapter 3

EXPERIMENTAL DIAGNOSTICS AND TECHNIQUES

Plasma sources used for radiation generation have an inherent challenge: the production of expanding matter together with the wanted photons. For this reason, in a complete study of an EUV source both radiation and debris have to be considered. The diagnostics and the techniques for the measurements of both the aspects are exposed in the following Sections. Furthermore, the technique for the characterization of the flow field from the mitigation system is exposed.

3.1 Radiation

Photons emitted by the plasma in a radiation source can be distinguished between those that are desired, and the rest, which are a byproduct. For an EUV sources, photons centered at a wavelength of 13.5 nm in a bandwidth of EUV +/-2% are desired, and the rest of the emission is known as out-of-band.
3.1 Radiation

3.1.1 EUV Radiation

**EUV Energy Monitor**

EUV photons are measured by an EUV energy monitor. The latter consists of an aperture, a Mo/Si ML mirror of $5^\circ$ incidence, a Zr filter and a x-ray photodiode with a receiving aperture of 3 mm in diameter, as shown in Fig. 10. The overall length of the EUV energy monitor is 170 mm, small enough to be mounted and moved inside the vacuum vessels to measure around the plasma.

![Fig. 10: Schematic of the EUV energy monitor. The instrument is composed by a first aperture, a ML mirror of $5^\circ$ AOI, a second aperture, a Zr filter and a photodiode.](image)

The Zr filter together with the ML mirror gives an overall detectable bandwidth (BW) of EUV $\pm 2\%$ centered at 13.5 nm, as shown in Fig. 11. The bandwidth is calculated as the value at FWHM of the transmission up to the photodiode. The difference between the curves with and without the Zr filter is a quasi-homogeneous reduction of the magnitude, which is given by the small variation of the transmission coefficient of the Zr filter between 12.5 nm and 14.5 nm (for a 200 nm filter it varies between 52% and 44%, respectively).
Alignment Procedure

For maintenance of the instrument, the ML mirror hosted at the back has to be inspected. When the signal from the instrument decreases, and a coating on the ML mirror is visible, the ML mirror should be rotated or replaced. One ML mirror can be used on four different spots, as the dimension chosen for the ML mirror is larger than the spot. A fresh spot is found by rotating the mirror of 90°. Every time that the mirror is inspected, rotated or replaced, the EUV energy monitor should be realigned.

The alignment procedure follows three steps. First, the instrument and a laser pointer are mounted and leveled on an optical table. Second, the laser beam is aligned with the EUV energy monitor. In order to have two apertures for the alignment, the ML mirror at the back of the instrument is not installed, in this way one aperture at the front and one at the back are visible. To facilitate the alignment of the laser beam, the pointer is mounted on a translational stage with tilt-adjust. The third and final step is the alignment of the ML mirror. Once the mirror is remounted at the back of the instrument, it is aligned to the photodiode position using the two screws at the back of the mirror holder. In order to see the beam during the alignment, the photodiode assembly is demounted and the beam
aligned at 35 mm from the housing at the marked position. Then the photodiode assembly is remounted, and the instrument is ready for use.

**Robotic Arm**

The EUV energy monitor was moved inside the vacuum chamber by a robotic arm. The latter moved the instrument on a sphere, leaving the distance between the irradiation point and the instrument aperture constant and equal to 185 mm.

To move the instrument on a sphere, two degrees of freedom had to be implemented. Two concepts were mainly considered, one with a rotational and a translational motion, and one with two rotational degrees of freedom. The latter was realized, as it was cheaper, and more robust against the debris generated by the plasma. In order to minimize the costs, and to leave the inside of the chamber as clean and free as possible, one rotational system was implemented outside the chamber and transmitted through a rotary feedthrough to a semi-circular holder in the inside. The system on the outside was in charge of longitudinal translation. Together with a motor and a rotary feedthrough, also an absolute encoder was installed, in order to monitor the position of the instrument during measurements. A picture of the longitudinal motion stage is presented in Fig. 12.

![Fig. 12: Outside the vacuum chamber. The motor and the rotary feedthrough for the robotic arm movement in longitudinal direction. An encoder mounted between the motor and the feedthrough monitors the longitudinal position of the arm.](image)

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The stage for latitudinal motion was located inside the vessel, and mounted on the semi-circular holder that was moved longitudinally from the outside. The materials of the components were chosen, in order to ensure vacuum compatibility and to avoid carbon contamination in the system. As in the outside, an absolute encoder continuously measured the position of the arm. The encoder was placed inside the cubic frame, which sustains the shaft of the arm, in order to protect it from debris. Fig. 13(a) shows the stage mounted in the vacuum vessel.

In the experiments shown in this thesis, the robotic arm was used to move the EUV energy monitor to measure the three-dimensional EUV emission from the droplet targets. The procedure of moving the robot to the predetermined positions with the feedback given by the encoders and the data acquisition were fully automated.

Fig. 13: (a) Picture of the robotic arm inside the vacuum chamber. The EUV energy monitor is mounted on the lower right. The arrows show the rotational movements in longitudinal and latitudinal directions. (b) Schematic of the experimental setup looking at the focusing lens. The laser irradiates the dispensed tin droplets, and the EUV energy monitor scans a predetermined range in one hemisphere. (c) Frame of reference for the measurement positions and for the Mercator projection. \( \lambda \) and \( \varphi \) are the longitude and the latitude, respectively. (d) Frame of
3.1 Radiation

reference for the two-dimensional representations. $\beta$ and $\theta$ are the angle versus the horizontal plane and the angle versus the laser axis, respectively.

In the experiments shown in Section 4.1 100 individual measurements were taken in 88 different positions. The EUV radiation was measured in a range of 1.8 sr in one hemisphere (shown in Fig. 13(b)). The robotic arm was moved along trajectories of constant longitude and constant latitude, defined in Fig. 13(c) by $\lambda$ and $\varphi$, respectively. The two-dimensional representations that are used in Chapter 4 are plotted in the frame of reference defined in Fig. 13(d), which is more convenient for data with an axis of symmetry around the laser axis. Together with the EUV signal, signals from UV photodiodes placed outside the vacuum chamber were simultaneously recorded. One photodiode monitored the laser pulse along the beamline, and three other photodiodes were placed at 45° with respect to the laser axis in three different directions. The EUV signal was captured by a 1 GS/s digitizer card, which was triggered by the laser photodiode. The main data processing steps included the computation of the average of the time integral, as well as its standard deviation. During post-processing, the absorption of EUV by the Ar background gas was taken into account and removed from the data.

3.1.2 Out of Band Spectrometer

To measure the plasma emission spectrum in the 280 nm to 1100 nm range a spectrometer with a resolution of 0.035 nm was used. The spectrometer was mounted at 90° to the laser axis. A deuterium lamp was used to calibrate the instrument in a previous work in the range of 200 nm to 400 nm. In this work also tin emission lines at larger wavelengths were considered for different operating parameters of the irradiation system and of the droplet generation system. The results are exposed with the EUV emission in Section 4.4 and with the ion load in Section 5.4.1. Experiments were performed at an irradiance of $2 \times 10^{11}$ W/cm$^2$ with different droplet diameters (30 µm and 50 µm), and with different ambient pressures (from $7 \times 10^{-4}$ mbar to 1.4 mbar).
3.2 Electrically Charged Particles

An imbalance of the flux of charged particles is detected by measuring the current flowing through a conductor exposed to the plasma. Depending on where the charged particles flux has to be measured different techniques apply. To measure close to the irradiation location the conductor has to be as small as possible in order to minimize the perturbation on the plasma, but having a surface large enough to collect a detectable signal. Normally, these probes are in form of wires with diameters of tens of micrometers (the material has to have a high melting temperature as tungsten) and are used to measure the charged particle flux at distances from the plasma from some millimeters to tens of millimeters. These detectors are called Langmuir probes, and by biasing the wire with different voltages, and by assuming a continuum in the plasma, the electron temperature and the charge density can be estimated. However, at distances of tens to hundreds of millimeters, where the collector mirrors of an EUV source are located, the charge flux on a wire becomes too small to be detected, and another measurement technique has to be used. The latter is based on Faraday cups, where instead of a wire a plate is used to increase the detected charge flux. Also the surface of Faraday cups is biased to measure either ions or electrons. An advantage of measuring at larger distances from the irradiation location is that the instrument can be larger and more features can be implemented. In particular, the detectors used in this work were shielded electrically with a grounded grid or casing. Having the advantage to minimize the perturbation of the expanding plasma through space charge effects, and to impose a clean and controlled electric field to the incoming particles. As in this work the load on the collector mirror is addressed, the measurement techniques are chosen for the detection at distances of hundreds of millimeters, as Faraday cups. Furthermore, an instrument for charge state and kinetic energy discrimination was developed and built. The latter as the Faraday cups was placed at distances from the irradiation site in the order of hundreds of millimeters.

3.2.1 Faraday Cups

The overall ionic population arriving at distances representative for collector mirror location were measured with Faraday cups at different operating conditions. A copper plate
was inserted in an insulator, which was surrounded by a metallic grid. The latter ensured that the electric field produced by the bias voltage on the plate was localized between the plate and the grid. The sensing area of the plate was 2 mm in diameter, and the maximum distance from the plasma was 150 mm, which corresponds to a solid angle of $1.4 \times 10^{-4}$ sr. The connection from the FC to the measurement card was a BNC cable with the shield connected to the grid surrounding the instrument, in order to minimize electrical disturbances. A circuit was built to bias the sensitive surface and measure the current due to the impact of the charged particles. This detector was mainly used for total ion current measurements, with a bias on the surface of -30 V, enough to repel electrons with a kinetic energy smaller than 30 eV, which is approximately one order of magnitude more than expected. The total ion current of the ions $i$ arriving up to the FC is proportional to the current measured by the digitizer card, and is given by the number of charged particles $N_i$ and to the mean ion charge $\bar{Z}$ of the plasma impacting the FC (Eq.(2)). Therefore, the time integral of the measured current is proportional to the total ion charge $Q$, which is proportional to the number of charged particles arriving up to the FC (Eq.(2)).

$$i = \frac{d(N_i e \bar{Z})}{dt},$$

$$Q = \int_{t_1}^{t_2} i \, dt = N_i e \bar{Z}.$$  \hspace{1cm} (2)

In Eq.(2) $t_1$ and $t_2$ define the interval of ion arrival on the sensitive surface.

### Multiple Faraday Cups

The FC design was replicated, and multiple FCs were used together to measure the symmetry of the plasma expansion together with EUV radiation. The aim of the experiment was to quantify the EUV emission profile versus the plasma expansion in off-design conditions. 8 FCs were placed around the plasma at 100 mm distance and covering a solid angle of around $1\pi$ sr. The signals from the FCs were stored synchronously with those from two EUV sensors.
Fig. 14: *Picture of the experimental setup with the 8 FCs and EUV sensors during operation.*

The experimental setup is shown during an experiment in Fig. 14, the eight FCs and one EUV sensor are visible in the picture; a second EUV sensor was placed on the other side of the chamber. For clarity Fig. 15 shows a sketch of the experimental setup. The FCs were placed on two planes, defined by rotating the equatorial plane of +45° and -45° with the laser axis as axis of rotation. FC1, FC3, FC5, FC7 were placed on the plane rotated of +45°, and FC2, FC4, FC6, FC8 were placed on the other plane rotated of -45°. On each plane, the FCs were placed at 45° and at 60° with respect to the laser axis taking the irradiation point as reference. More precisely, FC1, FC2, FC3, FC4 were at 45° from the laser axis, and FC5, FC6, FC7, FC8 were at 60° from the laser axis. Instead, the two EUV sensors were placed on the equatorial plane at a distance of 150 mm from the plasma, and oriented towards the plasma at an angle of +/- 75°.
3.2 Electrically Charged Particles

Fig. 15: Schematic of the position of the FCs and EUV detectors. The EUV detectors are placed on the equatorial plane. The FCs surround the lens, and are numbered to identify them. The locations of the FCs are on two planes rotated with respect to the equatorial plane with the laser axis as rotation axis of +/-45°.

3.2.2 Electrostatic Ion Energy Analyzer

A spherical sector electrostatic energy analyzer (ESA) was designed and built for the operation in the ALPS 2 vacuum vessel. The instrument is capable of filtering the ions depending on their charge to mass ratio and kinetic energy by imposing a constant electric field in a direction perpendicular to their trajectory. Fig. 16 shows a schematic of the horizontal cut of the instrument. The aperture below is directed to the plasma, ions entering in the casing travel up to a second aperture just before the electrodes, which is drilled through a stainless steel plate. The latter is grounded during operation, in order to leave the space between the two apertures with a negligible electric field, hence without accelerating the ions. After the second aperture, the ions enter in the spherical sector, shown in red. The latter is between two electrodes with constant curvature, leading to a constant electric field. The material of the electrodes was oxygen-free copper. When ions with the right charge to mass ratio and kinetic energy pass in the spherical sector, they arrive up to a third aperture that is just before the detector, otherwise they are deflected not enough or too much and hit one of the electrodes. During operation the electric field in the spherical sector was continuously changed, in order to scan across all the possible mass to charge ratios and kinetic energies.
Fig. 16: Schematic of the horizontal cut of the ESA. The ions enter in the instrument from below, travel to the spherical sector between the two electrodes (shown in red) and if their charge and velocity is compatible, are deflected up to the detector represented by the green box. The electrodes are two orange elements. The instrument is mounted in a casing for differential pumping. For alignment purposes a hole was manufactured in the casing and in the outer electrode.

A further feature implemented in the instrument was the channel, which crosses the whole instrument. The latter was essential for precise alignment of the instrument in the vacuum vessel at the desired angle from the laser axis. The hole on the back of the casing was then closed during operation with a screw. A picture of the instrument is presented in Fig. 17. A system of electrostatic lenses was placed together with an ionizer between the first aperture facing the plasma and the second one before the spherical sector. The ionizer was implemented in the instrument for the re-ionization of tin neutrals, however the electron impact cross-section on the neutrals was not large enough to see a significant difference in the detection. All the cables for voltage control of the electrodes and for signal acquisition were fitted through the flange for differential pumping.
3.2 Electrically Charged Particles

Fig. 17: Picture of the ESA. The spherical sector is shielded using grounded plates. Between the first aperture facing the plasma in lower part of the picture and the aperture of the spherical sector a system of lenses and an ionizer was inserted. All the cables are fitted through the flange used for differential pumping of the instrument.

Ion Filtering and Resolution

The trajectory of the ions in an ideal spherical sector is calculated analytically. An ion with kinetic energy $E_{\text{kin}}$ and charge $q$, is bended constantly up to the outlet of the spherical sector when the voltage difference $\Delta V$ applied to the electrodes is:

$$\Delta V = \frac{E_{\text{kin}}}{q} \left( \frac{r_2}{r_1} - \frac{r_1}{r_2} \right) = \frac{E_{\text{kin}}}{q} \varepsilon_{\text{ideal}} = 0.5 \varepsilon_{\text{ideal}} \frac{m}{q} v_{\text{ion}}^2,$$

where $r_1$ and $r_2$ are the radii of the inner and outer electrode, respectively; and $\varepsilon_{\text{ideal}}$ is the ideal geometrical factor dependent on the choice of the radii. The design of the spherical sector was such that $r_1 = 48.8$ mm and $r_2 = 60.2$ mm, leading to $\varepsilon_{\text{ideal}} = 0.42$. The ion kinetic energy is proportional to the ion mass $m$ and to the ion velocity $v_{\text{ion}}$ squared. The latter transformation is useful for the direct use with the measurements, as the ion velocity is calculated from the time-of-flight, with the laser pulse as reference.

In order to verify that the analytical relation that assumes a constant electric field across the whole spherical sector without influences from the surrounding, simulations were
performed. The whole instrument, from the first aperture facing the plasma to the detector was modeled using the software SIMION. The latter calculates the electric potentials in a system of electrodes solving the Laplace equation in the three-dimensional space by finite difference method. The geometry of the electrodes can be imported from common CAD file formats. Once the electric potential in space is evaluated, trajectories of the charged particles are predicted. Examples of the model results are shown in Fig. 18.

Fig. 18: Results of simulations using SIMION, the components of the instrument are visible in the horizontal cut. (a) The voltage difference in the spherical sector was set to allow access to Sn$^+$ ions with a kinetic energy of 500 eV. (b) The voltage difference in the spherical sector was set to allow access to Sn$^{2+}$ ions with a kinetic energy of 1000 eV.
3.2 Electrically Charged Particles

The model was then used to adjust the analytical analysis of deflection given the electric potential of the overall instrument, and not only of the two electrodes of the spherical sector (assumption of the analytical analysis). The plot in Fig. 19 is the comparison between the analytical model and the prediction of the software SIMION, where all the components in the instrument are considered and not only the spherical sector. The analytical relation overpredicts the kinetic energy of the particles that pass through the electrostatic filter by about 15%. The non-ideal geometrical factor of the instrument is $\varepsilon_{si} = 0.49$.

![Fig. 19: Relation between the voltage difference applied in the spherical sector and the kinetic energy over the charge state. The red continuous line shows the ideal result only dependent on the geometry of the spherical sector. The blue dashed curve considers the whole instrument in the SIMION simulation, reproducing a result closer to experimental condition with the complete instrument.](image)

The resolution of the instrument is defined by:  

\[
\frac{\Delta E_{\text{kin}}}{E_{\text{kin}}} = \frac{\omega}{R(1 - \cos \varphi) + \delta \sin \varphi},
\]

where $\omega$ is the aperture size, $R$ is the mean sphere radius, $\varphi$ is the angle described by the electrodes (127°), and $\delta$ is the distance from the third aperture at the exit of the spherical sector to the detector. The angle described by the electrodes was chosen to be 127° due to
the refocusing effect at the third aperture.\textsuperscript{125-127} The resolution of the instrument with respect to the kinetic energy is calculated with Eq.(4) and the physical dimensions of the designed instrument, and equals around 3%.

Inside the ESA, the spherical sector was enclosed in a Faraday cage in order to avoid any possible acceleration of the charged particles between the aperture and the spherical sector.

**Detector**

At the exit of the spherical sector a channel electron multiplier (CEM) was used to detect the selected particles. A CEM rather than a FC was selected, because the former amplifies physically the signal by several orders of magnitudes. The amplification mechanism is based on an electron cascade, which is generated by production of electrons at every impact onto the walls of the inner channels of the instrument, and by the acceleration of the electrons by the voltage difference imposed across the CEM. The surface of the CEM internal channels is typically chosen out of silicon oxide to maximize the production of secondary electrons at every impact. The CEM used in this work was a MAGNUM type from Photonis. The particular CEM incorporated six individual spiral channels to increase the gain of the signal, which was of around $10^7$. In the experiments, the active surface of the CEM was biased with 2.5 kV, and around the CEM a Faraday cage was built to contain the electric field and to avoid any possible discharge with the surrounding metal surfaces. Furthermore, the grounded grid placed in front of the CEM active surface defines the distance along which the charged particles are accelerated in the electric field caused by the bias. Since the distance was of 5 mm versus the overall distance from the plasma of 456 mm, the change in the time of flight, which would be dependent on the charge to mass ratio of the particle is neglected. The kinetic energy increase of the ions entering in the space between the grid and the bias follows directly from the bias voltage.
3.2 Electrically Charged Particles

![Image](image_url)

**Fig. 20:** Assembly of the detector. The channel electron multiplier (CEM) is hosted in insulated holder. The red cable attached towards the front of the instrument supplies the voltage for the bias. In operation the top is enclosed by the second mirrored piece of the holder. At the back the signal is received directly by a BNC connector, in order to minimize noise to the signal.

The detection efficiency of CEMs was evaluated by Krems et al.\(^{128}\), who showed that the detection efficiency of a CEM is mainly a function of the kinetic energy and of the mass of the incoming ions. The same method was successfully applied by Burdt et al.\(^{125}\) for tin ions by using the empirical relation from Krems et al.\(^{128}\) and is applied also in this work. The empirical relation for the detection efficiency \(\gamma\) is:

\[
\log(\gamma) = A + B \log\left( E_{\text{kin, CEM}} \right) + C \log\left( E_{\text{kin, CEM}} \right)^2 + D \log\left( E_{\text{kin, CEM}} \right)^3 + F \log\left( E_{\text{kin, CEM}} \right)^4, \tag{5}
\]

where the coefficients for tin are:\(^{125}\) \(A=-1.233, B=2.217, C=-2.094, D=1.427, F=-0.456\), and \(E_{\text{kin, CEM}}\) is not the initial kinetic energy of the ions but, the one of the ions impacting on the CEM surface, e.g. the kinetic energy of a singly charged ion would increase of 2.5 keV with respect to the initial value.

**Differential Pumping**

As the aim of the experimental campaign on the ion population measurements was to operate at regimes representative for EUV sources, the whole instrument was designed for operation at pressures up to tenths of millibars. Consequently, a differential pumping system was integrated in the instrument, in order to ensure that no collision happened during the energy filtering inside the ESA and to ensure normal operation of the CEM.
The ESA was connected through a stainless steel flexible pipe of 50 mm inner diameter to a tur
cromolecular pump, and the cables for control of the instrument and for signals passed through the same pipe up to a feedthrough by the turbomolecular pump. The pressure inside the ESA was measured to be $1.5 \times 10^{-5}$ mbar at the ambient pressure of $6 \times 10^{-2}$ mbar, which was the largest pressure, where ions could be detected. The mean free path calculated for the inside of the ESA with $1.5 \times 10^{-5}$ mbar was more than 50 times larger than the distance that the ions had to travel inside the instrument. Therefore no inertial collision is expected to happen between the tin ions and the background gas inside the instrument. Fig. 21 shows the differential pumping system on the outside of the vacuum vessel. A turbomolecular pump is placed at the bottom, and on the side a feedthrough is used to extract the cables for the voltage control and signal acquisition. The turbomolecular pump was capable of a maximum volumetric flow rate of 550 l/s.

Fig. 21: System for the differential pumping, with the electrical feedthrough for voltage control and signal acquisition of the ESA.
3.3 Sample Exposure

**Operation**

The ESA was located in a range between 45° and 120° from the laser axis, and at a distance from the irradiation region of 150 mm. At every angular position the instrument was aligned using different mirrors placed at the focus of the lens, and a laser pointer parallel to the laser axis.

The voltage difference in the spherical sector was changed from 0 to 500 V, and the voltage control and the signal acquisition was automated so that for every voltage setting, pressure and angle 40 plasma events were stored. The measurements were then filtered using the range of the EUV energy per pulse as criterion.

3.3 Sample Exposure

Another form of debris comes from fragments of the droplet target, which are emanated by the shattering of the droplet. This detrimental effect was quantified with sample exposure. The samples were Si polished plates with a dimension of 10 mm x 10 mm. Silicon was chosen as the surface is conductive, hence SEM imaging is possible without surface charging effects that would deflect the electrons used for imaging. Eight samples were mounted on a circular holder for each run as shown in Fig. 23. A stepper motor was used to control the exposure of the multiple samples during the runs. As also fragment ejection is dependent on whether the droplet targets are irradiated head on or not, a shutter was implemented in the measurement system. The shutter together with the motorized sample holder was enclosed in a box as shown in Fig. 22. A control system for the shutter was implemented in the main control system to open the shutter only when the EUV level was larger than a given level. The EUV level was measured by the EUV energy monitor, and the signal was integrated through a circuit to avoid any software real time post-processing that would have been too slow to detect every single plasma at 6 kHz. The EUV energy monitor during these measurements was located at 60° with respect to the laser axis. The trace of the EUV integrated signal was then recorded for every sample exposure. The reaction time of the control system was limited by the refresh rate and by the time needed
for the shutter to open and close, together was around 100 ms. Therefore the single plasma filtering could not be ensured during the sample exposure by the hardware.

On the outside of the box was mounted the debris mitigation system described in Section 7.1.1. The whole box was aligned to the plasma following the same procedure as used for the ESA as presented in Section 3.2.2.

![Image 1](image1.png)

**Fig. 22:** Setup for fragment measurements. The motorized sample holder is placed behind the shutter used to expose the sample only with source representative EUV level. The aperture size was of 2 mm in diameter.

![Image 2](image2.png)

**Fig. 23:** Si samples mounted on the holder before exposure.

90 samples were exposed at different ambient pressures, different operating conditions of the debris mitigation system and at different location in the vacuum chamber. The samples were exposed to a number of plasmas ranging from 1 million to 10 millions. After exposure, the samples were inspected under an optical microscope and with a SEM. As the dimension of the fragments was mainly in the micrometers range, the images from the
optical microscope were used and the image post-processing was automated. For every sample 5 images of fragments and 1 background image were captured. The background image was used in the post-processing to remove the background and set a color function for the fragment detection. Fig. 24 shows an example of the automated identification of the fragments from the microscope image to the data matrix. The latter was then further elaborated to detect dimension and amount of the fragments.

Fig. 24: Example of the post-processing from the image captured by the microscope (a) to the data matrix used for the quantification of the fragments (b).

3.4 Rarefied Flow Field

The performance of the debris mitigation system developed in this work had to be at first quantified. The debris mitigation system was based on the deflection of debris through inertial collisions with gas atoms or molecules from a high momentum layer of gas, called gas curtain. In order to accelerate the flow in the region of interest to produce the high momentum layer, a nozzle system was designed and implemented. The type of nozzle, which was implemented, was an aerospike and is described in Section 7.1.1. In order to quantify the performance of the system, the flow field at the exit had to be measured. As the pressure of the chamber is of tenths of mbars, and the characteristic dimension of the flow field is in the order of magnitude of the aerospike throat height (TH), the flow field is
rarefied with a $Kn$ number between 0.42 and 6.4, and the experimental procedure was planned accordingly.

Experiments were performed in continuous operation in the spherical vessel installed in ALPS 1. The ambient pressure in the vacuum vessel was varied from $7 \times 10^{-2}$ mbar to 1.5 mbar by changing mass flow of the injected gas, and by throttling the vacuum pump. The gas used in the experiments presented was argon, and the maximum mass flow at an ambient pressure of 0.12 mbar was around 0.1 g/s.

The gas flow was investigated by mapping the impact total pressure in the flow field downstream of the throat of the aerospike. The mapping points lay in two planes, one perpendicular to the other. By defining the coordinate system with the $x$-axis plotting streamwise direction, the $z$-axis plotting the height from the ideal exit of the aerospike, and the $y$-axis plotting the width of the aerospike (as it is depicted in Fig. 25(a) and Fig. 25(b)), one measurement plane corresponds to the $x$-$z$ plane, and the other one to the $y$-$z$ plane. In the $x$-$z$ plane, which is represented in red in Fig. 25(b), the resolution of the measurement points is finer, with 144 points over the measurement field of 56 TH in the $x$ direction, and 25 TH in the $z$ direction. The measurement points in the $y$-$z$ plane are taken in order to evaluate the strength of the three dimensional flow structures, and 35 measurement points are distributed homogeneously in 15 TH in $y$ direction and 20 TH in $z$ direction. The mapping was performed by moving a sting probe connected to a relative pressure sensor (see Fig. 25(a)). The distance from the aperture of the sting probe to the pressure sensor was kept as short as possible in order to minimize the time needed by the rarefied flow to fill the volume up to the sensor membrane. The sting probe had an inner tube diameter of 0.65 mm. Two different capacitive pressure sensors were used to measure the impact pressure. The difference between the two sensors was the range and the accuracy; in the first it was 30 mbar and $\pm 0.06$ mbar, respectively, in the second of 500 mbar and $\pm 0.75$ mbar. The latter type of pressure sensor was also used to measure the total pressure before the throat. The ambient pressure in the vacuum chamber was monitored by a pirani gauge.
3.4 Rarefied Flow Field

The Mach number was derived from the impact pressure using the Rayleigh relation, using the same method that was already applied at similar flow conditions. The assumption in the Rayleigh relation is that the measured pressure equals the total pressure after the normal shock, which is formed due to the sting probe itself. And, the shock is modeled as ideal and isentropic, leading to the following relation:

\[
\frac{p_{0,i}}{p_0} = \left( \frac{\gamma + 1}{\gamma - 1} \right) \left( \frac{2\gamma}{2\gamma - (\gamma - 1)} M^2 \right)^{\frac{\gamma}{(\gamma - 1)}},
\]

where \( p_{0,i} \) is the impact pressure, \( p_0 \) is the stagnation pressure supplied to the aerospike, and \( M \) is the flow Mach number in the measured position.

At low pressures, rarefaction and non-equilibrium effects might be present in the measurement of the impact pressure. These effects are due to an inadequate number of collisions to fully compress the shock that forms about the probe. Where this is the case, the pressure loss across the normal shock is less than that predicted by the Rayleigh relation. In order to take into account possible measurement errors due to rarefaction and non-equilibrium effects on the probe, a correction was considered. The correction depends on the Reynolds number of the sting probe \( Re_{sting} \) that is defined as

\[
Re_{sting} = \frac{\rho D_{sting} U}{\mu}.
\]

If \( Re_{sting} \) is lower than 5.6, the following relation should be used:

\[
\log\left( \frac{p_{0,corr}}{p_{0,i}} \right) = 0.089 - 0.129 \log\left( Re_{sting} \right),
\]

Fig. 25: (a) Sketch of the experimental setup. (b) 3D view of the measurement planes. The \( x-z \) measurement plane is represented in red, and the \( y-z \) plane in green.
where \( p_{0,\text{corr}} \) is the corrected value of the impact pressure. \( \rho, U, \) and \( \mu \) in the definition of \( \text{Re}_{\text{sting}} \) are evaluated from the computations (same method applied by Boyd et al.\(^{129}\)) and depend on the position of the probe in the flow field, and \( D_{\text{sting}} \) is the inner diameter of the sting probe. For the base case at 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure, \( \text{Re}_{\text{sting}} \) was smaller than the threshold only at locations immediately close to the surface after the exit of the aerospike, and the value at the core of the flow was around \( \text{Re}_{\text{sting}} = 70 \). A further possible cause of disturbance to the measurements was the interaction between the shock generated in front of the probe and the surface of the spike. In order to avoid unwanted measurement errors, the data points immediately close to the surface of the spike, at distances smaller than the outer diameter of the sting probe, were not considered, and therefore the correction was not needed either.
3.4 Rarefied Flow Field
Chapter 4

EUV EMISSION

The knowledge of how EUV is emitted is fundamental for the design of an EUV source. Depending on the EUV energy per pulse and on the extraction angle needed from the source, the EUV distribution around the droplet targets determines dimension, shape, and position of the collector mirror in the source.

In this Chapter first experimental results of the three-dimensional EUV emission at base operating condition are shown, then analytical models are presented to physically explain the measurements with the dynamics of plasma expansion. The model is validated with results from slab targets before being applied to the droplet target experiments performed in this work. The difference that the target geometry imposes on EUV emission is highlighted, and its dependence on plasma dynamics is shown. The analytical models applied to the droplet-based LPP predict a plasma expansion elongated towards the laser axis. Finally, the influences of changes in laser irradiance, droplet size and ambient pressure on the EUV emission distribution are presented. The analysis of the three parameters is based on experimental results and on physical models.
4.1 Experimental Results

Experiments were performed in the ALPS 1 laboratory, using the EUV energy monitor mounted on the robotic arm (described in Section 3.1.1), which moved the sensor around the droplet-based LPP.

The mean EUV energy emitted in a EUV +/-2% BW centered at 13.5 nm is qualitatively plotted in Fig. 26. The black points correspond to the measurement points, and the map is obtained by cubic interpolation. The two axes represent the longitudinal and the latitudinal angles and the origin corresponds to the laser axis (shown in Fig. 13(c)). The shape of the map is not continuous over the whole range because of geometrical constraints in the movement of the robotic arm. The map shows that the EUV emission decays from the center outwards in both directions.

![Fig. 26: Mercator projection of the EUV mean energy at EUV +/-2% BW centered at 13.5 nm versus the detector position. The origin corresponds to the laser beam intersecting the sphere of the mapping. The black points represent the positions of the measurements. The map is obtained by cubic interpolation and shows that EUV energy decreases moving away from the laser beam.](image)

In order to quantitatively show the emission distribution, multiple cuts are shown in Fig. 27 by different markers and colors. The cuts are taken following divergent lines, which have the plasma as their origin. Every cut is defined by the angle $\beta$ between itself and the horizontal plane (shown in Fig. 13(b) and Fig. 13(d)). The angle on the abscissa (defined
with $\theta$ (in the next sections) is equal to the angle between the laser axis and the line defined by the plasma and the position of the aperture of the EUV energy monitor (shown in Fig. 13(d)). The most important feature revealed by these experiments is that, by using 30 $\mu$m droplets as a target and 80 $\mu$m as a laser spot size, the EUV radiation is emitted up to angles from the laser axis larger than 120°, with approximately 50% of the maximum emission at 120°. The difference between the cuts is smaller than the standard deviation, therefore the emission is assumed to be axisymmetric. The discussion of these results follows in Section 4.3 with the help of the model developed in the next Section.

**Fig. 27:** EUV mean energy in the EUV +/-2% BW centered at 13.5 nm versus the angle from the laser axis for different cuts. The cuts are lines defined by the intersection between the mapped spherical surface and a plane defined by the laser axis and the angle $\beta$ from the horizontal plane, which is reported in the plot legend ($\beta$ is sketched in Fig. 13(b) and Fig. 13(d)). The blue circles refer to the cut at $\beta = 10^\circ$, the black squares to the cut at $\beta = 0^\circ$, the green triangles to the cut at $\beta = -20^\circ$, and the red crosses to the cut at $\beta = -40^\circ$. The error bars show the standard deviation of the measurements. The difference between the cuts is smaller than the standard deviation, therefore the emission is assumed to be axisymmetric.
4.2 Physical Processes and Analytical Models

The aim of the presented model is to develop an analytical function, which can be calibrated with the measurements and reveal details on the dynamics of plasma expansion.

4.2.1 Description

The model contains the absorption of EUV radiation by the plasma surrounding the EUV emission-dominant region (EDR) during the phase of EUV emission. This means that the opacity has to be estimated as a function of the density and temperature of the plasma, which are a function of the spatial position. The EUV absorption is calculated from the Beer-Lambert law as follows:

\[ \frac{I_{\text{EUV}}}{I_{0,\text{EUV}}} = \exp\left( -\frac{\tau_{\text{EUV}}}{\tau} \int_{EDR}^{\infty} \kappa \rho \, dl \right), \quad (8) \]

where \( I_{\text{EUV}} \) is the intensity of the radiation going through the plasma up to the EUV energy monitor, \( I_{0,\text{EUV}} \) is the intensity of the EUV emitted isotropically by the EDR, \( \kappa \) is the Planck mean opacity, \( \rho \) is the mass density, \( l \) is the coordinate along which the EUV ray is propagating, and \( \tau_{\text{EUV}}/\tau \) is the optical thickness in the EUV window versus the mean optical thickness. According to Murakami et al., the optical thickness ratio can be estimated to be \( \tau_{\text{EUV}}/\tau = 12 \). For the Planck opacity the following relationship is used:

\[ \kappa_v = a_m T^s \rho^{r_m}, \quad (9) \]

where \( a_m, s, r_m \) are the material constants (for tin respectively \( a_m=328.55, \ s=-1.588, \ r_m=0.228 \)) and \( T \) is the temperature. The Planck mean opacity is used, because the region of interest is located where the EUV is generated and manages to escape through the optically thin plasma. The opacity model is applied at the lower limit of its validity range (electron temperature from 30 eV to 1 keV and density from 0.1 g/cm\(^3\) to 10 g/cm\(^3\)). It is commonly used for EUV conditions with accurate results.

The density and the temperature are estimated by either assuming isentropic or isothermal expansion. The isothermal expansion is evaluated as the density and temperature in the inner region are high enough to satisfy the following condition: the characteristic distance
of heat diffusion into the plasma in an increment of time is larger than the distance covered by the expansion of the plasma, and the heat flux is large enough to compensate for the cooling due to radiation and expansion. This leads to a plateau in the temperature profile along the main expansion direction of the plume.\textsuperscript{136,137} Instead, when the characteristic distance of heat diffusion through electron conduction becomes smaller than the expansion distance, the closest thermodynamic approximation equals the adiabatic description of the plasma. Additionally, it is assumed that there is no entropy generation, hence that the plasma expansion is isentropic. This thermodynamic approximation is better suited to the outer region of the plume.

The model has to be 2D-axisymmetric in order to explain the spatial distribution of the EUV radiation. The laser beam propagates along the symmetry axis. It is assumed that the iso-density lines follow ellipsoids. The ellipsoidal plasma shape is chosen because it gives an additional degree of freedom compared to the isotropic 1D spherical case. Fig. 28 shows a sketch of the model set-up, where the axes are centered at the EDR. The parameters for the ellipsoidal shape and position, namely the semi-major axis \( a \) and the semi-minor axis \( b \), and the offset position \( z_0 \) are represented together with the angle \( \theta \), which is equal to the angle between the EUV ray and the laser axis. Moving the coordinate system from the middle of the ellipsoid to the EDR is advantageous for the formulation and solution of the integral in Eq.(8), where the lines of sight from the EDR to the EUV sensor are considered.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig28.png}
\caption{Schematic of the model setup. The EDR is at the origin of the cylindrical coordinate system. The form of the iso-density lines is ellipsoidal, with a semi-major axis \( a \), and a semi-minor axis \( b \). The distance between the EDR and the center of the ellipsoidal density distribution is \( z_0 \). The parameters of the model are \( a \), \( b \) and \( z_0 \).}
\end{figure}
The governing equations for the plume expansion are the gas-dynamic continuity and momentum equations. The use of gas-dynamic equations to model the plasma expansion in the first nanoseconds after the start of the laser pulse is justified by the dominant effect of the pressure gradient in the plume expansion.\textsuperscript{138-141} When considering the two fluids continuity and momentum equations, the ions are the dominant contributors to the density $\rho$ and are displaced by the electric field generated by the surrounding electrons. It can be shown that with a quasi-neutral plasma with an electron temperature that equals the ion temperature the two fluids approach reduces to the gas-dynamic equations (see Section 5.2.3).\textsuperscript{140-142} The gas-dynamic equations assume that the fluid is in a state of continuum. This condition is verified in Section 4.3.2. The gas-dynamic equations are:

$$\begin{align*}
\frac{\partial \rho}{\partial t} + \text{div}(\rho \mathbf{v}) &= 0 \\
\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla) \mathbf{v} + \frac{1}{\rho} \nabla p &= 0
\end{align*}$$

\hspace{1cm} (10)

where $\rho$ and $\mathbf{v}$ are the pressure and velocity, respectively. The solution of Eq.(10) is found by similarity solutions with the ellipsoidal ansatz. After transforming the solution to a cylindrical coordinate system, which is the simplest form for the axisymmetric assumption, the density in the isentropic case $\rho_s$ gives:\textsuperscript{139,140}

$$\rho_s = \rho_0 C \left(1 - \left(\frac{z - z_0}{a_s}\right)^2 - \left(\frac{r}{b_s}\right)^2\right)^{\frac{1}{\gamma - 1}},$$

\hspace{1cm} (11)

and the density in the isothermal case $\rho_T$ is:\textsuperscript{141,138}

$$\rho_T = \rho_0 \exp\left[-\left(\frac{z - z_0}{a}\right)^2 + \left(\frac{r}{b}\right)^2\right],$$

\hspace{1cm} (12)

where $\rho_0$ and $C$ are constants; $a$ and $b$ are the two semi-axis of the ellipsoidal iso-density line that has a density of $0.3678 \rho_0$. The heat capacity ratio is chosen to be equal to $\gamma = 1.25$ from previous studies,\textsuperscript{138,140,143} and it is kept constant for all the analytic results. The value of 1.25 differs significantly from the value for ideal monoatomic gases of 5/3. Indeed, the translational degrees of freedom are augmented by the possibility of the particles to store/release energy through ionization/recombination. This leads to an expansion, during
which the density profile decays faster. With this formulation of the density the semi-major axis \(a\) can be estimated by the plasma density length scale, which is defined as
\[
L_s = c_s \tau_{\text{las}} = \sqrt{Z T_e/M_i} \tau_{\text{las}},
\]
where \(c_s\) is the ion sound velocity, \(\tau_{\text{las}}\) is the laser pulse duration, \(Z\) is the average charge state, \(T_e\) is the electron temperature, and \(M_i\) is the ion mass.\(^{142}\) \(a_s\) and \(b_s\) correspond to the semi-axis for the isentropic case of the outermost iso-density line where \(\rho = 0\); and \(z_0\) is the offset of the EDR region from the center of the ellipsoids. In order to have comparable results between the isothermal and the isentropic assumptions, the eccentricity of the ellipsoids in \(\rho_s\) and \(\rho_T\) is kept constant, hence \(a/b_s = a/b\).

The constants \(\rho_0\), \(C\), \(a_s\) and \(b_s\) are chosen in order to satisfy two conditions. The first one is to have an electron number density \(n_{e,\text{EDR}}\) in the EDR region, where the relation between the electron number density and the mass density is given by \(n_{e,\text{EDR}} = \rho_{\text{EDR}} Z_{\text{EDR}} M_i\); \(Z_{\text{EDR}}\) is the average ion charge in the EDR, and \(M_i\) is the tin ion mass. \(Z_{\text{EDR}}\) is estimated through the collisional-radiative (CR) model for the given electron number density and temperature in the EDR.\(^{144}\) The second condition is that the isentropic and the isothermal density distributions have the same number of particles in the inner core of the plasma up to the EDR (the volume integral of the number density is the same).

**Fig. 29:** The density profiles for the isentropic (\(\rho_s\)) and for the isothermal (\(\rho_T\)) case are shown for a cut along the axis of symmetry together with the temperature for the isentropic assumption (\(T_s\)). The EDR is located at \(z=0\).
4.2 Physical Processes and Analytical Models

The temperature profile is a constant line in the isothermal case with \( T = T_{EDR} \). Instead, in the isentropic case the temperature profile follows the density with \( T_s = T_{EDR} \left( \rho_s / \rho_{EDR} \right)^{1/\gamma} \).

The density and the temperature profiles that satisfy the above-mentioned assumptions are shown in Fig. 29, where a cut along the symmetry axis is presented. It is noticeable in Fig. 29 that the density profiles for the isothermal and the isentropic case with \( \gamma = 1.25 \) are very similar, with a maximum difference in the center of the distribution of around 10%. Therefore it is expected that differences in the opacity distribution will be mainly due to the differences in temperature profiles.

The Planck mean opacity in Eq. (8) can be rewritten in spherical coordinates leading to Eq. (13) for the isentropic, and to Eq. (14) for the isothermal case.

\[
\tau_s = \int \kappa_s \rho_s \ dl = a_s \left( \frac{T_{EDR}}{\rho_{EDR}} \right)^{1/\gamma} \left( \rho_{\delta} C \right)^{(\gamma+1)/2} \int \left( 1 - l^2 \left[ \left( \frac{\cos \theta - z_0/l}{a_s} \right)^2 - \left( \frac{\sin \theta}{b_s} \right)^2 \right] \right)_{0}^{1/2} \ dl. \tag{13}
\]

\[
\tau_s = \int \kappa_s \rho_s \ dl = a_s T_s \rho_0^{\gamma+1} \int \exp \left[ -(r_m + 1)^2 \left( \frac{\cos \theta - z_0/l}{a} \right)^2 + \left( \frac{\sin \theta}{b} \right)^2 \right] \ dl. \tag{14}
\]

Due to the power-law formulation of the Planck mean opacity, the two integrals can be solved with analytic functions. The solutions are given by the Gaussian or ordinary hypergeometric function \( \text{}_2 F_1 \) in the isentropic case, and by the error function \( \text{erf} \) in the isothermal case. The solution for the isentropic assumption is:

\[
\tau_s (l, \theta) = P \left( \left( 1 - R \cdot l^2 - S \cdot l - T \right)^0 \right) \ dl = P \cdot M \cdot \_2 F_1 \left( -Q, Q + 1; Q + 2; \frac{A + 2R \cdot l + S}{2A} \right), \tag{15}
\]

where \( M \) is a function of \( l \) and \( \theta \); \( R, S \) and \( T \) are unique functions of \( \theta \); and \( Q \) and \( P \) are constants shown as:
\[ M = M(l, \theta) = \frac{1}{R(Q + 1)^2} 2^{Q-1}(A + 2R \cdot I + S) \left( \frac{A - 2R \cdot I + S}{A} \right)^{-Q} (1 - R \cdot I^2 - S \cdot I - T)^Q, \]
\[ A = A(\theta) = \sqrt{4R - S^2 - 4R \cdot T}, \]
\[ R = R(\theta) = a^2 \sin^2 \theta + b^2 \cos^2 \theta \]
\[ S = S(\theta) = -\frac{2b^2 \cos \theta \cdot z_0}{a^2 b_s^2}, \]
\[ T = T(\theta) = z_0 b_s^2 \]
\[ P = a_m \left( \frac{T_{EDR}}{\rho_{EDR}} \right)^{\gamma} \]
\[ Q = \frac{r_n + 1 + s(\gamma - 1)}{\gamma - 1}. \]

The isothermal assumption leads to:
\[ \tau_\gamma(l, \theta) = M_T \cdot \text{erf} \left( \frac{(r_n + 1)(J \cdot l - z_0 b^2 \cos \theta)}{a \cdot b \sqrt{J (r_n + 1)}} \right), \]
where the \( M_T \) and \( J \) are a unique function of \( \theta \):
\[ M_T = M_T(\theta) = \sqrt{\frac{\pi}{2}} \left( a_m T_0 P_{\gamma+1} \right) \cdot \left( -\frac{1}{2} \frac{z_0^2 (r_n + 1) (J - b^2 \cos^2 \theta)}{J \cdot a_s^2} \right) \]
\[ J = J(\theta) = b^2 \cos^2 \theta + a^2 \sin^2 \theta. \]

By substituting Eq.(15) for the isentropic or Eq.(17) for the isothermal case in Eq.(8) an analytic expression for the integrated transmission distribution as a function of the angle \( \theta \) is derived. This solution has two degrees of freedom, which are the shape of the plume \( b/a \) and the distance from the center of the ellipsoid to the EDR \( z_0 \).

### 4.2.2 Proof of the Models with Slab Targets

In slab targets, the emission ends at \( \theta = 90^\circ \) from the axis of the laser for geometrical reasons, while the emission from the droplets continues further than \( 120^\circ \) (Fig. 27). Between 0° and 90° the EUV emission decreases by around 10% in the droplet case (shown in Fig. 27 and Fig. 31). Instead, with slab targets the decrease is around 50%.
over the same range. This large difference comes from the plasma distribution around the EDR. In a slab target, the plasma is constrained by the surface in its expansion, leading to denser plasma on the sides of the irradiated region. Filevich et al.\textsuperscript{145} investigated the generation of this secondary plasma by simulating the hydrodynamic expansion with and without the ablation on the surface around the LPP due to high energy photons and to hot electrons from the central region. This secondary plasma in form of sidelobes affects the EUV transmission by locally increasing the opacity to EUV radiation escaping the EDR with a direction close to a parallel of the plane. In a droplet with a diameter smaller than the spot size of the laser (30\(\mu\)m vs. 80\(\mu\)m in this study) this does not happen, because there is no material at the sides that can produce the secondary plasma. This means that the energy, carried by hot electrons and photons, which produces the sidelobes in a slab target can propagate at larger angles for a droplet smaller than the laser spot. Therefore, the dominant factor that causes the different distributions of density and temperature in the plume is attributed to the ratio \(L_{\text{target}}/\phi_{\text{las}}\), where \(L_{\text{target}}\) is the dimension of the part of the target that interacts with the laser and with the plasma, and \(\phi_{\text{las}}\) is the laser spot size at FWHM. In the experimental study presented in this work \(L_{\text{target}}/\phi_{\text{las}} = 0.37 < 1\). Instead, in the case of planar slabs \(L_{\text{target}}/\phi_{\text{las}} > 1\), causing the production of sidelobes.

**Fig. 30:** EUV emission versus the angle from the laser axis for the slab target case. The experimental data come from the work of Ando et al.\textsuperscript{72}. The measurements and the results from
the models of the EUV transmission are plotted. The parameters used in the models are \(b/a = 1.2\), and an offset of \(z_0 = 1.55L\).

The generation of the secondary plasma is confirmed by the shape of the density profile of the plasma, which is calculated by calibrating the analytical functions on the experimental results from Ando et al.\(^{72}\) (shown in Fig. 30), where the used laser is a Nd:YAG at an irradiance of \(1 \times 10^{11}\) W/cm\(^2\). The constants of the models \(n_{EDR}, T_{EDR}\) are taken from Tao et al.,\(^{146,147}\) where the same irradiance was used. The parameter \(z_0\) is chosen to equal \(1.55L\), which is coherent with the profile of the EUV radiation in the space that was measured by Tao et al.\(^{147}\) The semi-axis ratio used to calibrate the models is of \(b/a = 1.2\), which means that the sidelobes are contained in the slightly flattened ellipsoidal shape of the density profile. This value is in agreement with the density maps derived from 532 nm and 266 nm interferograms by Tao et al.,\(^{147}\) which demonstrate that the analytic model is able to estimate the shape of the plasma from the angular distribution of the EUV emission.

4.3 Models Applied to Droplet Targets

4.3.1 Calibration of the Models with Experimental Data

With droplet targets, the constants of the model \(n_{EDR}, T_{EDR}\) are chosen based on the paper by Yuspeh et al.,\(^{96}\) who performed experiments with the same irradiance of \(2 \times 10^{11}\) W/cm\(^2\) as in this work. The laser irradiance rather than the pulse duration has been chosen as the matching parameter, because it has been shown that the conversion efficiency (CE) at an irradiance of \(2 \times 10^{11}\) W/cm\(^2\) is independent of the duration;\(^{10}\) and because by solving the CR model, the energy equation for incompletely ionized targets and small radiation losses becomes an expression with \(T_e\) as a function of the laser irradiance, and without a direct dependence on the pulse duration.\(^{24}\) \(n_{EDR}\) and \(T_{EDR}\) are the values at the peak of the EUV emission of \(1 \times 10^{19}\) cm\(^{-3}\) and 45 eV, respectively. The plasma density length scale \(L_s\), normalized by the laser pulse duration is used to define the semi-major axis \(a\). The density
profile is such that on the axis of symmetry it decreases of 63.21% at a distance $a$ from the position of the density maximum.

The experimental data points are shown in Fig. 31, where the results from the analytical functions are also displayed. The main result obtained from the model is the integrated transmission or $I_{\text{EUV}}/I_{0,\text{EUV}}$. The normalization of the integrated transmission to the experiments gives $I_{0,\text{EUV}} = 0.38 \text{ mJ/sr}$. The parameters found from the calibrations are $b/a = 0.8$ and $z_0/L_s = 1.55$ both for the isothermal and isentropic assumptions. The latter prediction of $z_0/L_s = 1.55$ is in agreement with the measurement of the position of the EDR in Yuspeh et al.,\textsuperscript{96} where the center of the EDR is located at around 190 $\mu$m from the droplet, and the plasma density length scale is $L_s = 125$ $\mu$m. The ratio of the two latter parameters results in $z_0/L_s = 1.52$.

Fig. 31: EUV emission as a function of the angle from the laser axis. The measurements and the results from the models of the EUV transmission are plotted. The parameters used in the models are $b/a = 0.8$, and an offset of $z_0 = 1.55L_s$. The amount of EUV energy emitted behind the target with respect to the laser (or $\theta > 90^\circ$) is 30% of the overall emitted energy.

The CE can be calculated by integrating the EUV emission from the calibration of the functions over the solid angle, leading to $CE_s = 1.33\%$, $CE_T = 1.3\%$ in EUV $\pm/2\%$ BW centered at 13.5 nm, where the subscripts $s$ and $T$ refer to the isentropic and isothermal
cases, respectively. Both the models, isothermal and isentropic with $\gamma = 1.25$, produce a very similar angular distribution of the EUV emission, and the closest assumption to reality could not be found.

In the hemisphere where the ML collectors are normally employed, from $0^\circ$ to $90^\circ$, the decay of the EUV emission detected by the EUV energy monitor is of 15%. The amount of EUV energy emitted behind the target with respect to the laser ($\theta > 90^\circ$) is 30% of the overall emitted energy. The models extrapolate the emission up to around $\theta = 150^\circ$. The steep decay of the EUV emission of 1.6%/° after $90^\circ$ is due to an increasing plume density in the region where the EUV ray passes, which causes an increase in the absorption.

Using the angular distribution of the in-band radiation, the integrated in-band energy for $\theta<90^\circ$ can be calculated and related to the measured value at $45^\circ$. The calculated solid angle to get the overall EUV energy directed towards the collector ($0^\circ < \theta < 90^\circ$) using the EUV emission at $45^\circ$ in mJ/sr is $1.98\pi$ sr. Therefore, the value at $45^\circ$ can be extrapolated with $2\pi$ sr to get the overall emission in the hemisphere towards the laser.

An anisotropic plasma expansion with $b/a < 1$, as the one found when calibrating the models with the experimental results presented in Fig. 31, is due to a larger pressure gradient in the axial direction than in the radial. This causes a larger expansion velocity in the axial direction, which is attributed to the fact that part of the plume exits the laser spot. Therefore, the laser energy does not deposit homogeneously in the whole plume, but locally close to the laser axis. It follows that the condition to have such an anisotropy is:

$$\frac{\phi_{\mathrm{las}}/2}{L_r} \approx \frac{\phi_{\mathrm{las}}/2}{\tau_{\mathrm{las}}},$$

where $\phi_{\mathrm{las}}$ is the diameter of the laser spot, $L_r$ is the characteristic expansion length in the radial direction, $c_s$ is the ion sound speed, and $\tau_{\mathrm{las}}$ is the laser pulse duration. The condition in Eq.(19) is satisfied in the experimental results presented here because of the relatively long laser pulse and small spot size. In this case, the laser spot rather than the plume dimension dominates the EUV source size. When the plume is irradiated homogeneously, the expansion is expected to be nearly spherical for the time scale of EUV generation. This
4.3 Models Applied to Droplet Targets

is valid under the assumption that a secondary plasma does not develop on the sides of the plume (this feature is discussed in Section 4.2.2).

4.3.2 The Continuum Assumption

In order to verify the continuum assumption, which is necessary for the application of the gas-dynamic equations Eq.(10), the gradient-length local Knudsen number ($Kn_{GLL}$) is calculated and shown in Fig. 32. The definition of $Kn_{GLL}$ is:

$$Kn_{GLL} = \frac{\lambda}{\left| \frac{dQ}{dl} \right|},$$

where $\lambda$ is the mean free path, $Q$ is a flow property taken here as the mass density and $l$ is a distance taken approximately along the line of steepest gradients in the flow properties (perpendicular to the ellipsoidal surfaces). There is not a clear boundary between the scenarios where the continuum assumption is valid and where the flow is rarefied. However, by defining the break-down $Kn_{GLL}$ as the value that indicates an error of the continuum approach of 5%, the threshold in question is located at $Kn_{GLL} = 0.05$. Below this threshold the flow is defined as continuum. Fig. 32 shows that $Kn_{GLL}$ is smaller than 0.05 up to $z = 0.44z_0$ of the treated domain (where the ellipsoids are centered at $z = -z_0$). At the limit where the condition fails the density is such that the absorption of EUV is negligible (from $0.44z_0$ to the outer domain border the decrease in the integrated transmission is 0.1%). As a consequence, the continuum assumption is valid to obtain the estimation of the overall EUV absorption.
Fig. 32: Gradient-length local Kn-number versus the axial position. The EDR is located at $z=0$. The value stays below 0.05 up to $0.44z_0$. The contribution to the EUV absorption of the region outside the continuum is 0.1%. The continuum assumption is valid to obtain the estimation of the overall EUV absorption.

4.3.3 EUV Transmission in Space

The models with the parameters found from the calibration with the experiments are used in Fig. 33 to qualitatively show the EUV integrated transmission as a function of the space around the target for the two assumptions. The integrated transmission is 100% in the EDR and decreases along the lines that propagate outwards. As in Fig. 31 the isentropic assumption shows a trend similar to the isothermal one, but the transmission loss is smaller from the EDR in the direction of the droplet. This is due to the temperature profile, which remains constant for the isothermal case, but increases with density in the isentropic one. The higher temperature leads to a smaller $\tau_{\text{EUV}}$ (in Eq. (9) $T$ is at the power of $r_m$, which is negative), and hence to a smaller transmission loss.
4.3 Models Applied to Droplet Targets

**Fig. 33:** *Spatially integrated EUV transmission. The laser comes from the right and the EDR is at the origin of the coordinate system. The EUV transmission is integrated in the space along lines that propagate outwards, the lines have the EDR as their origin. The isentropic case (above) shows a slower decrease in the transmission than the isothermal case (below) because of the larger temperature between the droplet and the EDR.*

**Fig. 33** shows also that the EUV radiation expanding towards the droplet (in negative z direction) is completely absorbed. This is known, and is due to the absorption from the region where most of the laser is absorbed (the critical number density for the laser wavelength $\lambda_L = 1064$ nm is $10^{21}$ cm$^{-3}$), which is optically thick to EUV. With a CO$_2$ laser the absorption between EDR and the target is expected to be less pronounced, because the critical density (of around $10^{19}$ cm$^{-3}$ for photons from CO$_2$ lasers) is not optically thick to EUV, leading potentially to a larger portion of EUV that escapes the plasma at $\theta > 90^\circ$, and to debris with lower kinetic energy.
4.4 Dependence on Laser Irradiance and Droplet Diameter

![EUV emission measurements as a function of the angle from the laser axis. The blue dots represent the base case with $D_{dr} = 30 \, \mu m$ and $I_a = 2 \times 10^{11} \, W/cm^2$; the green squares with $D_{dr} = 50 \, \mu m$ and the same irradiance as the base case; and the red triangles with $I_a = 0.5 \times 10^{11} \, W/cm^2$ and the same diameter as the base case.](image)

**Fig. 34:** EUV emission measurements as a function of the angle from the laser axis. The blue dots represent the base case with $D_{dr} = 30 \, \mu m$ and $I_a = 2 \times 10^{11} \, W/cm^2$; the green squares with $D_{dr} = 50 \, \mu m$ and the same irradiance as the base case; and the red triangles with $I_a = 0.5 \times 10^{11} \, W/cm^2$ and the same diameter as the base case.

The measured EUV emission distribution for three different cases is plotted in Fig. 34. The mapping where 30 $\mu m$ droplets are irradiated with $2 \times 10^{11} \, W/cm^2$ is the base case.\textsuperscript{151} At the same irradiance but with 50 $\mu m$ droplets, the EUV emission decays faster at larger angles compared with the profile using the smaller droplets. In the analytical model, the faster decay was reproduced by increasing the density scale length $L_s$, hence by imposing a slower decay of the density profile on the space around the emission dominant region (EDR). A change in the density scale length for different ratios of droplet diameter to laser spot size (constant at a diameter of 80 $\mu m$ in this work) was already observed in experiments with droplets larger than the laser spot size.\textsuperscript{95} The change in $L_s$ indicates a difference in the expansion dynamics, from a more three-dimensional expansion with a short $L_s$ to a more two-dimensional expansion where the density decays more slowly and $L_s$ is longer. A longer $L_s$ implies that more EUV is absorbed around the EDR by the plasma, which is more opaque to EUV. Indeed, the opacity for tin was found to be proportional to
4.4 Dependence on Laser Irradiance and Droplet Diameter

The density at the power of 0.228.\(^{134}\) On the other hand, EUV emission close to the axis is stronger with the larger droplets (by approximately 1.5 times). The latter observation is justified by considering that the in-band emissivity increases with the density.\(^{152}\) The overall conversion efficiency (CE) increases by using 50 \(\mu\)m droplets from 1.3% to 1.4%. However, depending on where the EUV radiation is extracted, the local gain in using larger droplets is larger than 1 at angles from the laser axis of less than around 90°, and it is smaller than 1 at larger angles.

The EUV emission from 30 \(\mu\)m droplets irradiated at 5 \(\times\) \(10^{10}\) W/cm\(^2\) is also illustrated in Fig. 34. The emission measured by the energy monitor remains, in this case, nearly constant at around 0.24 mJ/sr until around 90°, where it begins to decay. The undisturbed emission at larger angles is due to the fact that the plasma around the EDR is less opaque to EUV compared with the base case. The opacity is proportional to the density and inversely proportional to the temperature.\(^{134}\) Since a higher temperature around the EDR with a lower irradiance is not expected, the main cause of change in the emission profile is attributed to the density. For this scenario in the analytical model the density scale length was changed, and in this case reduced, causing faster decay of the density around the EDR, hence a more transparent plasma. The change in \(L_s\) is in agreement with previous measurements: for instance, comparing Ref. 97 with Ref. 95, \(L_s\) increases from 115 \(\mu\)m to 135 \(\mu\)m for an increase in irradiance from 2 \(\times\) \(10^{11}\) W/cm\(^2\) to 6 \(\times\) \(10^{11}\) W/cm\(^2\). The lower emission at angles close to the laser axis is instead due to the lower production of EUV, which is the opposite of what happens with the larger droplets, caused by the lower density in the EDR.\(^{152}\) The conversion efficiency (CE) for the latter case with an irradiance of 5 \(\times\) \(10^{10}\) W/cm\(^2\) was 2.5%, which is approximately double the CE calculated for the base case at 2 \(\times\) \(10^{11}\) W/cm\(^2\). The increase in CE for the lower irradiance is due to the reduction in overheating of the plasma on the one hand,\(^{96}\) and to the lower absorption of EUV around the EDR on the other hand.
Fig. 35: Ratio of the photon counts at selected wavelength for Sn ions emission. The green squares represent the ratio between the photon counts of plasma generated by irradiating a 50 μm droplet versus 30 μm with $2 \times 10^{11}$ W/cm$^2$. The red triangles represent the ratio between the photon counts of plasma generated by irradiating a 30 μm droplet with $5 \times 10^{10}$ W/cm$^2$ versus $2 \times 10^{11}$ W/cm$^2$.

It is interesting to point out that the dependence of the irradiance on the CE was already measured for shorter pulse durations of 1.2 ns, and the same qualitative conclusion holds for pulse durations of up to 40 ns long, as in this work. The similar behavior for the different laser pulse durations is explained by the fact that the expansion process is comparable, i.e. during irradiation the plasma expands isothermally. This means that the time scale for the heat diffusion through the electrons is faster than the time scale for the expansion, and it is strong enough to equilibrate the plasma around the EDR in the scenarios considered in this work.

The different density between the three cases is also confirmed by the measurements performed with the spectrometer shown in Fig. 35. In the figure, the ratio of photon counts for the variations with respect to the base case (with $D_{dr} = 30 \mu m$ and $L_e = 2 \times 10^{11}$ W/cm$^2$) are plotted for selected wavelengths of tin ions. The photon count measured by irradiating 50 μm droplets is around $5/4$ of the one with 30 μm droplets, which is consistent with the increase in density and the emissivity predicted by the model. Similarly, the lower laser irradiance causes a reduction of around 40% in the photon counts.
4.4 Dependence on Laser Irradiance and Droplet Diameter

The differences in plume density are produced by differences in the amount of mass ablated from the droplet. In addition, considering scenarios where $D_{dr}$ is smaller than the spot size (as in this work), the ablated mass flow is proportional to $D_{dr}^2$, i.e. to the cross section of the irradiated area. In the first order approximation considered here, the possibility of droplet spallation as a function of droplet diameter during EUV generation is neglected.

Where the droplet diameter is larger than the spot size, already the subject of a number of publications, the scenario is complicated by the generation of secondary plasma next to the main spot, known as sidelobes. The latter contribute to the ablation of mass, and were studied for the slab target as limiting case in Ref. 145. Because of the presence of sidelobes around the main irradiated site, the ablated mass flow from droplets larger than the spot size is no longer expected to be linearly dependent on $D_{dr}^2$ or to the laser spot size squared.

Together with the droplet diameter, the ablated mass flow $\dot{m}_{abl}$ also combines the dependency with the laser irradiance, as the following relation describes:

$$\dot{m}_{abl} \approx 3 \times 10^3 \left( \frac{I_a}{10^{15} \text{W/cm}^2} \right)^{5/9} \left( \frac{\lambda_L}{1 \mu \text{m}} \right)^{-4/9} Z^{3/8} \left( \frac{\pi}{4} D_{dr}^2 \right),$$

where the proportionality with $I_a$, the laser wavelength $\lambda_L$, and the atomic number $Z$ and the constant were taken from Burdt et al.; $D_{dr}$ is in cm; and $\dot{m}_{abl}$ is in g/s. By calibrating the analytical model with the measurements, it is found that $\rho_{EDR} L_s^3 \varepsilon_{EUV} \propto \dot{m}_{abl}$, as
shown in Fig. 36, where $\rho_{\text{EDR}}$ is the density in the EDR, and $\varepsilon_{\text{EUV}}$ is the in-band emissivity. The latter was calculated by comparing the overall transmission of EUV in the plasma with the measured value. The linear relations are useful when predicting the EUV emission distribution as follows. $L_s$ and $\varepsilon_{\text{EUV}}$ are estimated by $\hat{m}_{\text{abs}} = f\left(D_{\text{dr}}, I_s, \lambda_L\right)$, $L_s$ is then used to calculate the EUV transmission distribution, and this is multiplied by $\varepsilon_{\text{EUV}}$ to give the EUV emission distribution. Fig. 37 shows profiles of the EUV transmission distribution for different $L_s$, and for every curve the correspondent value of the droplet diameter or the laser irradiance is shown.

Furthermore, the effects of changing the laser wavelength on the EUV emission distribution can also be estimated with Eq.(21). In particular, using a CO$_2$ laser ($\lambda_L = 10.6$ µm) would have a similar effect to reducing the diameter by more than 40%, leading to a plasma with a transmission greater than 90% up to around 120°.
4.5 Dependence on the Ambient Pressure

A further parametric study was performed to quantify the effect of the ambient pressure on the emission of EUV radiation. The aim was to understand whether the presence of ambient gas is important in the plasma expansion dynamics already during EUV emission or if the influence of the ambient gas on the emission comes mainly from the absorption of part of the radiation. Plasma was generated using droplets with a diameter of $D_{dr} = 30 \, \mu m$ and laser pulses with an irradiance of $I_a = 2 \times 10^{11} \, W/cm^2$ at three different ambient pressures, 0.17 mbar, 0.55 mbar, and 1.4 mbar. The background gas was for all the experiments argon. Fig. 38 shows the result of the experiments without taking into account the absorption of EUV from the argon ambient as in the post-processing of the previous Sections. The increase in pressure from 0.17 mbar to 1.4 mbar leads to a decrease in the EUV emission of more than 80%.

![EUV emission measurements as a function of the angle from the laser axis. The blue dots represent the base case at 0.17 mbar, the green dots 0.55 mbar and the red dots 1.4 mbar. For all the three emission measurements the ambient gas was argon. Absorption of argon was not taken into account in the post-processing.](image)

The influence of the EUV absorption in the ambient gas was removed from the measurements with the data of the transmission of EUV in argon atmosphere. In all the
measurements the path between plasma and photodiode was 385 mm in total, 185 mm from plasma to the instrument aperture and 200 mm inside the instrument. Fig. 39 shows the result of the absorption removal at the three different ambient pressures. Thus, the EUV emission at the irradiation site (IS) is estimated. The emission distributions at the different pressures converge in one, the match is particularly good between 0.17 mbar and 0.55 mbar. At the largest pressure the noise, which is amplified together with the EUV emission is too large to clearly distinguish a profile, however, the data points are spread around the profile of the other pressures.

Fig. 39: EUV emission measurements as a function of the angle from the laser axis. The colors of the dots represent mappings at different pressures. The absorption of EUV was removed. The data are representative for the EUV energy just outside the plasma, which is approximately constant for the different pressures.

The range of pressures was extended at a constant angle of 60° from the laser axis at a minimum pressure of $1.5 \times 10^{-2}$ mbar. Fig. 40 shows the measurements together. All the measurements were normalized by the EUV transmission at the given ambient pressure. The transmission of EUV in Ar was larger than 80% for the measurements up to 0.17 mbar, and then decayed at 12% for 1.4 mbar. The EUV mean energy in the EUV +/-2% centered at 13.5 nm produced at IS remains around 0.34 mJ/sr for all the pressures from $1.5 \times 10^{-2}$ mbar to 1.4 mbar. The error bars becomes larger with the pressure, as they are
4.5 Dependence on the Ambient Pressure

Scaled together with the EUV mean energy, in particular for 1.4 mbar the magnification of the EUV mean energy was of one order of magnitude.

![Graph showing EUV mean energy vs. pressure](image)

**Fig. 40:** EUV mean energy in the EUV +/-2% BW centered at 13.5 nm at IS. The measurements were normalized by the transmission of EUV in the ambient gas, in order to estimate the emission at IS. The angle from the laser axis was 60°, and the ambient gas was argon. The measurements remain around the same value of 0.34 mJ/sr.

The match of the profiles of EUV energy emission at the IS shows that the EUV emitted energy is not changing significantly with the increase of the pressure in the regime studied. Therefore, the EUV distributions exposed in the last Sections are scalable with the transmission of EUV in the ambient gas to find the distributions at different pressures. Additionally, as the EUV distribution is linked to the density of the plasma around the EDR (see Section 4.2), the constant EUV emission at the IS at different pressures means that the plasma dynamics in the region opaque to EUV is not being affected significantly by the ambient gas during EUV emission at least up to ambient pressures of 1.4 mbar.

In order to visualize the region of the plasma plume that has an impact on the absorption of EUV from the plasma, the local decrease of EUV transmission (the gradient of the integrated EUV transmission along the propagation directions, which is shown in Fig. 33) is plotted in Fig. 41. On the upper half of the figure the result of the calibrated model with
isentropic assumption is shown, and on the lower half is the result of the isothermal model. The EDR is centered at the origin of the plot and the initial droplet location is at $-z_0$.

**Fig. 41**: Local decrease of EUV transmission coefficient. The contour map shows where the plasma absorbs most of the EUV radiation. The upper part of the plot was calculated from the calibrated model with isentropic assumption, and the lower one with isothermal.

Where the local decrease of EUV transmission goes to zero (shown by the dark blue contour) plasma does not absorb EUV radiation anymore. Indeed, plasma surrounding the EDR region absorbs EUV photons in a very confined region, mainly where the density is larger than around $0.1\rho_{_{\text{EDR}}}$, which is three orders of magnitude larger than the density of the argon ambient gas at 0.1 mbar. In the pressure range that was considered the compression of the plume, which has been documented to become important with increase in ambient pressure, is affecting the plasma at distances where the opacity to EUV is already negligible compared to the one immediately around the EDR. And, at larger ambient pressures the ambient gas itself would be the main absorber of EUV photons from the EDR to the detector or to the collector mirror.
4.6 Summary

A further prove that the EDR formation and the EUV emission is difficult to influence externally is reported in the work of Roy et al.,\textsuperscript{156} where no significant difference in CE is measured with the addition of a 0.5 T magnetic field at the plasma formation region.

4.6 Summary

The EUV spatial distribution with regenerative targets in the form of droplets was measured. EUV is emitted in EUV +/-2\% BW centered at 13.5 nm up to angles from the laser axis larger than $\theta = 120^\circ$ and the models applied to the measurements extrapolate the range of EUV emission up to around 150°. In the region where ML normal incidence collectors are employed, from 0° to 90°, the decay is 15\%. The amount of EUV energy emitted behind the target ($\theta > 90^\circ$) is about 30\% of the total emission and the overall in-band CE is 1.3\%.

The results from the analytical models are in agreement with the results from the slab targets, where the shape of the density profile of the plume is of $b/a = 1.2$ and the position of the EDR is $z_0 = 1.55L_s$. These parameters are consistent with experimental results from interferometry and EDR imaging. The flattened form is attributed to a ratio $L_{\text{target}}/\phi_{\text{las}}$ greater than 1, hence to the formation of sidelobes. Instead with the droplet target and laser spot size used in the base case, the ratio is $L_{\text{target}}/\phi_{\text{las}} = 0.37 < 1$ and no sidelobes can be produced. The shape of the density profile of the plume in the latter case is defined by $b/a = 0.8$ and $z_0 = 1.55L_s$. The position of the EDR zone is confirmed by EUV imaging and the elongated form of the plume is explained as being generated by an inhomogeneous irradiation of the expanding plasma, whereas the non-dimensional number that identifies the anisotropic expansion towards an elongated form is $\phi_{\text{las}}/2\tau_{\text{las}}c_i < 1$.

Changes in laser irradiance and droplet diameter influence the EUV emission distribution. Larger droplet diameters lead to a faster decay of the emission with the angular position, and to larger emission close to the laser axis. EUV emission with smaller irradiance, instead, has a slower decrease with the angular position, and smaller emission close to the axis. The two effects are explained by a change in the density scale length of plasma expansion. With larger droplets more mass is ablated, $L$, and $\epsilon_{\text{EUV}}$ are larger, and the opacity around the EDR
region is also larger. Smaller irradiance and longer laser wavelength ablate instead less mass, $L$, and $\varepsilon_{\text{EUV}}$ are smaller, and plasma around the EDR is more transparent, leading to the more constant EUV emission with the angular position. The dependence between $L$, and the ablated mass allows the estimation of EUV emission profiles with given irradiance, droplet size, and laser wavelength.

The influence of the ambient pressure on the EUV angular distribution comes mainly from the absorption along the path of the EUV rays between the EDR and the sensor or collector mirror. During EUV generation the ambient gas does not influence significantly the plasma dynamics at least up to 1.4 mbar.
4.6 Summary
Chapter 5

MATTER EXPANSION

Analogously to the previous Chapter, experimental results are presented first, followed by the description of the dominant physical process in LPP dynamics. Particular emphasis is given to the angular dependence and to the influence of the ambient pressure on the load of matter at distances representative for collection of radiation. Experiments were performed to measure the total ion population, and the contribution of the different ion species. The experimental results are completed with physical models and computations to consider the plasma dynamics from generation to acceleration and expansion up to distances four orders of magnitudes larger than the focal spot size of the laser, where experimental data was collected.

5.1 Experimental Results

In this Section the results of experiments performed with FCs and the ESA are presented. The first are used to quantify the dependence of the total ion flux on the pressure and on the ambient gas type; and the second are used to quantify the contribution of the ion species around the plasma at different angular positions and pressures.
5.1.1 Total Ion Load

The negatively biased Faraday cups measure the current due to ion charges impacting on the surface of the cup. The measured current is proportional to the mean ion charge and to the flux of ions arriving at the sensitive surface (Eq. (2)). A typical measurement is shown in Fig. 42, represented as time-of-flight (TOF) diagram. The ambient pressure was $2.6 \times 10^{-2}$ mbar, the ambient gas was argon, the laser irradiance was $I_a = 2 \times 10^{11} \text{ W/cm}^2$ and the droplet diameter was $D_{dr} = 50 \mu\text{m}$. The first peak of the curve is due to prompt electrons, and to the photoelectric effect caused by the high-energy photons, which arrive at the surface of the FC (the sampling rate on the digitizer card used in the measurements was not large enough to discriminate the two contributions). The peak was coincident with the EUV signal, and can be taken as reference for the beginning of the plasma expansion. After the first peak, the current detected by the FC goes to zero again, until the fastest ions arrive, which are gradually followed by the slower ones.

![Fig. 42: Ion current measured with a Faraday cup at 45° with respect to the laser axis, at a pressure in argon background of $2.6 \times 10^{-2}$ mbar. The blue shaded area shows the current due to ions, and the red shaded area is the contribution from prompt electrons and high energetic photons.](image)

At the same laser irradiance and droplet diameter, data from the FC and EUV energy monitor placed at 45° from the laser axis was collected at different argon ambient pressures ranging from $1.8 \times 10^{-2}$ mbar to $6.2 \times 10^{-2}$ mbar. In the data analysis, the EUV
measurements were used to filter the ionic signals with an EUV energy in the same range. The error bars represent the standard deviation calculated from the filtered measurements. The total ion charge $Q$ arriving at the FC is shown in Fig. 43, and is calculated as the time integral of the ion current, which is displayed in the blue shaded area in Fig. 42. $Q$ gives a measure of the amount of ions $N_i$ arriving at the sensitive surface of the FC (Eq. (2)). The amount of ions reaching the FC decays over the whole measured range up to $6.2 \times 10^{-2}$ mbar, which was the highest pressure at which ions could be detected.

![Fig. 43: Total ion charge detected by the FC at different ambient pressures. Argon was used as background gas. The FC was placed at 150 mm from the IS, at 45° from the laser axis, and the 50 μm diameter droplets were irradiated with $2 \times 10^{11}$ W/cm². EUV energy was used to filter the data within the same EUV energy range.](image)

The same measurements were taken also for helium, which has a high transmission coefficient in the EUV range, and nitrogen. The results are plotted in Fig. 44. The different ambient gases produce a different decay in the number of ions arriving at the FC. The behavior of nitrogen is similar to the one of argon, but helium is more transparent to ions. The measurement points are fitted with a linear function in the semi-logarithmic space, where the slope is the same for all the three gases. The same function is shown in the linear space in Fig. 43. With argon the maximum pressure at which ions were detected at 150 mm distance from the IS was $6 \times 10^{-2}$ mbar. With nitrogen the maximum pressure at
which ions are detected is extrapolated to be lower, around $4.5 \times 10^{-2}$ mbar, and with helium significantly higher, around 0.12 mbar.

![Graph](image)

**Fig. 44:** Total ion charge detected by the FC at different ambient pressures, and with different ambient gases. The FC was placed at 150 mm from the IS, at 45° from the laser axis, and the 50 µm diameter droplets were irradiated with $2 \times 10^{11}$ W/cm². EUV energy was used to filter the data at the same EUV energy range.

### 5.1.2 Contribution of the Ion Species

The contributions of the single ion species at different angles with respect to the laser axis and at different ambient pressures could be investigated using the ESA with differential pumping system as presented in Section 3.2.2. At every angle and ambient pressure ion traces were stored with different voltages applied to the electrodes of the electrostatic filter. In the experiments presented in this Section, 40 ion traces were stored at each of the 36 different voltage settings on the electrodes of the ESA, and this was repeated at every angular position and ambient pressure. The voltage difference between the electrodes of the electrostatic filter was automatically varied from 0 to 500 V. For every ion trace, the laser pulse and the EUV energy were also measured. The laser pulse was measured with a photodiode placed on the beamline, and was used as reference for the TOF calculation (the first peak used with FCs is not measured here, as there is no line-of-sight between detector
and plasma). The EUV energy was measured with the EUV energy monitor located at 60° from the laser axis, and was used to filter the signals, in order to consider only ion traces at the same range of EUV emitted energy.

![Typical ion traces from the ESA, plotted together with the laser pulse. The voltage difference between the electrodes was 240 V, the argon ambient pressure was 2 × 10⁻² mbar, and the pressure inside the ESA was 8 × 10⁻⁶ mbar. The ionic peak occurring first in time is due to Sn²⁺ (shown by the green area), and the second one to Sn⁺ (shown by the blue area).](image)

**Fig. 45:** Typical ion traces from the ESA, plotted together with the laser pulse. The voltage difference between the electrodes was 240 V, the argon ambient pressure was 2 × 10⁻² mbar, and the pressure inside the ESA was 8 × 10⁻⁶ mbar. The ionic peak occurring first in time is due to Sn²⁺ (shown by the green area), and the second one to Sn⁺ (shown by the blue area).

A typical ion trace is shown in **Fig. 45**. Around 10 µs after the laser pulse and plasma generation the first Sn²⁺ ions arrive at the detector, followed by Sn⁺ ions. The trajectory inside the electrostatic filter depends on the charge-to-mass ratio of the ions and on their kinetic energy as shown in Eq.(3). Indeed, an ion with a higher kinetic energy is less deflected by the electric field. The accuracy of the kinetic energy filtering was estimated in Section 3.2.2 to be around 3% of the absolute value. The population distribution in the kinetic energy space dN/dE_{kin} (where N is the number of particles of the specie in consideration) is obtained combining the data at all the voltage settings. The result is shown in **Fig. 46** for an argon ambient pressure of 2 × 10⁻² mbar, and at 45° from the laser axis. The distribution is approximated by three Gaussian functions, and the error bars illustrate the standard deviation of the measurements at the same conditions.
5.1 Experimental Results

**Fig. 46:** Distribution of the ion population with respect to their kinetic energy. The angle with respect to the laser axis was 45°, and the argon ambient pressure was $2 \times 10^{-2}$ mbar. The distribution is approximated by the sum of three Gaussian functions that are shown for the distribution of the doubly charged tin ions with the discontinuous lines.

The kinetic energy distribution of the single species is integrated in the kinetic energy space to obtain the overall number of Sn$^+$ and Sn$^{2+}$ ions emitted by the plasma and entering the ESA $N_i(Z)$, as shown in Eq.(22).

$$
N_i(Z) = \int_0^\infty \left( \frac{dN_i(Z)}{dE_{\text{kin}}} \right) dE_{\text{kin}}.
$$

(22)

$N_i$ is shown in **Fig. 47** for different angles from the laser axis at a constant argon ambient pressure of $2 \times 10^{-2}$ mbar, and in **Fig. 48** as a function of the argon ambient pressure at 45°, 90°, and 120° from the laser axis. The error is calculated propagating the single errors in the integration over the kinetic energy. The total Sn$^{2+}$ population decreases with larger angles from the laser axis. The Sn$^{2+}$ ions propagating at 120° and arriving at 150 mm from the IS are around one third of the amount of Sn$^{2+}$ ions at 45°. Instead, the population of Sn$^+$ ions remains approximately constant, showing a more isotropic expansion.
Fig. 47: Total ion population at different angles from the laser axis. The argon ambient pressure was $2 \times 10^{-2}$ mbar. The contribution of $\text{Sn}^{2+}$ decreases with the angle, and the expansion of $\text{Sn}^+$ ions is more isotropic.

Fig. 48: Total ion population at different argon ambient pressures. (a) At $45^\circ$ from the laser axis. (b) At $90^\circ$ from the laser axis. (c) At $120^\circ$ from the laser axis. The ambient pressure similarly affects both the ion stages.
5.1 Experimental Results

The effect of an increase in ambient pressure is a gradual loss of ions arriving at the instrument, 150 mm from the IS, as already measured with FC and shown in Section 5.1.1. The ambient pressure similarly affects both Sn$^{2+}$ and Sn$^+$ at the different angles.

**Fig. 49:** Mean kinetic energy at different angles from the laser axis. The argon ambient gas pressure was $2 \times 10^{-2}$ mbar. The mean kinetic energy of Sn$^{2+}$ decreases with the angle, instead Sn$^+$ shows a more isotropic behavior.

The mean kinetic energy $\bar{E}_{\text{kin}}$ for the two ion species was calculated as the average of $E_{\text{kin}}$ weighted with the ionic distribution, shown by the following equation.

$$
\bar{E}_{\text{kin}}(Z) = \frac{\int_{0}^{\infty} E_{\text{kin}} \left( \frac{dN_i(Z)}{dE_{\text{kin}}} \right) dE_{\text{kin}}}{N_i(Z)}.
$$

The dependence of $\bar{E}_{\text{kin}}$ on the angle from the laser axis is shown in Fig. 49. Similar to the amount of ions, also the kinetic energy is changing with the laser axis for Sn$^{2+}$, and remains close to isotropic for Sn$^+$. More Sn$^{2+}$ ions with a larger kinetic energy are directed towards the laser, followed by an almost isotropic cloud of Sn$^+$ ions. Analogous to the total ion population, the dependence on the pressure is depicted in three plots at different angles from the laser axis in Fig. 50. The increase of ambient pressure does not significantly change the mean energy of the ions for both Sn$^+$ and Sn$^{2+}$. The mean ion energy of both
ion species is enough to sputter a Si, Mo or Ru surface, as the energy threshold to sputter at normal incidence is of 83 eV, 46 eV and 44 eV, respectively.\textsuperscript{157} For example, the sputtering yield (number of atoms of the surface sputtered by the incoming debris ions) of Sn$^-$ ions impacting on a Ru surface decreases from around 3 to around 1 for energies from 1.3 keV to 400 eV.\textsuperscript{157,158}

Fig. 50: Mean kinetic energy at different argon ambient pressures. (a) At 45° from the laser axis. (b) At 90° from the laser axis. (c) At 120° from the laser axis. The ambient pressure between 1.5 x 10\textsuperscript{-2} mbar to 5 x 10\textsuperscript{-2} does not significantly affect the mean kinetic energy.

5.2 Physical Processes and Models

The dominant physical processes, and the developed models are presented in this Section. The aim of the models is to explain the experimental results considering the phenomena in
laser-matter interaction, in laser-plasma interaction, and in plasma expansion, which are dominant in the regime that is experimentally studied. As the measurements are taken at large distances from the IS (four orders of magnitude larger than the focal spot size), different stages in the plasma generation and dynamics have to be considered with particular emphasis on the spherical geometry of the target. The main physical processes are: ablation (laser-matter interaction) and ionization of the spherical target, absorption of the laser radiation (laser-plasma interaction), acceleration and expansion of the charged particles.

5.2.1 Ablation and Ionization of the Spherical Target

In LPP the ablation process is the initial step in the generation of plasma from the target. The laser photons are partially absorbed by the target, vaporize and eventually ionize the material creating plasma. In metallic targets, the electrons of the conduction band absorb the laser light, and once the electrons are heated the electron-phonon coupling transfers the energy to the lattice. During this phase, a tin surface reflects around 70% of the incoming radiation at 1064 nm. Once enough energy is absorbed by the surface layer, phase transition occurs, which depending on laser fluence, laser wavelength and material might be melting, boiling (explosive boiling if the temperature is larger than the thermodynamic critical temperature), and sublimation. Ablation is reached when the phase transition lead to a vaporization of the target material. The threshold for ablation for tin with a 1064 nm laser pulse is documented to be at a fluence of around 2 J/cm$^2$. As time passes and more energy is transferred, the vapor ionizes and plasma is formed on the surface of the target. The threshold for ionization for tin with a 1064 nm laser pulse is at a fluence of around 5 J/cm$^2$. At this point, plasma expands at velocities in the order of km/s and starts to interact directly with the laser pulse. At ionization the average expansion velocity of the forming plasma is expected to be approximately 3 km/s. Then, the expansion velocity rapidly increases with time as more energy is absorbed, and $T_e$ increases. When plasma with a number density equal to the critical density (around $10^{21}$ cm$^{-3}$ for 1064 nm photons) is formed, the laser beam is completely absorbed by the plasma, and does not arrive at the
surface anymore. However, the surface continues to be heated by the plasma itself. More details on the laser-plasma interaction are in the next Section.

![Sketch of the spatial laser profile together with a 30 µm droplet. The dimensions are scaled accordingly.](image)

**Fig. 51:** Sketch of the spatial laser profile together with a 30 µm droplet. The dimensions are scaled accordingly.

As the laser profile is not constant in time and space, the ablation process and the time needed to reach ablation is expected to be highly dependent on the position on the target surface. In addition, the target geometry is spherical, which complicates the dependence, as the normal of the surface \( \mathbf{n} \) irradiated by the laser is not parallel to the laser beam. The time needed for the ablation process to happen and the percentage of the surface being ablated is estimated using the laser profile determined in the experiments. The latter is shown in Section 2.2 to be approximated by a Gaussian function in space and time. In cylindrical coordinates the laser irradiance is expressed as:

\[
I(r,t) = I_0 \exp \left\{ -\left( \frac{t-t_0}{\sigma_t} \right)^2 - \left( \frac{r}{\sigma_r} \right)^2 \right\},
\]

where \( \sigma_t \) and \( \sigma_r \) are parameters of the laser pulse, \( t_0 \) is an offset introduced to have the laser beginning at \( t = 0 \), and \( I_0 \) is a constant. The latter is found knowing that the energy delivered by the laser per pulse is equal to the integral in time and space of \( I(r,t) \). The start of the laser pulse is defined when 0.5% of the maximum of the laser profile in time is reached.

The space dependent time integral of the local laser irradiance is calculated, which results in the space dependent fluence. The latter is then transformed from the frame of reference of
the laser spot to the one of the droplet surface to estimate the time needed for ablation and ionization, and the location on the droplet surface where the process of ablation starts.

5.2.2 Laser Absorption by the Plasma

After ablation and ionization, the plasma begins the expansion at maximum ion velocities in the order of 10 – 40 km/s.\(^{160}\) This regime is hydrodynamic, and the expansion is governed by the gas-dynamic equations (see Section 4.2.1).\(^{139,141,151}\) The plasma, rather than the surface of the target (as during ablation), is absorbing the energy of the laser in a region which is approximated by the isothermal state.\(^{136,137}\) The laser energy is primarily absorbed by electrons in the plasma (similarly as when the laser beam irradiates a metallic target). Electrons are accelerated by the electromagnetic field imposed by the laser pulse, and collide eventually with the ions around them, transferring part of the kinetic energy that they gained in the electromagnetic field and scattering away. The result of electron acceleration and electron-ion collisions is the transformation of laser energy into thermal energy of the plasma as the collisions increase the magnitude of random motion. The process is known as inverse Bremsstrahlung, and is the most important mechanism for the absorption of radiation in a plasma.\(^{161}\)

In a plasma with electron number density \(n_e\), electron temperature \(T_e\), and mean ion charge \(Z\), which is irradiated by photons with wavelength \(\lambda_i\), the coefficient for inverse Bremsstrahlung absorption \(\alpha_{IBA}\) is:\(^4\)

\[
\alpha_{IBA} = \frac{13.49}{\Lambda Z} \left( \frac{n_e}{n_c} \right)^2 \frac{\ln \Lambda}{\sqrt{1 - n_e/n_c}} \frac{1}{T_e^{3/2}},
\]

(25)

where \(n_c\) is the critical number density of the electrons in the plasma (normally called critical density), and \(\ln \Lambda\) is the Coulomb logarithm, which is typically around 5 for laser plasmas.\(^{143}\) The critical density is dependent on the wavelength of the incoming radiation, and identifies the electron number density at which electrons resonate with the electromagnetic field. In that state absorption is maximized, and the beam does not penetrate significantly further in the plasma. The critical density equals:\(^{142}\)

\[
n_e \left( \text{cm}^{-3} \right) = 1.1 \times 10^{21} \lambda_i^{-2} \left( \mu \text{m} \right).
\]

(26)
With a laser irradiation wavelength of $\lambda_L = 1064$ nm (Nd:YAG lasers operating at the fundamental frequency) $n_e$ is approximately $10^{21}$ cm$^{-3}$, instead with 10.6 $\mu$m (CO$_2$ lasers) $n_e$ is around $10^{19}$ cm$^{-3}$.

**Fig. 52:** Schematic of the absorption region in the ellipsoidal density profile with semi-axes $a$ and $b$. The dimension of the absorption region depends on the laser spatial profile and on the expansion of the plasma. The EDR of EUV radiation is also in the absorption region.

In order to estimate the location of the absorption region, calculations in the three-dimensional space were performed with the three-dimensional density profile derived from the mapping of EUV emission around droplet targets (Chapter 4). The absorption region, sketched in Fig. 52, was calculated as the region where 99% of the laser beam intensity is absorbed. The absorption of laser beam power density (irradiance) through inverse Bremsstrahlung in the ellipsoidal plasma profile is integrated along the propagation direction leading to:

$$
\frac{I_n(z,r,t)}{I(r,t)} = \exp\left(-\int_{\infty}^{z} \alpha_{BA} dl\right),
$$

where $I_n$ is the laser beam irradiance transmitted in the plasma, and $I$ is the initial beam irradiance that is Gaussian in space and time (Eq.(24)), and the laser pulse travel parallel to the axis of symmetry starting at $z \to \infty$ and directed towards the droplet surface centered at $z=0$. As the electron number density in the EDR region for EUV radiation is around $10^{19}$ cm$^{-3}$, which is smaller than the critical density of the plasma to 1064 nm wavelength photons, the EDR is expected to be in the absorption region. In the absorption region itself
the energy density that is being absorbed is not constant, but it increases strongly with the electron number density and decreases with the electron temperature ($\alpha_{IBA}$ is proportional to $n_e^2, T_e^{3/2}$).

The collisional-radiative (CR) model was included in the computation, in order to estimate the mean ion charge state needed in the calculation of $\alpha_{IBA}$. The CR model considers the rate of ionization and recombination due to radiative and to collisional processes to estimate the change in the population density of each species present in the plasma. The rate of change of the population of an ion specie can either be due to recombinations or ionizations:

$$\frac{dn_i(Z)}{dt} = n_e n_i(Z) S(Z-1, T_e) - n_e n_i(Z) \left[ S(Z, T_e) + \alpha_r(Z, T_e) + n_e \alpha_{3b}(Z, T_e) \right] + n_e n_i(Z+1) \left[ \alpha_r(Z+1, T_e) + n_e \alpha_{3b}(Z+1, T_e) \right].$$

(28)

where $S(Z)$ is the photoionization cross-section for ions with ion stage $Z$, $\alpha_r$ is the electron-ion radiative recombination rate, and $\alpha_{3b}$ is the three-body recombination rate. In radiative recombinations an electron recombines with an ion, and the excess of internal energy is released through a photon. Instead, in three-body recombinations, the excess in internal energy is removed by a second electron. The relations for the ionization and recombination coefficients are reported in a number of publications (see Colombant et al.$^{144}$ and Kolb et al.$^{162}$). Eq.(28) considers the change in ion population $Z$ over time as due to ionization of the specie $Z+1$ (first term on the right side of the equation), to ionization and recombinations of the specie $Z$ (second term), and to recombinations of the specie $Z+1$ (third term). As the time scale of the laser pulse is longer then the time scale needed to reach equilibrium in the ionic population,$^{89}$ the CR model is treated in its stationary limit (CRE model). The ratio between two successive ionization stages in stationary state is derived directly from Eq.(28):

$$\frac{n_i(Z+1)}{n_i(Z)} = \frac{S(Z)}{\alpha_r(Z+1) + n_e \alpha_{3b}(Z+1)}. \quad (29)$$

The mean ion charge is calculated using Eq.(29) for all the charge states $Z$ that might be present in the plasma. The CRE model has a wider application range in the temperature-density space compared with the local thermodynamic equilibrium (LTE) model and the
corona equilibrium (CE) model, but it has still limitations due to the following assumptions. The assumption of the CRE model are a Maxwellian velocity distribution of the electrons (electron-electron relaxation time smaller than the electron heating time), a population of the ions with charge $Z+1$ that does not change significantly during the time needed to establish the equilibrium with the ion population with charge $Z$, and that the plasma is optically thin to its radiation. In the absorption region of the LPP, the first two assumptions are valid as the characteristic time for thermalization is much shorter than the laser pulse duration. The last assumption is not valid everywhere and for every photon wavelength. Indeed, in the region where most of the laser radiation is absorbed the density is high enough to have a plasma opaque for example to EUV. However, the CR model was already used for the estimation of the mean ion charge in LPP from solid state lasers with reliable results. A finite opacity has the effect of suppressing the radiative decay, hence the results of the CR model could underestimate the mean ion charge in the denser regions, leading to an underestimation of the local absorption of the laser radiation.

**Characteristic Length Scale for Expansion and Heat Transfer**

In order to understand the role in the plasma of conduction, storage of the laser energy, and cooling from expansion, a dimensional analysis is conducted in this Section. Two length scales can be distinguished: an expansion length scale $l_s$ and a heat diffusion length scale $l_h$. In regions where $l_s > l_h$ the thermal diffusion in the plasma is slower than the expansion, and the cooling of the plasma through expansion is dominant (isentropic expansion as limiting case). Instead, the opposite case means that the expansion is slow enough to accommodate thermal differences within the plasma. The latter scenario includes isothermal expansion as limiting case, but the more general condition is a temperature profile that is dependent on the local absorption of the laser radiation.

The characteristic length scale for the expansion $l_s$ is given by the sound speed $c_s$ in the absorption region:

$$l_s = c_s t_{c,i} = \sqrt{\frac{T_i Z}{M_i}} t_{c,i},$$

(30)
where $t_{c,i}$ is the self-collision time between the ions, in other words it is the time needed to substantially reduce any anisotropy in the velocity distribution and to reach a Maxwellian distribution of the kinetic energy. The general form of the self-collision time is derived by Spitzer,\textsuperscript{11} and leads to:

$$t_{c,i} = \frac{M_i^{1/2} (3k_B T_i)^{3/2}}{8 \times 0.714 \pi n_i e^4 Z^4 \ln \Lambda}.$$  \hfill (31)

The characteristic length for heat diffusion is calculated from the energy transport equation:

$$\frac{\partial e_{\text{int}}}{\partial t} + \nabla \cdot q_{\text{cond}} = \nabla \cdot (\Phi_a - q_{\text{rad}}),$$  \hfill (32)

where $e_{\text{int}}$ is the internal energy density, $q_{\text{cond}}$ is the heat flux, $\Phi_a$ is the portion of the laser flux absorbed by the plasma, and $q_{\text{rad}}$ is the heat lost via radiation. As $t_{c,i}$ is several orders of magnitude shorter than the pulse duration, and the length scales are shorter than the laser focal spot size, change of absorbed laser flux and of the radiation heat loss over the space tend to 0, and in the infinitesimal limit the thermal conduction is the dominant mechanism for energy transport. As the time constant is small, and as a first order estimation is needed, the equation is linearized leading to:

$$\Delta T \left( \frac{c_v}{t_h} - \frac{k_s}{l_h^2} \right) \sim \frac{q_{\text{ext}}}{l_h},$$  \hfill (33)

where $c_v$ is the specific heat at constant volume, $k_s$ is the thermal conduction coefficient (from the substitution of the $q_{\text{cond}}$ with the Fourier’s law for thermal conduction), $t_h$ and $l_h$ are the characteristic time and length of the heat diffusion, and $q_{\text{ext}}$ is the heat flux being absorbed by the plasma minus the heat flux that the plasma loses through radiation. In plasma the thermal conduction coefficient was estimated by Spitzer to be:\textsuperscript{11}

$$k_s = 32 \left( \frac{2}{\pi} \right)^{1/2} \frac{T_e^{5/2}}{2 \pi Z e^4 m_e^{1/2} \ln \Lambda},$$  \hfill (34)

And the specific heat for an ideal plasma is:\textsuperscript{143,168}

$$c_v = \frac{3}{2} k_B n_i (1 + Z),$$  \hfill (35)
From Eq.(33), during plasma heating with a positive external heat flux, the term in parenthesis should be larger than zero, leading to:

\[ l_h > \frac{k_s}{c_v} t_h \approx \frac{k_s}{\sqrt{3/2 k_B n_i (1 + \bar{Z})}} t_h. \] (36)

The inequality is equivalent to have a Fourier number \( F_o \) smaller than 1, hence a larger energy storage rate compared with dissipation through conduction (the temperature increases). Eq.(30) and Eq.(36) are combined to have the condition for plasma where the heat diffusion is fast enough to dominate over the expansion during heating:

\[ \sqrt{\frac{k_s}{3/2 k_B n_i (1 + \bar{Z})}} t_h \geq \sqrt{\frac{T_c \bar{Z}}{M_i}} t_{c,i}, \] (37)

and the physical times \( t_h \) is taken equal to \( t_{c,i} \) for comparison.

The relation for the specific heat used in Eq.(35) was chosen to be in agreement with the thermodynamic analysis of an ideal plasma.\(^{143,168}\) Another way to estimate \( c_v \) is to use \( \gamma \). The ratio of the specific heat that is experimentally used is not the one for an ideal monoatomic particle, but is instead lower (\( \gamma = 1.25 \), instead of \( \gamma = 1.667 \)),\(^{140,143}\) an estimation for \( c_v \) can be found through the degrees of freedom:

\[ f = \frac{2}{\gamma - 1}, \] (38)

where \( f \) are the degree of freedom, and \( f(\gamma = 1.25) = 8 \), which is larger than \( f(\gamma = 1.667) = 3 \) for the case of the ideal monoatomic gas. Indeed, plasma has more degrees of freedom that can absorb and release energy (e.g. through ionization and recombination) compared with an ideal monoatomic particle where translation in the three dimensions is the only mechanism. A relation of the specific heat closer to experiments should be \( c_v = 8/2 k_B n_e \).

The comparison with Eq.(35) would lead to a mean ion charge of 1.667 to equal the two expressions. Therefore, the estimation of \( \gamma = 1.25 \) seems appropriate for regime were recombinations already reduced the mean ion charge, but close to the absorption region, were the mean ion charge is larger, \( \gamma \) could be overpredicted.
5.2.3 Ion Expansion and Acceleration

The gas-dynamic model (Eq.(10)) is based on the assumption that the plasma is quasi-neutral during the expansion, and that the ion temperature is equal to the electron temperature. The model is derived from the two-fluid description of the plasma, which in turns comes from the more general Vlasov equation. The continuity and momentum equations of the two-fluid description are given in Eq.(39), where $n_j$, $u_j$ and $p_j$ are the number density, the velocity, and the pressure of the specie with mass $m_j$ and charge $q_j$.

$$\frac{\partial n_j}{\partial t} + \nabla \left( n_j u_j \right) = 0,$$

$$m_j n_j \left( \frac{\partial u_j}{\partial t} + u_j \left( \nabla u_j \right) \right) = n_j q_j \left( \vec{E} + \frac{\vec{u}_j \times \vec{B}}{c} \right) - \nabla p_j. \quad (39)$$

In LPP at the irradiance regime considered in this work, the effects of the magnetic field $B$ are not dominant and $B$ is usually neglected.$^{11}$ At larger irradiances (in the order of $10^{19}$ W/cm$^2$) the magnetic field imposed by the laser becomes important by exerting the ponderomotive force on electrons in the skin layer.$^{169}$ And, at the same irradiance range, $B$ is generated by current in the plasma, leading to a pinching of the underdense plasma.$^{170}$

The gas-dynamic equation is derived as follows. In the momentum equation of the electrons the inertial component is neglected because of the small mass of the electrons leading to:\footnote{142}

$$n_e e \vec{E} = -\nabla p_e, \quad (40)$$

and, substituting $E$ from Eq.(40) in the momentum equation for ions, under the quasi-neutral assumption, leads to:

$$m_j n_j \left( \frac{\partial u_i}{\partial t} + u_i \left( \nabla u_i \right) \right) = -\nabla p_e - \nabla p_i = -\nabla p, \quad (41)$$

where $p$ is the total pressure with ion and electron contribution. The momentum equation of the gas-dynamic model in the form of Eq.(10) is found with $m_j n_j = \rho$. The latter approximation is justified, as the contribution of electrons to the density is negligible due to the small mass in comparison with the ions (216937 times smaller for tin ions). The gas-dynamic continuity equation comes directly from the continuity equation of the ions,
multiplying with the ion mass, and neglecting again the electron contribution because of their small mass. The solution of the gas-dynamic equations depends on the assumption on the equation of state as discussed in Section 4.2.1.

The gas-dynamic model leads to a density and temperature distribution close to experimental results in the first phase of the plasma expansion and in the core of the plasma, where the quasi-neutral assumption is valid. For this reason, it could be used for EUV emission distribution and laser absorption. But, to model the ion expansion at further distances, one main phenomenon is missing. Indeed, the model alone would predict ion velocities close to the sound speed in the plasma, which is smaller than what is measured. For example, in the expansion once the temperature is 2 eV, the Sn$^{2+}$ sound speed would be 1.8 km/s, which is more than one order of magnitude smaller than what is measured experimentally.

The missing phenomenon is the electrostatic acceleration of the different ion species. The acceleration is due to an electrostatic field generated by the non-homogeneous motion of electrons and ions. It is interesting to consider the simple case of a plasma in a slab with quasi-neutral assumption. The number of ions $N_i^{\text{esc}}$ that would escape from the slab of a given temperature and number density is given by $N_i^{\text{esc}} / N_i = r_D / L$, where $r_D$ is the Debye length of the plasma. With number density and temperature present in the EDR at laser peak for an irradiance of $2 \times 10^{11}$ W/cm$^2$, $N_i^{\text{esc}}$ is in the order $10^{-5} N_i$ because of the short Debye length (electrons shield electrically the ion charges in a very short distance), and is therefore negligible. From this simple analysis it can be concluded that others are the dominant mechanisms that cause the motion of fast electrons leaving the plasma and creating the electrostatic field in the scenario considered in this work. Bulgakova et al. considers in particular laser absorption through inverse Bremsstrahlung, and three-body recombinations. Both the processes transfer energy to electrons, increasing their velocity. However, the dependence between ion acceleration and laser irradiance suggests that inverse Bremsstrahlung is the dominant process.

In inverse Bremsstrahlung absorption, electrons directly absorb the energy in terms of thermal motion. As a consequence, electrons are accelerated to velocities much larger than the ones of ions and the fastest leave the plasma volume, creating an electric field between ions and electrons and space charge that leads to repulsion between the ions. The result is
an acceleration of the ions proportional in magnitude to their charge. The process of electrons escape and acceleration is repeated a number of times during the expansion, leading to the ionic velocity profile measured experimentally. The acceleration due to the decoupling between ions and electrons is the main cause of the shift in the Maxwellian distribution of the velocity, also observed in Fig. 42. And, it is represented by the analytical function:

\[ f(E_{\text{kin}}) = A_m E_{\text{kin}} \exp \left[ -\left( E_{\text{kin}} - \mu \right)^2 / (2\sigma)^2 \right], \] (42)

where \( A_m \) defines the amplitude of the function, \( \sigma \) the width, and \( \mu \) the shift. The shifted Maxwellian distribution of ions caused by the electrostatic acceleration in LPP was also called Boltzmann-Coulomb-shifted distribution. In the construction of the fitting to the experimental data Gaussian functions were used instead of the more correct shifted Maxwellian distributions, because the distribution of the populations of ions in the kinetic energy space were too close and wide to allow also the fitting of \( \mu \).

Fig. 53: Development of the non-quasi-neutral region leading to the electrostatic ion acceleration. A first layer of escaping electrons (red) expands followed by a region of excess of ions (blue). In both regions quasi-neutrality is not valid anymore, and the electric field accelerates the charged particles. The quasi-neutral core is not in the center of the density distribution, as the absorption region is shifted towards the laser.
The development of the non-quasi-neutral region caused by the escaping electrons is shown in the schematic in Fig. 53. The quasi-neutral core is not located in the center of the density distribution, as the absorption region that mainly produces the escaping electrons is shifted towards the laser. The escaping electrons lead the expansion, and are followed by a region with an excess of ions. This binary structure can be periodically repeated in space; in the schematic only one period is shown for clarity. Both the regions (escaping electrons, and excess of ions) are not quasi-neutral anymore, therefore the simplification of the two fluid model applied to find the gas-dynamic equation is not valid.

The scenarios considered in this work are not in high vacuum, since the minimum pressure was $6 \times 10^{-3}$ mbar. The presence of a background gas influences the expansion dynamics, as the ambient gas exchanges momentum (deceleration of the expansion) and energy with the plume (cooling of the ablated mass). The range of ambient gas pressures that was used in the experiments is between the low pressure regime ($p_{\text{amb}} < 1.3 \times 10^{-2}$ mbar), and the moderate pressure regime ($p_{\text{amb}} < 2 \times 10^{-1}$ mbar) considered in Ref. 93. In the first regime the plasma expands hydrodynamically up to rarefaction, and then freely at a constant velocity (the plasma position is proportional to the elapsed time $t$); in the second one the plasma expansion is represented by a shock model, where the plasma position is proportional to $t^{2/5}$, and in a later stage by a drag model, where the plasma position grows asymptotically. However, during the first phase of the expansion when the laser is absorbed and EUV photons emitted (see Section 4.5), the ambient gas has a negligible influence on the plasma expansion. Only at around 40 ns ($5\tau_{\text{laser}}$) from laser irradiation a deceleration becomes visible in the moderate regime, as shown also in Section 4.5.

In order to accommodate the ion acceleration with the initial quasi-neutral region, and the possible change in expansion velocity of the plasma, a new model was developed. The latter is based on the gas-dynamic solution for density and temperature with the ions modeled as expanding from the boundary of the absorption region (found from the model described in Section 5.2.2) with the mean velocity measured experimentally (see Section 5.1.2). Ions are modeled as expanding from the border of the absorption region, as it is where fast electrons are generated during irradiation. Instead, the density profile is modeled in time as expanding at the sound speed until a minimum velocity $a_p$ is reached. Then, the expansion continues at $a_p$. The latter is determined in the calibration of the model with experimental
data. Without this condition the model would predict an expansion velocity that tends to 0 as the plasma cools during expansion. The calibration is introduced also to account for rarefaction and deceleration of the plasma because of the ambient gas.

The density of the plasma is modeled in time and space from the isentropic ellipsoidal profile with the semi-axes ratio \( b/a = 0.8 \) (as derived in Chapter 4). The density \( (\rho_0 \text{ in particular}) \) is adjusted in order to respect mass conservation, hence the integral of the density over the volume occupied by the plume \( V(t) \) is equal to the mass ablated by the laser up to the instant in time \( t \).

\[
\int_0^t \dot{m}_{abl} \, dt = \int_{V(t)} \rho(z,r,t) \, dV. \tag{43}
\]

The ablated mass is modeled with Eq. (21), and the density profile follows an exponential decay as in Eq. (12). Once the laser pulse ends, the source term vanishes and the density decays everywhere. The second condition to model the density in time is that the characteristic length of the plume increases with the sonic speed. In the ellipsoidal density distribution this is done with the growth of the semi-axes at the local sound speed. The latter is a variable in the calculation, as it is a function of temperature and of the mean ion charge \( (c_s = \sqrt{ZTe/M_i}) \), and changes with time and space. Once \( a_s \) is reached, the expansion velocity remains constant.

During expansion and cooling, the ion front experiences recombinations, which are modeled in a first phase using the CR model (see Section 5.2.2) hence considering both collisional and radiative recombinations. Then, as temperature decreases three-body recombinations (proportional to \( T_s^{-9/2} \)) become dominant over radiative recombinations (proportional to \( T_s^{-1/2} \)). The condition to have negligible contribution from radiative processes is:

\[
n_e \gg \frac{3 \times 10^{13} T_s^{3.75}}{Z}, \tag{44}
\]

where \( n_e \) is in \( \text{cm}^{-3} \), and \( T_s \) is in eV. Once the condition is true, three body recombinations dominate, and the rate of change of the mean charge state becomes:
\[ \frac{d\bar{Z}}{dt} = -8.75 \times 10^{-27} \frac{n_e^2 \bar{Z}^3}{T_e^{9/2}}. \] (45)

The recombination coefficient of the three body recombinations rapidly increase with the cooling of the plasma because of the proportionality with \( T_e^{-9/2} \), then decreases again due to the decrease in \( n_e \).

Recombinations not only change the mean charge state, but also release energy in the plasma. For each recombination, a part \( E^* \) of the ionization potential \( I \) is returned to the electron. Not all the potential energy \( I \) is transferred to the electrons as there are also radiative losses from line emission in radiative transitions. The gain of energy in the plasma is transformed in internal energy and expansion work, leading to the following energy balance:\textsuperscript{143}

\[ \frac{d\varepsilon}{dt} + p \frac{dv}{dt} = E^* \left( -\frac{d\bar{Z}}{dt} \right), \] (46)

where for a quasi-neutral plasma with equal ion and electron temperature, \( \varepsilon = 3/2(1+Z)k_B T \) is the particle internal energy, \( p \) is the pressure in the plasma \( p = N(1+Z)k_B T \), \( v \) is the specific volume \( v = 1/N \), and \( N \) is the number of particles. In the computation, at each time step the internal energy is updated with the number of recombinations calculated with the recombination model and with \( E^* \) calculated from:\textsuperscript{143,179}

\[ E^* = 1.1 \times 10^{-3} \bar{Z}^{2/3} n_e^{1/6} T_e^{1/12}, \] (47)

where \( n_e \) is in \( \text{cm}^{-3} \), and \( T_e \) is in eV. With the exposed processes, the ion front is computed up to distances representative for collector optics, and where the ion sensors were placed. The model results are used in particular to study the angular distribution of the ion charge state from the droplet targets.
5.3 Discussion on the Phases of Matter Expansion

5.3.1 Ablation and Ionization of the Spherical Target

When the laser starts to irradiate the target, the irradiance is still lower than the peak value, and where the normal to the surface is parallel to the laser axis the ablation begins. The ablated region expands at a rate smaller than the one for slab target, as a larger surface is irradiated due to the third dimension in space. Location and time needed for ablation of the droplet target surface is estimated in this Section, in order to quantify the impact of this process on the overall plasma generation. The time needed to ablate the droplet target is shown depending on the position on the target surface in Fig. 54 for droplets with a diameter of $D_{dr} = 30 \, \mu m$ and $D_{dr} = 50 \, \mu m$, and a laser focal spot at FWHM of $\phi_{las} = 80 \, \mu m$.

![Time for ablation](image)

**Fig. 54:** Time needed for the ablation on the surface of the droplet target with the laser pulse used experimentally with 24 ns duration at FWHM, $\phi_{las} = 80 \, \mu m$ focal spot size at FWHM, and a mean irradiance of $2 \times 10^{11} \, W/cm^2$. The view is perpendicular from the laser axis, as seen from the focusing lens. (a) 30 $\mu m$ droplet diameter, (b) 50 $\mu m$ droplet diameter.

The picture shows that after around 2 ns (8%$\tau_{las}$) from the beginning of the deposition of energy on the surface of the droplet surface, ablation starts to occur at the center of the spot, where the normal to the surface and the laser axis are parallel. The whole droplet receives enough energy for ablation after around 7 ns (30%$\tau_{las}$). The difference of the ratio
\( D_{\omega}/\phi_{\text{las}} \) is a slightly faster ablation for a larger \( D_{\omega}/\phi_{\text{las}} \) at the same radius, which is due to the increasing radius of curvature of the surface.

**Fig. 55:** The time-dependence of the ablated surface fraction is shown together with the laser pulse. Ablation starts at around 8\%\( \tau_{\text{las}} \), and the whole surface is ablated at 30\%\( \tau_{\text{las}} \), when the laser pulse is at around 5\% of its maximum power.

The time-dependence of the surface ablation is shown in **Fig. 55** together with the laser pulse. At 8\%\( \tau_{\text{las}} \), the surface starts to be ablated, and after around 20\%\( \tau_{\text{las}} \), the complete surface reach ablation. After complete ablation more than 99\% of the energy of the laser pulse is available for ionization and heating of the plasma. Therefore, the start of the ablation process is negligible in terms of energy consumption of the overall laser pulse. The scenario changes when the droplet is not aligned with the laser, as exposed in Section 6.2.1.

### 5.3.2 Laser Absorption by the Plasma

During the laser pulse, plasma expands and absorbs laser energy via inverse Bremsstrahlung. The region where the laser energy is absorbed at the peak of the laser irradiance for the case considered experimentally is shown in **Fig. 56**. The green lines show the contours that contain the ablated mass. The region of absorption is narrow with respect to the plasma size as the laser pulse has a long duration (23.9 ns at FWHM), and plasma has time to expand.
5.3 Discussion on the Phases of Matter Expansion

up to distances from the target that are in the order of hundreds of micrometers, thus larger than the focal spot size. This leads to an inhomogeneous source of thermal energy in the plasma, which causes the anisotropic expansion (see Section 4.3). The flow exiting the absorption region in lateral directions expands and cools down, as the heating from the laser energy absorption is no more present, and as the sound speed, which is proportional to the square root of the temperature, decreases. Therefore, the plume has a larger velocity along the laser axis than perpendicular to it. The condition for anisotropy with droplet targets smaller than the laser focal spot is mainly due to the laser focal spot versus the pulse duration (Eq. (19)).

Fig. 56: Region of absorption of the laser pulse at the maximum irradiance in time. (a) and (b) are at two different views. The droplet target is in the origin of the coordinate system. The green contours are isodensity lines, and identify the boundary where the ablated mass is contained. The blue lines correspond to the position where 99% of the laser radiation has been absorbed, and the red dot is the center of the EDR of EUV radiation from Chapter 4.

In the three-dimensional plots of Fig. 56, the center of the EDR of EUV radiation is shown with the red dot. The position was derived from the experiments on the spatial distribution of the EUV emission in Section 4.3. The position of the EDR is contained in the absorption region, however the local absorption at the center of the EDR is not at its
maximum, as shown in Fig. 57. In the figure, the local absorption normalized with the value at the center of the EDR is shown on the plane \( y=0 \). The maximum absorption of laser radiation happens at around \( 1 \) \( L \) from the droplet target, which is confirmed by experiments,\(^9\) and is around \( 1/2 \ L \) away from the center of the EDR.

![Fig. 57: Local absorption of the laser pulse at the maximum irradiance. The laser arrives from the right. The region of maximum absorption is at around 1 \( L \), in axial position, whereas the EDR (shown by the red dot) is centered at 1.55 \( L \).](image)

The fact that the EDR does not correspond to the maximum absorption of the laser is due to the absorption of EUV radiation by plasma with a number density close to \( n_e(\lambda_L = 1064 \text{ nm}) \) as shown in Fig. 41. Indeed, the maximum absorption of EUV from Fig. 41 is at around 0.9 \( L \), from the target surface, almost coincident with the peak in laser absorption, which is at 1 \( L \). Only once the plasma expands to densities of around \( 10^{19} \text{ cm}^{-3} \), EUV radiation escapes from the plasma, and is transmitted in a measure that depends on the region where the EUV photons pass through, as measured and discussed in Chapter 4. Then, after further expansion the plasma state in terms of temperature and density is not anymore suitable for EUV generation.\(^{152}\)

The shape of the contours of the local absorption is due to the combination of the ellipsoidal density profile with the Gaussian spatial laser profile. A larger \( \phi_{\text{las}}/(2\tau_{\text{las}}) \) is expected to lead to a broader absorption region in radial direction.
5.3 Discussion on the Phases of Matter Expansion

The characteristic length scales $l_s, l_h$ were calculated in the absorption region and it was found that in the whole absorption region $l_h > l_s$. Therefore the heat transfer is always fast enough to keep pace with the expansion of the plasma, and the temperature distribution is dominated locally by magnitude of absorption of the laser energy. The boundary with absorption of laser radiation equal to the one in the center of the EDR is taken as reference to start the expansion of the plasma as exposed in the next Section.

5.3.3 Ion Expansion and Acceleration

The complete calculation for ion front expansion is first calibrated, as it integrates different models for the three-dimensional expansion from the absorption region, for the charge state estimation, and for the energy conservation from the release of energy due to recombinations. After calibration, the model is applied to the results presented in Section 5.1.

Calibration of the Model

The model was calibrated with data from experiments performed with Langmuir probes in ALPS 2 facility by Gambino et al.\textsuperscript{180}. The experimental condition were the same as in the angular mapping using the ESA presented in Section 5.1. The probe was positioned at 82° with respect to the main laser beam direction and at different distances from the IS (20 mm, 30 mm, 40 mm). The experimental ion density was estimated from TOF signals in the ion saturation regime using the Bohm theory, which assumes a Maxwellian distribution of the velocity.\textsuperscript{181}

The comparison between the model and the experiments is shown in Fig. 58. The model predicts the decay of the ion number density within the measurement errors. The ion number density was chosen as physical quantity to calibrate and test the model, because it includes the effects of the expansion, and of the recombinations happening during expansion. And, the mean ion charge (not measured at the distances under consideration for the comparison with the ESA) was not needed as input in the analysis of the experimental data.
Fig. 58: Comparison between model and experiments. Experimentally the ion number density was measured with Langmuir probes and evaluated with the Bohm relation by Gambino et al.\textsuperscript{180} The angle from the laser axis in the experiments and in the computation was 82°.

The model was calibrated by adjusting the minimum value of the sound speed $a_\rho$. The result of the calibration is $a_\rho = 9.1$ km/s, which is used for all the model results presented.

**Angular Distribution of the Mean Ion Charge**

From the relative ion population shown in Fig. 47 the angular distribution of the mean ion charge was directly calculated as the weighted average of the two populations, and is shown in Fig. 59. The mean ion charge decays with the angle from the laser axis, from around $\bar{Z} = 1.9$ at 45° to $\bar{Z} = 1.6$ at 120°. The model for ion expansion, developed in Section 5.2.3 and calibrated in the previous Section, is applied to the angular distribution of the charge state as shown by the red line in Fig. 59. The model predicts the experimental results within the experimental error, and extends the measurement range showing a monotonic decay.
5.3 Discussion on the Phases of Matter Expansion

**Fig. 59**: Mean charge state versus the angle from the laser axis at $p_{\text{amb}} = 2 \times 10^{-2}$ mbar. The blue squares are the measurements, and the red line is the model, which predicts the decay of the mean charge with the angle within the experimental error.

The main cause of the anisotropy in the mean charge state is due to the fact that the region where energy is absorbed by the plasma is not centered with the source of mass. As shown in **Fig. 57**, the absorption reaches the maximum at $1 \ L_i$ from the source of mass (the droplet surface) on the side of the droplet irradiated by the laser. An isotropic mean ion charge distribution would be expected for a hypothetical explosion where the location of the energy source and of the mass source (droplet surface in the experiments and models) would be the same.

The fact that the maximum mean ion charge is around 2 and that larger ion stages were not detected is probably due to the presence of the background gas, as discussed in Section 5.4. It is interesting to compare the results with slab target experiments in the same range of irradiance and with tin as target material. The distribution of the ions from the slab target showed a similar decay with the angle for all the ion species measured (up to Sn$^{5+}$ as the condition was in high vacuum). And, the presence of a quasi-isotropic ion population with a lower charge is not observed. Therefore the measured distribution of the singly charged ions is a specific feature of droplet targets. The reasons might be the formation of the sidelobes already described in Section 4.2.2. Ions would also be impacted by the sidelobes as it is an additional region where recombinations might happen, leading to a modification.
of the ion charge profile. With droplet targets, instead, the plasma is free to expands in all direction, provided that:

\[ D_{dr} / 2 < c_s \tau_{las}, \]  

(48)
as in the case of the measurements performed in this work. And, the higher charge close to the laser axis is the evidence of localized laser absorption.

![Mean ion charge distribution](image)

**Fig. 60:** Spatial distribution of the mean ion charge state. The laser beam irradiates the plasma from the right. The mean ion charge is larger towards the laser, and decays monotonically at larger angles from the laser axis.

The mean ion charge from the model in a cut along the laser axis is shown in **Fig. 60.** The mean ion charge is larger at the absorption region, because of the heating from the laser, and the asymmetry caused by the shifted energy source compared to the mass source is visible from the beginning of the expansion.

**Effect of the Position of the Quasi-Neutral Region**
The model was used to study qualitatively the sensitivity of the position of the quasi-neutral region on the distribution of the mean ion charge by moving the quasi-neutral region along the laser axis. Thus, moving the border where the accelerated ions are
initialized in the simulation with a smaller or larger offset with respect to the center of the density distribution.

**Fig. 61**: Spatial distribution of the mean ion charge state. (a) The quasi-neutral region is shifted 1 $L_s$ towards the droplet surface, closer to the center of the density distribution, resulting in a more symmetric mean ion charge in the space. (b) The quasi-neutral region is shifted 1 $L_s$ towards the laser, away from the center of the density distribution, resulting in a more forward-peaked mean ion charge spatial distribution.

**Fig. 61(a)** shows the mean ion charge distribution with the quasi-neutral region shifted 1 $L_s$ towards the droplet surface, hence closer to the center of the density distribution. The result is a more homogeneous distribution of the ion population in space. **Fig. 61(b)** shows, instead, the opposite case, with the quasi-neutral region shifted 1 $L_s$ away from the center of the density distribution. The mean ion charge distribution is in this case more forward-peaked, with larger charges towards the laser. The larger mean ion charge along the axis is due mainly to a shorter residence time of the ions in densities were recombinations are important. Indeed, the rate of three-body recombinations is proportional to $T_e^{-3/2}$, which promotes recombinations at lower temperatures, but it is also proportional to $n_e^2$. And, the fast decrease of $n_e$ in the three-dimensional expansion implies a freezing of the mean ion charge.
As the quasi-neutral region is linked with the region of absorption of laser radiation, the shift of the quasi-neutral region is directly related to a shift of the absorption region. In reality, this could be reproduced with different laser photon wavelength, as the critical density would change. A shorter wavelength (for example from frequency-doubled or tripled Nd:YAG or excimer lasers) would lead to an absorption region shifted towards higher densities, and to a more uniform spatial distribution of the mean ion charge (Fig. 61(a)). Instead, a longer wavelength (for example from CO$_2$ lasers) would lead to a shift of the absorption region towards lower densities, and to a more forward-peaked distribution of the mean ion charge (Fig. 61(b)). And, the mean ion charge of ions expanding towards the incoming laser would have a larger mean charge for longer laser wavelength and a smaller one for the shorter wavelength.

**Ion Acceleration**

The acceleration of the ion species depends on their charge $Z$ and on the initial electron and ion temperatures (assumed equal) $T$ in the quasi-neutral region. A scaling law that account for the electrostatic acceleration in the ion final kinetic energy is:

$$E_{ion} = nZT + T.$$  \hspace{1cm} (49)

The physical meaning of $n$ is the number of times that the acceleration is repeated during the laser pulse. $n$ and $T$ were fitted with experimental data performed at conditions similar to the one of this work by Burdt et al., leading to $n=4.7$, and $T=50$ eV. The choice of the value for $T$ is in agreement with the isentropic model that estimates plasma with a temperature of around 50 eV in the absorption region, between the EDR region and the maximum absorption region (Section 5.3.2). The ion kinetic energy is measured in this work (see Section 5.1.2), and Eq.(49) is used in this Section to derive the ion charge during acceleration. As the distribution of the ion kinetic energy is broad in the measurements due mainly to recombinations, only the dominant Gaussian distribution is considered, and its mean kinetic energy is taken to calculate the ion charge during acceleration $Z_0$. The result is shown in Fig. 62 together with the amplitude of the dominant Gaussian for both Sn$^{2+}$ and Sn$^+$ arriving at the ESA.
5.4 Dependence on the Ambient Pressure

The spatial distribution of the ion charge has the same trend of what is observed directly at 150 mm from the IS (Fig. 47). Therefore, already during the acceleration phase highly charged ions are present towards the laser. Initial charges up to around 6 are in accordance with what was measured by Burdt et al.\textsuperscript{89}. The difference between the initial charge states calculated from the Sn\textsuperscript{2+} dominant population compared with the one from Sn\textsuperscript{+} is an indication that the two populations come from different waves of ions expanding from the quasi-neutral region. And, at larger angles from the laser axis the two waves become closer in space, up to the region at an angle from the laser axis larger than about 90°, where they are coincident. The fact that particles with initial charge from 6 to 3 are equally reduced to a double charge state is considered in Section 5.4 in relation with the background gas effects, and with the dependence on the pressure.

5.4 Dependence on the Ambient Pressure

The question of what happens to the matter expansion in the presence of ambient gas is important for the estimation of the damage of ionic and atomic debris on surfaces in terms of coating and sputtering. The experimental results presented in the previous Sections show
that at around \(5 \times 10^{-2}\) mbar the flux of ions at distances of 15 cm is not anymore measurable. Therefore, the question becomes whether no more particles from the laser plasma arrive at 15 cm, or if ions recombined along the way and became invisible to the detectors. A number of works have shown that the second scenario takes place, and are in agreement with evidences of coating during sample exposure performed at \(p_{\text{amb}} > 5 \times 10^{-2}\) mbar.

Different observations from experiments and from computations are used to analyze the influence of the ambient pressure on matter expansion. In particular, the following experimental data and their dependence with the ambient pressure are considered: the total collected ion charge, the contributions of the different ion species, and the spectra in the 200 nm – 1100 nm range. Computations are introduced and used to estimate the influence of the inertial collisions on the ion propagation in the ambient gas up to the sensor or collector position.

### 5.4.1 Experimental Results

The dependence of the ambient pressure on the total ion charge, and on the ion charge states is presented in Fig. 44, Fig. 48, and Fig. 50, respectively. These experimental data are completed with the mean ion charge and with the spectra in this Section. The former is derived from the measurements presented in Fig. 48, and the latter were captured with the experimental setup exposed in Section 3.1.2.

The mean ion charge versus the argon ambient pressure is shown in Fig. 63, where also the EUV mean energy is shown. The EUV mean energy remained within the measurement error over the measurement range. The figure shows that the increase in ambient pressure does not have a large impact on the mean ion charge that remains between 1.9 and 1.8. However, the last point at 0.04 mbar has a large error (around 20%), and the mean ion charge could be lower. The measurement error becomes larger with the pressure due to the smaller amount of ions arriving up to the detector (see Fig. 48).
5.4 Dependence on the Ambient Pressure

**Fig. 63**: Mean ion charge and EUV mean energy as a function of the pressure of the argon ambient gas. The angle from the laser axis was constant at 45°. The EUV mean energy remains approximately at the same level in the pressure range, hence changes in the dynamics of the expansion during EUV generation is not expected.

The last experimental results used in the discussion are the ratios of selected lines of the spectra of tin and argon, which are plotted in **Fig. 64**. The ratio is calculated with respect to the spectrum at 0.17 mbar. In the lines of both the elements, the trend of the amplitude is the same: ascending with the pressure. In order to highlight the ambient pressure dependence, the mean of the count ratios at every pressure was calculated and plotted together in **Fig. 65**. The amount of radiation emitted by both tin and argon excited states first increases strongly with the pressure, and then flattens at pressure levels in excess of 0.1 mbar. And, in the range measured by the ion detectors, the sensitivity of the emission to the background pressure (the slope of the curve) is approximately at its maximum.
Fig. 64: Ratio of the photon counts of selected lines with contribution of argon, in the plot on
the left and tin, in the plot on the right. The different colors refer to different pressures, and show
a monotonic increase of the contributions with the ambient pressure of both Sn and Ar.

Fig. 65: Dependence of the emission of tin and argon relative to the argon ambient pressure. The
red continuous line is a guide for the eye.

5.4.2 Rarefied Expansion of LPP in Ambient Gas

Computational Setup

After acceleration, expansion, and recombinations the ion front propagates through the
ambient gas up to solid surfaces as the vacuum chamber, ion detectors, or ML mirrors. In
order to model this last step of the expansion, a particle in cell code coupled with a direct
5.4 Dependence on the Ambient Pressure

Simulation Montecarlo code (PIC-DSMC) was used in combination with a hydrodynamic code (POLLUX). The coupled code is called hydro-particle (HP) code, and was developed by the author of this work. The PIC-DSMC part of the code was validated with two different experimental cases. In its original form the HP code was restricted to applications in high vacuum, as the initialization of the background gas computational particles (CP) was limited due to insufficient virtual memory available at the nodes used in the simulations. The HP code was extended in order to include the background gas by Frick. The upgrade is based on a modified adaptation of the computational domain, and is called background-gas-hydro-particle (BGHP) code. In the original HP code, the computational domain was homogeneously enlarged as soon as tin ions were arriving at the border. Instead, only a section of the domain is considered in the BGHP, and the simulation is transformed from two-dimensional axisymmetric to two-dimensional coordinates as shown in Fig. 66. The reduced domain has the length \( l = 150 \) mm needed for comparison with experiments, and is inclined of an angle \( \alpha \) from the laser axis.

**Fig. 66:** Schematic of the computational domain during the simulation. The initial domain is enlarged as soon as tin ions arrive at the border in the normal domain adaptation shown with dashed lines. The reduced domain of length \( l \) and inclined of an angle \( \alpha \) versus the laser axis was used to simulate also the background gas.
The advantage of a reduced domain size is that the ratio between the number of tin CP and
the number of background gas CP is larger than the case with a larger domain where the
majority of the CP are of the background gas specie. Furthermore, the transformation from
2D-axisymmetric to 2D allows the simulation of a gas curtain entering on one side of the
domain perpendicular to the ion direction.

In the simulations, the velocity of the different ion species coming from the hydrodynamic
code POLLUX was adjusted to correspond to the experimental data and fed in the PIC-
DSMC part of the program.

Two kinds of collisions are treated in the BGHP code: Coulomb and inertial collisions. The
former are simulated with the PIC part of the code, and the latter with the DSMC part. The inertial collisions between charged and uncharged CP, and between uncharged
CPs are modeled with the variable hard sphere (VHS) model. The simulations were
tailored to quantify the influence of the inertial collisions during the expansion of tin ions
and atoms, therefore no mechanism was implemented for the recombinations of Sn
particles and for charge exchange collisions.

Results

The expansion of ions and atoms was simulated at different ambient pressures from 1.5 x
10^{-2} mbar to 0.12 mbar centered at an angle from the laser axis of 45°. Fig. 67 shows the
pressure of the argon ambient gas with the tin ions and atoms at 2.2 µs from the beginning
of laser irradiation. The argon population was initialized and maintained at 1.5 x 10^{-2}
mbar. The green dots are doubly charged ions that are faster, lead the propagation front,
and are followed by Sn^+ represented by the red dots. The atomic tin expands at a lower
velocity (mean velocity of around 5 km/s) and is still in the first 30 mm of the domain at
2.2 µs.
5.4 Dependence on the Ambient Pressure

Fig. 67: Position of tin ions and neutrals in argon ambient gas at $1.5 \times 10^{-2}$ mbar. The conditions (laser and droplet) were the same as in the experiments, the domain was reduced centered at 45°. The time elapsed from the plasma generation was 2.2 µs. The green dots are Sn$^{2+}$ CPs, the red ones are Sn$^+$ CPs, and the black ones are Sn CPs.

All the CPs exiting from the computational domain at $x > 0.15$ m were stored, and the amount of those particles is shown in Fig. 68 for the different simulations, and different species. The amount of particles was integrated in time up to 200 µs, in order to ensure that all the ionic CPs exited the computational domain. As in the simulations the only mechanism that was implemented was inertial collisions, the recombinations along the way do not transform the ion population in neutrals. However, the results are representative for the stopping capability of the background gas on the debris particles, as inertial collisions are the only way to deviate or stop debris particles expanding in the background gas. The background gas stops a different amount of debris particles depending on the pressure, i.e. the number of background gas particles in the domain. The particles with the smallest velocity (Sn) are those that are stopped more efficiently, followed by the particles with velocity of Sn$^+$, and of Sn$^{2+}$. The amount of particles at the different pressures deviates from a clean trend, because of the small amount of debris particles present in the domain, which is limited by computational time and virtual memory. Indeed, the number of CPs exiting the domain at the lowest ambient pressure was only around 100.
Fig. 68: Number of ions arriving up to 0.15 m from IS, where the ESA was placed, at different pressure of the argon ambient gas. The data is integrated in time up to 200 µs. The number of particles is normalized for each specie with the value at 1.5 x 10^{-2} mbar. The fastest particles (Sn^{2+}) are less affected by the presence of the ambient gas, instead the slowest atomic tin is stopped after around 6 x 10^{-2} mbar.

5.4.3 Discussion

Detector Sensitivity

The first aspect to consider is whether the sensitivity of the ESA is too coarse to measure a flux of ions significant as coating. The whole instrument is not calibrated, however the sensitivity is estimated based on the detector characteristics. The CEM used for the measurements had a gain (from the manufacturer) of \( G = 1.1 \times 10^7 \) at the operating voltage of 2.5 kV, and the signal was measured through a resistance of \( R_c = 1 \, \text{kΩ} \). The signals were not measured directly through the 50 Ω in order to amplify the signal without broadening it more than the features of interest, as with 1 MΩ. The signals had a typical noise level of 2 mV, which sets the minimum detectable signal \( U_{\text{min}} \). The minimum detectable charged particle flow is therefore given by:

\[
\left( \frac{dN_i}{dt} \right)_{\text{min}} = \frac{U_{\text{min}}}{G \gamma R_c},
\]

(50)
where $\gamma$ is the detection efficiency defined in Section 3.2.2. Substituting the values Eq.(50) with $\gamma=0.6$ (for a 500 eV singly charged ion accelerated in the 2.5 kV detector bias to 3 keV) results in a minimum detection threshold of $1.9 \times 10^6 \#/$s (charged particles per seconds). The minimum aperture of the ESA is of 2 mm at the inlet (150 mm from the IS), therefore the minimum particle flux is of $1.35 \times 10^{10} \#$/s$^{-1}$sr$^{-1}$. In perspective of coating, the time needed to cover uniformly a surface with one atomic layer of tin ions would be $2.6 \times 10^7$ s (304 days of continuous operation at 6 kHz and $1.6 \times 10^{11}$ plasmas). Therefore, the sensitivity of the ESA is enough to detect significant ion load for the coating of surfaces.

**Collisions**

At a constant pressure and constant ambient temperature the number of atoms or molecules present along the trajectory of the ions is approximately the same for all the gases. Therefore, comparing the total ion charge arriving at the FC at a constant pressure (Fig. 44) gives the efficiency of neutralizing the ion charge and/or of deflecting the ions from their initial trajectory per number of ambient gas particles. The particle species – tin ions from the LPP and ambient gas particles – interacts through collisions. The collisions between ions and ambient gas particles are at first not Coulomb collisions, as only ions are electrically charged. Collisions might be inertial and/or charge-exchange (CEX) collisions. Inertial collisions exchange momentum between the species, in this scenario ions have the largest momentum that is partially absorbed by the ambient gas. If the relative velocity is large enough, part of the energy can be absorbed by the ionization of one collision partner, leading to the release of an electron. And, in CEX collisions electrons of the outer shell can be exchanged during the event.

In the exchange of momentum the most important factors are the relative velocity between the encountering particles, the dimensions of the particles, and their masses. As the motion of the ambient gas particles is thermal, hence with random direction and of one to two orders of magnitude smaller than the ion speed, the relative velocity between the ions and the ambient gas particles is mainly given by the ion velocity. The dimension of the particles gives the collision frequency, and the mass is directly proportional to the amount of momentum exchanged per collision. The values for the background gases and for tin are
given in Tab. 1. Following the VHS model as in the computations shown in the previous Section, the total collision cross-section $\sigma_T$ is given by:

$$\sigma_T = \pi d_{12}^2 = \pi \left( \frac{d_1 + d_2}{2} \right)^2,$$

(51)

$$d = d_{ref} \left( \frac{v_{rel,ref}}{v_{rel}} \right) \eta,$$

(52)

where $d_1$ and $d_2$ are the diameters of the collision partners calculated with Eq.(52); $v_{rel}$ is the relative velocity, $\eta$ is the power law coefficient and the subscript $ref$ indicates the reference values. Therefore, the largest and heaviest ambient gas atoms are expected to stop more efficiently debris. Indeed, between the gases used experimentally, helium is the smallest and lightest and is the worst gas to stop debris as seen in Fig. 44.

<table>
<thead>
<tr>
<th>Covalent radius [pm]</th>
<th>Mass [amu]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Sn</td>
<td>He</td>
</tr>
<tr>
<td>140</td>
<td>28</td>
</tr>
</tbody>
</table>

Tab. 1: Covalent radius and atomic mass of the elements used as background gas.

Inertial collisions were studied computationally in Section 5.4.2, and resulted to be as expected dependent from the kinetic energy of the particles. In particular the inertial collisions are effective against slow Sn atoms coming from the vaporization of the droplet target, already with an ambient pressure of $6 \times 10^{-2}$ mbar the motion of the tin atoms was mainly through diffusion in the ambient gas. Instead, Sn$^+$ and Sn$^{2+}$ were still able to reach the end of the computational domain (0.15 m) with only a portion stopped or deviated by the ambient gas. Therefore, inertial collisions alone are not expected to explain the decay of the ionic flux measured experimentally.

In the CEX collisions that might take place in the expansion of tin in the ambient gas, the charge transfer is asymmetric, as it involves different atoms. Rapp et al. treat the asymmetric charge transfer analytically, coming to an approximation for the charge transfer
5.4 Dependence on the Ambient Pressure

cross-section, which is proportional to \( (v/\Delta I)^4 \) where \( v \) is the relative velocity in the encounter and \( \Delta I \) is the difference in ionization potential between the two species. Therefore, charge transfer at the same relative velocity with larger \( \Delta I \) becomes rapidly less probable. The values for ionization potentials of the most abundant tin species, and of the background gases is given in Tab. 2.

<table>
<thead>
<tr>
<th>Ionization potential [eV]</th>
<th>Dissociation potential [eV]</th>
<th>Metastable potential [eV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Sn</td>
<td>Sn'</td>
<td>Sn''</td>
</tr>
<tr>
<td>7.34</td>
<td>14.73</td>
<td>30.5</td>
</tr>
</tbody>
</table>

Tab. 2: Potential of ionization, dissociation and metastable excitation of the elements used as background gas in comparison with the ionization potential of tin.

It is interesting to point out that only tin ions with charge state larger than 2 receive enough energy from the recombination with one electron to ionize the argon or helium atom. Charge transfer with Sn' and Sn''\(^+\), which are the only species observed in the experiments, need energy to happen. Therefore, the contribution of CEX collisions in the tin plasma expansion could explain the almost constant mean charge state with increasing ambient pressure observed in Fig. 63.

**Recombinations**

Data from the spectrometer presented in Fig. 64 and Fig. 65 show that with the pressure both argon and tin become more excited. The increasing number of transitions from Sn' and Sn''\(^+\) is coherent with the decrease of ion flux arriving at the detectors. And, the presence of Ar' increases with the pressure. The fact that more Ar' ions are present, means that more electrons were at disposal of tin ions to recombine, and hence produce line emission and become invisible to the ion detectors. Argon ions could be produced by the fast electrons that cause the electrostatic acceleration, by photons emitted by the plasma, by collision with the high kinetic energy tin ions, or by CEX collisions. The ESA could not
detect Ar ions (which have the same mass-to-charge ration as Sn\(^{3+}\)) therefore the motion of Ar\(^+\) is most probably thermal, with a negligible centripetal drift velocity.

The region were recombinations become important is expected to be after the acceleration region, because the kinetic energy of Sn\(^+\) and Sn\(^{2+}\) ions remains approximately constant over the pressure range (see Fig. 50), hence the charge state during acceleration is also expected to stay constant.

**Fig. 69:** Schematic of the region of recombination with ambient gas together with the acceleration region. The recombination region is expected to surround the acceleration region, with little or no overlap, since the kinetic energy of the ions measured experimentally did not diminish significantly with the pressure increase, i.e. the acceleration process was not significantly affected.

Recombination with the electrons supplied by the ambient gas could also explain the different behaviors seen with the different background gases in Fig. 44. Indeed, the atom with the lowest ionization potential (hence the most generous electrons donor) is the one with the steepest decay of ionic signal, instead helium that is the most difficult gas to ionize (see Tab. 2) between those studied produce a slow decay with the ambient pressure. And, with helium the measured ion signal was the highest.

Thus, the mechanisms that are behind the decrease of the ionic flux measured experimentally are at least three: inertial collisions, which are responsible of a decrease of the ion flux of around 30% over the measurement range; CEX collisions that explain the
almost constant value of the measured mean ion charge of the tin ions; and, recombinations with the electrons donated by the background gas.

5.5 Summary

Several experiments were conducted to investigate the angular distribution of ions in LPP generated by droplet targets, and the effects of the background gas – atom/molecule and pressure – on the plasma expansion. The experimental results show an angular distribution characterized by a fast expansion of Sn\(^{2+}\) ions with a preferential direction towards the incoming laser beam, followed by a slower and quasi-isotropic Sn\(^+\) population. Since the data stored during the experiments is the result of multiple processes happening from the irradiation site up to the ion sensors, models have been developed to take into account the main processes: ablation, hydrodynamic expansion, laser absorption, electrostatic acceleration of the ions, and recombinations.

The energy needed to completely ablate the droplet target aligned with the laser beam is negligible, around 5% of the overall energy of the laser pulse. The largest amount of laser energy is absorbed behind the EDR, at around 1 \(L\), from the droplet surface, in a region opaque to EUV. The dimension of the absorption region is dominated by the ratio between the radius of the laser focal spot size and the duration of the laser pulse. Once calibrated with Langmuir probe measurements, the overall model predicts the angular distribution of the mean ion charge within the experimental error. The parameter that dominates the angular distribution of the ion charge is the position of the quasi-neutral region, from where the acceleration of electrons and ions starts. The model finds agreement with experimental results. The quasi-neutral region is coincident with the absorption region, which is expected to be the source of the electrons that lead to the ion acceleration.

A shorter laser wavelength, e.g. from excimer lasers, would move the absorption region closer to the droplet surface, creating a more isotropic distribution of the ion charges. Instead, a longer wavelength, e.g. from CO\(_2\) lasers, would move the absorption region further away from the target, creating a more forward peaked (towards the incoming laser) mean ion charge distribution. The angular distribution of the mean kinetic energy of the...
dominant Gaussian functions that contribute to the overall ionic population is used to estimate the charge state during acceleration, leading to mean charge states during acceleration between 5 and 2. The angular distribution of the charge states during acceleration is similar to the one measured by the instrument. It shows a forward peaked distribution for fastest ions with a larger charge state, and a more isotropic expansion for the ions that during acceleration have a charge state between 2 and 3, which are then detected as singly charged from the ESA.

The influence of an increasing pressure of the ambient gas is a decrease of the ion flux measured by the instruments. The mean charge state varies only of around 5% between \( 1.5 \times 10^{-2} \) mbar and \( 4 \times 10^{-2} \) mbar. The emission from the singly and doubly charged ions of tin and of the ambient gas (argon in the experiments) increases with the ambient pressure. The in-house BGHP code was used for the estimation of the influence of inertial collisions between background gas and tin ions. Results show that inertial collisions alone stops around 30% of the tin ions in the considered pressure range, and that fully stop atoms from \( 6 \times 10^{-2} \) mbar. Apart from the elastic collisions, the ion flux decreases with the ambient pressure because of CEX collisions with the highly charged ions, and because of recombinations with electrons donated by the background gas.
5.5 Summary
Chapter 6
NON-IDEAL DROPLET IRRADIATION

When droplets are not located in the focal spot of the laser beam, the irradiance is lower with respect to the spot, and when the displacement is not purely along the beam, the dynamic of the plasma becomes fully three-dimensional. The knowledge of what happens in this scenarios to EUV emission and debris load is important, in order to tailor the design and the operating condition of source components. And, as part of plasmas generated during an EUV source lifetime might be at non-ideal irradiation condition, the study of the source at off-design condition improves the estimation of the lifetime of source components.

6.1 Experimental Results

6.1.1 Correlation between Ion Load and EUV Emission
An array of FCs was used concomitantly with 2 EUV sensors, in order to correlate the spatial profile of ion expansion and EUV emission of each plasma. Details of the experimental setup are given in Section 3.2.1. During the experiment the operating condition of the irradiation system was at an irradiance of $2 \times 10^{11}$ W/cm$^2$ and the
repetition rate was 6 kHz. The droplet generation system was intentionally moved to reproduce non-ideal alignment between droplet targets and laser focal spot size. The ambient gas was argon at a pressure of \(3 \times 10^{-2}\) mbar.

![Graph showing ion charge flux versus time for four different FCs.](image)

**Fig. 70:** Ion charge flux versus time of four different FCs distributed around the plasma (see schematic in Fig. 15). The signals were measured synchronously.

**Fig. 70** shows an example of the measurement of the ion current of four FCs placed around the plasma at an angle from the laser axis of 45°. The signal acquisition was synchronized, so that all the signals were from the same plasma. In the particular case shown in **Fig. 70,** the main expansion of the plasma was closer to FC1 than to the other FCs (see schematic in **Fig. 15** for the location of the different FCs).

The data of around 400 plasmas are combined in **Fig. 71** to show the correlation between the main propagation direction of the plasma and of EUV radiation. The main propagation directions were calculated using the profile of ion expansion and EUV emission measured and exposed in the previous Sections to the data points in the space (+/45° and +/-60° for the ion distribution, and +/-75° for the EUV emission distribution). The ion main propagation direction correlates with the one of EUV close to 1:1. And, the maximum deviation from the laser axis of the two distributions was around 20°.
6.1.2 Correlation between Electron Load and EUV Emission

Electrons were measured with the same FC array used for ion measurements, but with a different bias setting. The bias in electrons measurements on the surfaces of the FCs was +60V. It has to be noticed that the signal from the electrons from the FC array is not expected to be as clean as the ones of the ions, since normal incidence FCs were used, and the voltage applied as bias was of +60V, insufficient to deflect all the ions with kinetic energies up to 1 keV. However, the signal gives still an indication of where the majority of the electrons are expanding.

An example of the measurements using the FC-array is shown in Fig. 72, where the ion currents from the FCs of the inner circle of the array at 45° from the laser axis are presented (as in Fig. 70). FC2 was the closest to the main propagation direction.
6.1 Experimental Results

**Fig. 72:** Electron flux versus time of four different FCs distributed around the plasma, and measured synchronously.

The correlation between the main propagation directions of electrons and EUV emission was calculated in the same way as for ions, and is depicted in **Fig. 73.** Electrons are less sensitive to the offset compared to ions. Indeed, the maximum deviation of the main propagation direction from the laser axis is around 10°, whereas for ions it is more than double (see **Fig. 71**).

**Fig. 73:** Correlation between the main propagation direction of the EUV emission and the electron expansion.
6.1.3 Droplet Fragments

The fragments of the droplets accelerated and shattered by the pressure of the plasma on the droplet surface were quantified in terms of number, dimension and occurrence. The experimental procedure was based on the exposure of Si samples to the plasma and is explained in details in Section 3.3. In the following Sections the experimental results of the distribution of the fragments versus the angular position with respect to the laser axis, and the dependence on the ambient pressure is presented. For all the samples the operating condition of the irradiation system was at an irradiance of $2 \times 10^{11} \text{ W/cm}^2$ and a repetition rate of 6 kHz. The samples were placed at 200 mm from the irradiation region and the fragments expanding in a solid angle of $1.4 \times 10^{-4} \text{ sr}$ were captured.

Angular Distribution

The load of the fragments at 45°, 90° and 120° angles is shown qualitatively in Fig. 74, Fig. 75 and Fig. 76, respectively. The exposure was not constant in order to have a detectable level of fragments on the sample and was taken into account in the quantification. Already from the pictures it is clear that the load of the fragments increases dramatically with the angle from the laser axis.

![Image](image.png)

Fig. 74: Picture of a sample exposed at 45° for 10 Mshots.
6.1 Experimental Results

Fig. 75: Picture of a sample exposed at 90° for 6 Mshots.

Fig. 76: Picture of a sample exposed at 120° for 1 Mshots.

As the effect of the fragments on collector optics is a localized loss of reflection of EUV in correspondence with the area occupied by the fragment, data are quantified as portion of the surface occupied during a number of shots. The portion of the surface occupied by the fragments \( k \) was calculated using the image post-processing introduced in Section 3.3. Data at different angles from the laser axis are combined accounting for the exposure at a different number of pulses and plotted in Fig. 77.
Fig. 77: Percentage of the part of the surface occupied by fragments in 1 Mshot versus the angle from the laser axis. The pressure in the argon ambient gas was of 0.12 mbar. The fragments expand mainly towards the direction opposite to the side of the droplet that is irradiated by the laser.

Fig. 77 shows that the fragments expand mainly towards the direction opposite to the side of the droplet irradiated by the laser. Indeed, the vertical scale is logarithmic, in order to display all the results. At 120° the surface is covered three orders of magnitude faster than at 45°.

From the images of the samples the diameter of every fragments was also extracted by the image post-processing. The combination of pixel size and magnification used in the microscope allowed to detect fragments larger than 0.6 \( \mu \text{m} \). The data are combined in Fig. 78 in a map were the occurrence of different fragment dimensions is shown at different angles. The contribution in the forward direction (towards the laser) is mainly from fragments with a dimension on the sample smaller than 20% of the droplet diameter. Instead, at larger angles the contribution of larger fragments becomes increasingly important.
6.1 Experimental Results

Fig. 78: Map of the occurrence of the fragment diameter depending on the angle from the laser axis. The pressure of the argon ambient gas was of 0.12 mbar. The occurrence of larger fragments increases with the angle from the laser axis.

Influence of the Ambient Pressure

Fig. 79: Percentage of the part of the surface occupied by fragments in 1 Mshot versus the pressure of the argon ambient gas. The angle from the laser axis was 120°. Increase of pressure leads to a larger portion of the surface occupied during the same number of laser shots.
Analogously to the angular distribution, the ambient pressure was parametrically studied from $2 \times 10^{-2}$ mbar to 0.6 mbar. The samples were exposed at 120° for 1 Mshot. The result of the exposure of the samples is shown in Fig. 79, where the portion of the surface occupied by the fragments is plotted against the ambient pressure. The trend is a rapid increase in the fragment formation and ejection from around 0.1 mbar up to a flattening of the trend at around 0.4 mbar.

The dimension of the fragments is shown in Fig. 80. However, at the different pressures no monotonic change in the occurrence of fragments with different dimensions is visible at pressure larger than 0.12 mbar. The sample exposed at $2 \times 10^{-2}$ mbar had instead smaller fragment dimensions.

**Fig. 80**: Map of the occurrence of the fragment diameter depending on the argon ambient pressure. The angle from the laser axis was 120°. No clear structure is present in the occurrence of the fragment dimensions. Only at the lowest pressure of $2 \times 10^{-2}$ mbar the fragment size was smaller.

**Diameter**

The droplet diameter was varied from 50 µm to 30 µm, and samples were exposed at 120° for comparison. The exposure lasted 166 s as for the other samples at 120°, equivalent to 1
6.2 Discussion on Effects of Offset on Matter Expansion

Mshot. The surface deposition rate for the case with 30 µm droplets was around 0.7 %/Mshot, whereas it was 3.4 %/Mshot for the larger droplet diameter.

6.2 Discussion on Effects of Offset on Matter Expansion

In the following Sections the LPP and its effects are considered as a function of the offset Δ/ between the focal spot of the laser and the droplet position. The offset is defined as the distance between the focal spot and the droplet in a direction perpendicular to the laser axis as shown in Fig. 81. Because of the offset, plasma expands in the vacuum chamber with a main propagation direction different than the laser axis, as observed from the measurements in the previous Sections.

![Fig. 81: Sketch of the main plasma propagation direction and of the offset between droplet targets and laser focal spot on the plane defined by laser axis and droplet center of mass. The droplet irradiation is not centered with the laser axis, and generates a three-dimensional plasma expansion, which is characterized by the main plasma propagation direction.](image)

6.2.1 Ablation of the Spherical Target with Offset

The calculation presented in Section 5.2.1 was extended with the addition of an offset between the droplet position and the laser spot. This allows studying the influence of the offset on the region where ablation starts, and on the time needed for the ablation process to occur. The measured laser profile in time and space was used in the computations.
Fig. 82 shows the time needed for ablation on the droplet surface seen from the laser axis at different offsets in \( x \) direction. The effect of moving the droplet targets out of focus is a progressive shift of the region where ablation begins. In the frame of reference of the droplet, the ablated region shifts towards the laser focal spot. With the position of ablation on the droplet surface, also the direction of the initial expansion of vapor and plasma changes as it mainly follows the normal to the surface.\(^{86,187}\) The time needed for the ablation process to occur also changes with the offset, going from the minimum of 2 ns in the aligned case (see Fig. 54) up to around 20 ns for an offset of 125 \( \mu m \) (1.56 \( \phi_{\text{laser}} \)). Therefore, the offset has two main effects: it generates plasma with a preferential direction different from the laser axis, and it delays the ablation process. The former causes asymmetry in the plasma formation, with the direction of expansion naturally pointed towards the focal spot, and a pressure force on the droplet target that is not aligned with the laser axis. The effect of the latter is a loss of the laser radiation during the longer period needed for ablation. The analysis presented in this Section is valid provided that the droplet target does not shatter during ablation. This could lead to change in the geometry of the target facing the laser pulse and to enlargement of the droplet. Experiments are not yet available for liquid droplets irradiated with a laser with \( D_{dr}/\phi_{\text{laser}} = 0.625 \). However, similar conditions show that the timescale for deformation of irradiated droplets at the same irradiance is of microseconds,\(^{105}\) which, compared with 45 ns that is the maximum time needed for ablation with offset, goes in the direction of sustaining the analysis performed here.
### 6.2 Discussion on Effects of Offset on Matter Expansion

**Fig. 82**: Time needed to start the ablation on the surface of the droplet at different offsets. The view is perpendicular from the laser axis, as seen from the focusing lens. The droplet diameter was 50 µm ($D_{dl}/\phi_{lw}=0.625$). The offset between droplet and laser was varied parametrically with $\Delta l = 25$ µm, 50 µm, 75 µm, 100 µm, and 125 µm in the figures (a), (b), (c), (d) and (e), respectively. In the droplet frame of reference, the region where the surface is ablated moves towards the laser spot with increasing offset. And, the ablation process takes longer to happen.
Fig. 83 shows the direction of the initial pressure force acting on the droplet $\beta$ (opposite of the vapor and plasma expansion) as a function of the offset together with the time needed for ablation. The pressure force decreases from $180^\circ$ for the aligned case up to around $115^\circ$ for the last possible ablation event at around 3.5 laser radii from the axis, which is the scenario where the droplet target need all the laser pulse for ablation. The curve is not linear and flattens with the offset, as the gradient of the Gaussian spatial profile decreases. The time needed for ablation increases with the offset, since longer is needed to deposit enough energy to reach the threshold of 2 J/cm$^2$.\textsuperscript{160} The trend is exponential, and up to around 2.5 focal spot radii the time needed for ablation is below 10 ns.

Both an angle $\beta$ and a main plasma propagation direction $\delta$ have been defined also if initially the two parameters (both positive per definition) are expected to be related ($\beta = 90^\circ + \delta$), because the laser energy deposited on the plasma during expansion from the droplet surface could alter the main propagation direction, as discussed in Section 6.2.2. The asymmetry in the start of the ablation on the droplet surface depends on the ratio between the droplet diameter $D_\phi$ and the laser focal spot $\phi_{las}$. Indeed, a large droplet with
an offset from the laser axis sees a larger variation of the laser spatial profile on its surface compared to a smaller one. At the limit of $D_{dr}/\phi_{las}$ going towards 0 (infinitesimally small droplet compared to the laser spot size, or infinitely large focal spot size compared to the laser), the droplet would not see any difference in the energy flux arriving on its surface, and the ablation would always happen head-on irrespective on the offset. This trend is shown in Fig. 84(a), where $\beta$ is plotted against the offset for different $D_{dr}/\phi_{las}$ ratio. The minimum $\beta$ to have ablation decreases with an increasing $D_{dr}/\phi_{las}$. And, the maximum offset at which ablation might take place increases with $D_{dr}/\phi_{las}$. For the laser focal spot size used in this work of 80 $\mu$m, the range of droplet diameters on the plots in Fig. 84 is from 10 $\mu$m to 70 $\mu$m.

![Graph showing the effect of offset on matter expansion](image)

**Fig. 84**: Effects of the offset on the initial direction of the pressure force (a), and on the amount of energy remaining in the laser beam after the start of the ablation process (b). The different colors refer to different ratio of droplet diameter and laser focal spot size $D_{dr}/\phi_{las}$.

In Fig. 84(b) the energy of the laser radiation, which is available for processes happening after ablation, is plotted. Smaller $D_{dr}/\phi_{las}$ are more sensitive to the amount of energy remaining after ablation. However, more than 95% of the laser energy remains available for all the studied $D_{dr}/\phi_{las}$ at $\Delta l < \phi_{las}$. The available energy of the laser after ablation is not expected to be completely used, since the plasma develops in time, occupies the region where the energy flows (focal spot), and reaches critical density.
The maximum offset to have ablation visualized in Fig. 84(b) gives also directly an estimation for the minimum distance between the droplets not to have ablation of more than one droplet at the same time, which is between $1.5 \phi_{\text{las}}$ and $2 \phi_{\text{las}}$ depending on $D_{dr}$.

### 6.2.2 Plasma Expansion with Offset

After ablation other processes, e.g. EUV generation, need a threshold of fluence to start. In head-on irradiation, the scenario is optimized to deliver in the shortest time the highest fluence in the core of the plasma, but with an offset $\Delta l > 0$ the source of mass (center of the density distribution) is not aligned with the source of energy (along the laser axis). With offset, the head-on condition is approached when the plasma expands up to the center of the spot, where the local laser irradiance is at its maximum. The time needed to expand up to the center of the laser spot and to reach the critical density depends on the expansion speed of the plasma and on the offset. Masnavi et al.\textsuperscript{160} find a maximum expansion velocity just after ionization produced by Nd:YAG irradiation between 10 km/s and 40 km/s. But, the average velocity after ionization is only around 3.5 km/s, which rapidly increases with the fluence up to around 25 km/s at the peak of the laser pulses used in this work.\textsuperscript{160} For a first estimation of the time needed for the plasma to travel up to the center of the spot size, the average value of 14.25 km/s is taken. 5.6 ns are needed for $\Delta l = 1 \phi_{\text{las}}$, which is short compared to the overall duration of the laser pulse of around 60 ns (see Fig. 9).

The absorption of the laser radiation at the maximum irradiance in time for the case with $\Delta l = 1 \phi_{\text{las}}$ is shown in Fig. 85. In the calculation, the same relations exposed in Section 5.2.2 were used with the addition of a term for the offset, in order to give a qualitative description of what happens in the irradiation with offset. The first main feature in laser absorption with an offset is that the absorption region (defined as the region where 99% of the laser energy is absorbed), marked by the blue lines, is deeper on the side of the offset. This is due to the lower density compared with along the axis, which lead to a smaller absorption of the Gaussian laser profile. The second main feature is that also with the offset $1 \phi_{\text{las}}$ of the plasma expands enough to reach the critical density at the maximum irradiance of the laser beam.
6.2 Discussion on Effects of Offset on Matter Expansion

**Fig. 85**: Region of absorption of the laser pulse at the maximum irradiance in time the laser beam is at an offset $\Delta l = 1 \phi_{\text{las}}$. (a) and (b) are at two different views. The droplet target is at $x=-1 \phi_{\text{las}}, y=0, z=0$. The green contours are isodensity lines, and identify the boundary where the ablated mass is contained. The blue lines correspond to the position where 99% of the laser radiation has been absorbed.

The scenario presented in Fig. 85 is dependent on the laser profile in time and space. Indeed, a laser pulse with a long duration (as the one used in this work) leads to a plasma generation and expansion, which is less sensitive to offset. For example a laser pulse with the same spatial profile and an overall duration of only 6 ns would lead to a faster ablation, but the time needed to arrive up to the center of the focal spot would remain approximately the same, therefore the largest part of the laser pulse would be lost compared to only around 5% with the 60 ns overall pulse duration. The condition to reach the center of the focal spot is:

\[
\frac{c_s}{\sqrt{2}} > \frac{\Delta l}{\tau_{\text{las}}},
\]

\[
\tau_{\text{las}} > \frac{\sqrt{2} \Delta l}{c_s},
\]  

(53)

where the factor $\sqrt{2}$ was added to take into account an average initial plasma propagation direction at an angle with respect to the laser axis. The longer the laser pulse, the larger the
offset can be for the similar amount of energy absorbed by the plasma. Eq.(53) describes the sensitivity of plasma formation and of laser energy absorption on the offset. Recalling the anisotropy condition of Eq. (19), another effect of a longer laser pulse with the same spatial profile is an increase of the anisotropy in the plasma expansion, leading to a density profile with a smaller \( b/a \).

An evidence of the fast plasma expansion versus the laser pulse duration in the experiments is the main propagation angle of ions and electrons, which is close to the laser axis. For ions the maximum was of around 20°, and for electrons around the half. Electrons are most probably even less sensitive to the offset then ions, because of their larger mobility due to their smaller mass.

The correlation between EUV and ion main propagation direction is expected from the model developed in Chapter 4. Indeed, the emission of EUV was found to depend on the absorption around the EDR region by the plasma, therefore a rotation of the plasma directly influences the amount of EUV that filter through.

### 6.2.3 Ejection of Fragments

When droplets are irradiated with \( \Delta l > 0 \), hence not perfectly head-on, the process of droplet shattering and acceleration becomes three-dimensional as opposed to axisymmetric. Under these circumstances, the fragments expand with a preferential direction different than the laser axis.

In order to compare samples exposed to similar conditions, a shutter was installed between the aperture and the sample. The opening of the shutter was controlled with the EUV emission from the plasma, only within a given range of EUV emission the shutter was turned on. EUV emission was measured and stored with the EUV energy monitor placed at 60° from the laser axis. **Fig. 86** shows the mean EUV and its standard deviation for the study on the angular distribution. For all the samples the mean generated EUV changed less than 10%. And, the standard deviation during exposure was of around 20% for most of the samples, with an exception at 60°, where the standard deviation was larger, of around 30%. The EUV standard deviation was not smaller in the experiments by purpose because the off-design operating condition was tested.
Discussion on Effects of Offset on Matter Expansion

**Fig. 86:** Mean EUV integrated level and its standard deviation during exposure of samples at different angle from the laser axis. The mean level of EUV changed less than 10%, and the standard deviation was of about 20%, with an exception at 60° with 30% standard deviation.

The fact that the most of the fragments are directed opposite to the irradiated surface (Fig. 77) is explained by the large pressure caused by laser irradiation, which acts on one side of the surface of the droplet. The large pressure gradient in space and time (as the plasma generates in few nanoseconds) is due to the vapor and plasma generated on the surface of the droplet, and has mainly two effects. On one hand, it induces the shattering of the droplet. On the other hand, the high-temperature and high-pressure fluid on the droplet surface transfers momentum to the free-falling droplet in the direction normal to the surface.

Considering again the angular distribution of the fragments (see Fig. 77), the large increase of fragments that was measured at 120° (43 times larger than at 90°) could be due to the pressure force acting on the droplets with offset. Indeed, the direction of the initial pressure force shown in Fig. 83 flattens at around 120° where the offset is estimated to be around 2 laser focal spot radii. A second observation from the measurements on the samples is that the amount of larger fragments in the distribution of the fragment diameter increases after 105° with the angle from the laser axis. Both the greater amount of fragments and the larger observed cluster size, are indicators that the angular distribution is dominated by the
Chapter 6 NON-IDEAL DROPLET IRRADIATION

acceleration of the droplet due to vapor and plasma on the surface in offset condition. And, the location where the pressure force acts is around the region where the ablation begins.

The smaller fragments (splash size < 10%$D_d$) that are more uniformly spread in space come instead mainly from spallation of the droplets and phase explosion. The droplet shattering was already observed in LPP in a number of works,\textsuperscript{103-105,188,189} but only imaging single events close to the IS.

In tin, the spallation is characterized by cavitation through the opening and growth of holes or bubbles in the fluid.\textsuperscript{190} The fragment size $s$, the time needed $t_s$, and the critical tension for fragmentation $P_s$ can be estimated using relations based on energy balance developed by Grady:\textsuperscript{190}

\[
\begin{align*}
  s &= \left( \frac{48\gamma_s}{c_0 D^2} \right)^{1/3}, \\
  t_s &= \frac{1}{c_0} \left( \frac{6\gamma_s}{\rho_0 D^2} \right)^{1/3}, \\
  P_s &= \left( 6\rho_0^2 c_0^2 \gamma_s D \right)^{1/3},
\end{align*}
\]

(54)

where $\gamma_s$ is the surface tension, $c_0$ the bulk sound speed, $\rho_0$ the reference mass density, and $D$ is the dilatation rate. For tin $\gamma_s = 0.5$ N/m, $c_0 = 2'480$ m/s, and $\rho_0 = 6'500$ kg/m\textsuperscript{3}. And, for shocks produced by lasers $D$ is in the order of $10^8$.\textsuperscript{191} Eqs. (54) lead to $s = 0.7$ µm, $t_s = 0.15$ ns, $P_s = 43$ MPa. The spallation should happen if the pressure acting on the surface is larger than the critical tension. In order to estimate the pressure from the plasma on the surface, the relation of Eliezer et al. was used,\textsuperscript{168} leading to a peak pressure of around 7 GPa, which is confirmed also by Yuspeh et al., who performed experiments on droplets at similar conditions.\textsuperscript{105} The fragmentation theory of Grady was used successfully for tin in a number of previous works.\textsuperscript{192,191} Therefore, spallation is expected to happen. The size of the small splashes at the different angles that were measured is in the order of magnitude of the estimation of the size of the fragments produced by spallation (also considering that $s$ is the diameter of the fragment during flight, which is smaller than the diameter of the splash measured on the sample). The fact that the small fragments are observed at all angles could be due to the geometry of the droplet and to an ejection of the fragments normal to the surface. Indeed, it was observed that fragments are ejected with a preferential direction.
normal to the surface, and was used in PLD system to deposit the thin film without fragments by having the substrate off-axis with the ablated surface.\textsuperscript{193,194}

Not only the pressure on the surface of the droplet has a role on processes happening, but also the local increase of temperature leads to disruptive changes in the structure of the droplet surface. Heating of the surface up to temperature close to the thermodynamic critical temperature $T_c$ (around 0.9 $T_a$) lead to homogenous bubble nucleation, and the liquid becomes a mixture of vapor and liquid homogeneously spread in the volume. The large pressure of the vapor transfers momentum to the material that remained liquid leading to explosive boiling or phase explosion.\textsuperscript{195} Indeed, experimental and numerical studies have reported the presence of explosive boiling at laser fluences smaller than the one considered in this work, at shorter pulse duration, and with target materials with higher melting temperature as tin;\textsuperscript{195-198} therefore the phenomena is expected also in the droplets irradiated in this work.

The increase of the amount of fragments with the ambient pressure means that the ambient gas participates to the dynamics of the plasma generation and droplet fragmentation. The ambient pressure is important already from the ablation process. The rate of evaporation from the surface of the droplets can be estimated with the Hertz-Knudsen equation. In the general form, the equation is:\textsuperscript{199}

\[
\frac{1}{A_e} \frac{dN}{dt} = a_v \left( p_{eq} - p_{amb} \right) \sqrt{\frac{N_A}{2\pi M k_B T_{sf}}} ,
\]

where $dN$ is the number of atoms evaporating from the surface area $A_e$ in a time $dt$; $a_v$ is a coefficient, $M$ is the molar mass, $N_A$ is the Avogadro’s number, $T_{sf}$ is the surface temperature. And, the pressure terms are the equilibrium pressure $p_{eq}$ and the ambient pressure $p_{amb}$. The equilibrium pressure is the pressure given from the maximum possible amount of atoms leaving the surface through evaporation, and is estimated by the Clausius-Clapeyron relation:

\[
p_{eq} = p_{amb} \exp \left( \frac{H_M}{k_B N_A} \left( \frac{1}{T_{eq}} - \frac{1}{T_{sf}} \right) \right) ,
\]
where $H_v$ is the enthalpy of vaporization and $T_{eq}$ is the equilibrium boiling temperature. Substituting Eq.(56) in Eq.(55) leads to the direct relation between the evaporation rate and the ambient pressure: $dN/dt \propto p_{amb}$. The evaporation rate is proportional to the ambient pressure, therefore the rate of evaporation on the surface of the free-falling droplet increases with the ambient pressure. The consequence is larger force acting on the droplet, and potentially a more disruptive phase explosion, which could be the cause of the larger amount of fragments with the increasing ambient pressure observed in Fig. 79.

The comparison with the different droplet diameters, which led to a smaller amount of fragment for the smaller droplet, is explained with the mass available for to the shattering. Indeed, the ratio of the two deposition rates is 4.9, which is approximately the ratio of the droplet volumes (or mass) that is $50^3/30^3 \approx 4.6$.

### 6.3 Summary

The non-ideal irradiation of droplets was experimentally investigated. The simultaneous distribution of ions and electrons was measured with an array of FC together with two EUV sensors. The fragments ejected from the free-falling droplet targets were studied with the exposure of Si samples to the plasma at different angles from the laser axis, at different ambient pressures and with two droplet diameters.

The ablation of the droplets begins in a location, which is a function of the laser spot size, the droplet diameter and the offset. As the offset increases, the ablation in the frame of reference of the droplet begins closer to the focal spot. Initially, plasma expands normal to the droplet surface, hence towards the laser focus, then as the laser duration is long, plasma expands enough to absorb part of the energy of the laser, and the expansion straightens in the direction of the laser axis, as it is where thermal energy increases. The condition to absorb part of the laser radiation, and reach or approach condition for EUV generation, is to have a laser pulse duration longer than $\sqrt{2} \Delta l / c_s$, where $\Delta l$ is the offset, and $c_s$ is the plasma mean sound speed. In the experiments, the condition was satisfied and led to a plasma expansion towards a direction close to the axis with maximum main propagation direction 20° off the laser axis. The ion main propagation direction and the EUV main
emission direction are correlated, which is due to the opacity to EUV of the plasma around the EDR region. Electrons are less sensitive to the offset than ions and EUV.

In order to avoid the ablation of more than one droplet per pulse the distance between the droplets should be larger than 1.5 $\phi_{\text{las}}$ to 2 $\phi_{\text{las}}$, depending on the droplet diameter.

On the droplet surface and in its interior a number of processes take place and lead to the ejection of fragments. A net acceleration is experienced by the droplet because of the expansion of vapor and plasma. The main direction of the force acting on the droplet is normal to the droplet surface in the ablation region and cause the rapid increase of large fragments ejection detected experimentally at angles larger than 105°. The angle is expected to depend on the ratio $D_{\text{dr}}/\phi_{\text{las}}$, the smaller the ratio the larger the angle of main fragment ejection is expected to be. The smaller fragments detected on the samples are expected to come from the spallation of the droplet and from phase explosion on the droplet surface.

The dimension of the fragments is in agreement with the spallation theory. From a debris perspective, the largest difference between ideal and non-ideal irradiation is the presence of the large fragments (splash diameter larger than 10%$D_{\text{dr}}$), which are accelerated at angles up to 105° from the laser axis as opposed to 180° for the aligned case. The ejection of fragments increases with the pressure, which is explained by the proportionality between the ambient pressure and the evaporation rate. And, a smaller droplet diameter leads to fewer fragments, proportionally to the droplet volume.
Chapter 7

MITIGATION

Chapters 5 and 6 show that in droplet-based LPP, matter in form of ions, neutrals and fragments is ejected together with the generation of EUV radiation. In radiation sources matter expansion is a byproduct of the wanted radiation and has to be managed, in order to avoid detrimental effects. In particular, in an EUV source the main detrimental effects are the deposition of the collector mirror and on the components of the stages upstream as illumination optics, reticle stage, and projection optics.

Therefore, matter expanding towards EUV collector mirrors and towards the upstream stages has to be stopped or deflected. Thus, a force has to be exerted on it. In classical physics a moving object can be accelerated in one of the following three ways: through inertial collisions, imposing an electric field, or imposing a magnetic field. The first way include fluid-dynamic forces, as the drag given by the motion of a flow, which microscopically is the result of a multitude of collisions between the flow and the particle or object in it. The second and third way, through an electric or magnetic field, implies that the particle or object to be accelerated is electrically charged. Therefore, atoms would have to be ionized, and fragments would need an imbalance of electrons in the metallic structure. The important parameter for the motion of a charged particle by an electric or magnetic field is the charge to mass ratio, as the force is proportional to the charge and the acceleration is inversely proportional to the mass. The charge to mass ratio becomes a limiting factor for the fragments. In order to deflect a fragment of the same distance as an
ion, the charge on the fragment surface would have to be so high that the fragment would undergo to fission. The process is known as Coulomb fission, and describes the detachment of small (sub-micron diameter) droplet from the main droplet.\textsuperscript{200-202} The process equilibrates the repulsive force of the charges on the surface of the droplet with the surface tension force, which in the case of Coulomb fission is not enough to contain the droplet. Therefore, fundamentally a fragment with a diameter of micrometers cannot follow a trajectory similar to the one of the ions. Regarding acceleration through an electric field, it should be noticed that the motion of charged particles compensates the imposed electric potential within the Debye length. Therefore, in order to affect the charged particles the distance between the electrodes should be shorter than the Debye length. Furthermore, the presence of plasma in the space between the electrodes could lead to unwanted discharges between the electrodes.

\begin{equation}
\Delta s_\perp = \frac{1}{2}\left(\frac{\Delta s_\parallel}{v_\parallel}\right)^2 \left[ a_{in} + a_{el} + a_{mag} \right] = \frac{1}{2}\left(\frac{\Delta s_\parallel}{v_\parallel}\right)^2 \left[ a_{in} + \frac{q}{m} (\Delta E + v_\parallel B) \right], \tag{57}
\end{equation}

Eq.(57) integrates the possible mechanisms to deflect of a distance $\Delta s_\perp$ a particle travelling with velocity $v_\parallel$, which gets accelerated tangentially to the initial velocity along a distance $\Delta s_\parallel$ of $a_{in}$, $a_{el}$, $a_{mag}$ due to inertial collisions, electric field $\Delta E$, and magnetic field $B$, respectively. Electric and magnetic fields contribute to the deflection only if the charge of the particle to deflect $q$ is larger than 0, therefore would be applicable only against ions. For example a Sn\textsuperscript{+} ion with a kinetic energy of 500 eV would be deflected of 10 mm along a distance of 50 mm by $\Delta E = 4 \times 10^4 \text{ V/m}$ (400 V over 1 cm distance between the electrodes) or by $B = 2\text{T}$. Deflection through magnetic field would not be subject to electrostatic shielding, but it could be challenging to generate the needed magnetic field for continuous operation.

In this thesis the mitigation strategy that was chosen is based on inertial collision, since this affects all kind of debris – ions, neutrals and fragments – without the need of re-ionize tin atoms, and charge the fragments. The disadvantage of this strategy is that the mitigation efficiency depends directly on the number of collisions during the flight of the debris particle, and the latter depends on the ambient pressure. However, the ambient pressure
cannot be increased as needed from the mitigation system alone as it affects directly the EUV transmission. Because of this trade-off a device has to be designed, which locally deflects debris with a layer of particles that have the highest possible momentum, hence with the largest possible velocity for a given mass flow. The device has to accelerate the shielding particles up to the pressure in the vacuum chamber, in rarefied condition. And, the layer of shielding particles has to be confined in space to avoid negative effects on the free-fall of the droplets from the droplet dispenser to the IS.

In the next Sections, design and testing of such a device are presented, followed by the interpretation of the experimental and computational results to derive design guidelines for similar devices at rarefied operating conditions. And, finally results of the use of the device in the radiation source are exposed.

7.1 Aerospike Performance

In order to accelerate a flow to supersonic velocity, diverging surfaces forming a nozzle are normally employed. Nozzles that operate close to or at rarefied condition ($Kn > 0.01$) are becoming increasingly important for a number of applications, including as propulsion systems for microspacecraft,$^{111,112,203}$ for refueling high temperature fusion plasmas,$^{131,204}$ and as protection mechanisms against third particles.$^{49,60}$ The goal of these systems is to generate a jet with the maximum possible momentum in a desired direction. In the first case, the optimization of the momentum maximizes the thrust directly, in the second it maximizes the fueling efficiency, and in the third it maximizes the debris mitigation efficiency and the confinement of the jet.

Among the different nozzle designs that are considered, linear aerospikes are investigated because of two main advantages over conventional bell shaped geometries: that the flow self-adjusts to the background pressure,$^{205,206}$ and that the manufacturing is easier and cheaper compared to conventional bell-shaped nozzles, in particular for micronozzles.$^{112}$ The challenge is that the flow field is significantly affected by the boundary effects along the nozzle walls.$^{207,209}$ Here, linear aerospikes are considered for the study of the kinetic boundary layer (KBL) or Knudsen layer, because they provide the ideal scenario for a
diverging two-dimensional surface designed to accelerate the flow by bounding it on only one side.

In the following Sections, the expansion of the flow in the transition regime over a diverging surface is addressed both experimentally and computationally. The phenomenon, which has been identified in this work to affect mainly the flow field at the transition regime, is the formation of a KBL. A KBL forms immediately, attached to a surface in a region where the near equilibrium and the continuum assumptions are not enforced.\textsuperscript{210,211} In principle, a KBL forms close to every surface, but it is negligible in most scenarios. The KBL becomes important where the collisions close to the surface are not enough to equilibrate thermodynamically the flow. In this situation, the gas-surface interaction directly influences the KBL, and has an increasingly important impact on the whole flow field. The experimental and computational results are used to build a phenomenological model for the growth of the KBL along the surface of the aerospike. Furthermore, a method of modifying the shape of the diverging surface to compensate for the non-ideal expansion of the rarefied flow is presented and experimentally evaluated.

The structure of this Section is as follows: first the design of the aerospike is presented, followed by the computational setup. Then, the experimental results and the validation of the simulations are exposed. Successively, the growth of the KBL is discussed considering local thermodynamic equilibrium close to the surface and gas-surface interaction models; and a phenomenological model is presented. Finally, the influence of the geometrical correction of the diverging surface on the global flow field is discussed, and considerations on source design are exposed.

\subsection*{7.1.1 Design}

The aerospike system, sketched in Fig. 87, was designed in a modular way that enabled an easy exchange of the diverging surfaces, which are referred to as spikes. Four different spikes with different geometries and different roughness averages were tested. The geometry of the spike for the base case was designed following the method of Angelino,\textsuperscript{212} which is based on the continuity equation and on the Prandtl-Meyer function for flow expansion. The design Mach number was 4.5, and the contour of the spike was cut at 65\% of the ideal length,
where the distance projected in vertical direction (z-axis) from the end of the contour to the ideal end is half of the throat height (TH). In order to minimize the three-dimensional flow structures, the spikes were designed with a width of 30 TH. Furthermore, walls were added at the ends of the spike contour in order to constrain the expansion of the flow in the third dimension. A heater and a heat exchanger were implemented close to the minimum cross section to heat the flow up to a stagnation temperature of 500 °C. The electrical heater had a power of 800 W, and the heat exchanger was a bed filled with stainless steel spheres with a diameter of 2 mm. The temperature of the flow was monitored by two thermocouples: one placed by the convergent section and one by the Ar inflow. Furthermore, the temperature was estimated and cross-checked against the stagnation pressure. Since in the temperature study presented in Section 7.1.3 the massflow rate and the gas element were kept constant, $p_0/\sqrt{T_0}$ is an invariant, and from the change in stagnation pressure the change in stagnation temperature can be estimated. The material used for all of the components of the aerospike was stainless steel.

Fig. 87: Sketch of the aerospike system. The design is modular and allows the exchange of the spike contour. The stagnation pressure is measured in the section immediately preceding the throat of the aerospike. Argon is injected in the system and passes through a heat exchanger surrounded by the heater. Two temperature sensors are integrated to estimate the stagnation temperature of the flow.
7.1 Aerospike Performance

7.1.2 Computational Setup

The direct simulation Monte Carlo (DSMC) method is the most widely used method of simulating a flow field in rarefied conditions. The Navier-Stokes equations that are solved in computational fluid dynamics (CFD) are not used in rarefied conditions, because they begin to break down.\textsuperscript{213} In a rarefied flow the assumptions that lead from the Boltzmann equation to the Navier-Stokes equations are no longer valid, namely continuum, and near equilibrium. Instead, the DSMC method simulates the motion and collisions of a large number of particles, reproducing the underlying physics of the Boltzmann equation. The computational code used in this work is based on the DS2V program of G.A. Bird.\textsuperscript{185} For collisions between particles, the variable hard sphere model (VHS) and the no time counter (NTC) collision technique are employed.\textsuperscript{214} Atomic argon was simulated with a reference diameter of $4.17 \times 10^{-10}$ m and a viscosity-temperature power law coefficient of 0.81. For collisions between particles and surface, two models are considered: the CLL model and the Maxwell model,\textsuperscript{185} which are the most widely used models of gas-surface interaction. The code separates collision cells and sampling cells. The former have a finer resolution, and their physical dimension was continuously adapted to the flow to optimize the accuracy of the simulation.\textsuperscript{214} In a standard simulation performed in this work, around 600,000 collision cells were used. Furthermore, the physical dimension and distribution of the sampling cells were also continuously adapted to the flow, in order to resolve the local gradients in the flow structure.

**Fig. 88:** Part of the computational domain close to the diverging surface. The blue dotted area is the grid for flow quantities sampling, the green line is the solid diverging surface, and the red line is the flow inlet, which corresponds to the throat of the aerospike.
All of the simulations are two-dimensional, because linear aerospike are considered. The dimension of the computational domain was 90 TH in length and 24 TH in height. The simulations are unsteady and start with only the presence of argon gas at the pressure of the vacuum chamber for the given case. At the margins of the computational domain the boundary condition was a constant pressure, equal to the initial pressure in the domain. The part of the computational domain around the spike is shown in Fig. 88. The argon computational particles inlet corresponded to the position of the throat of the aerospike (in two dimensions it is the shortest line between the spike and the corner that initiate the Prandtl-Meyer expansion) and is depicted by the red line in Fig. 88. Sonic speed $a^*$ at the throat was imposed on the injected particles, with the superposition of the thermal speed given by the temperature at the throat $T^*$. The number of injected particles per iteration was calculated with the number density at the throat $n^*$. $T^*$, $n^*$ and $a^*$ were calculated with isentropic relations dependent on the stagnation values $p_0$ and $T_0$, which were the input for every simulations.

7.1.3 Measurements

Far Field Measurements

A parametric study was performed with the experimental setup described in Section 3.4. The ambient pressure $p_{amb}$ was changed from $7 \times 10^{-2}$ mbar to 1.5 mbar, and the pressure ratio across the aerospike between $p_0$ and $p_{amb}$ was kept constant. Using the classical definition of the Knudsen number to describe the level of rarefaction of the flow, with the height of the throat of the aerospike as the characteristic dimension, results in values of $Kn = 6.4$ to $Kn = 0.42$, respectively. Therefore, according to definitions the flow is in transition regime. This provides a rough picture of the experimental conditions, but a more precise calculation of non-equilibrium parameters based on local changes of the flow quantities is given in Section 7.1.5. The measurements in Fig. 89 show a qualitative comparison between the horizontal velocity of two cases at the same pressure ratio: with an ambient pressure of 0.11 mbar and a stagnation pressure of 18.4 mbar (Fig. 89(a)), and 0.85 mbar
7.1 Aerospike Performance

ambient pressure with 145.3 mbar stagnation pressure (Fig. 89(b)). All the dimensions in the plots are normalized with the TH. The two contour plots show different behaviours: when the ambient pressure is lower (0.11 mbar), the jet lifts with an angle of around 14.4°, whilst the higher ambient pressure (0.85 mbar) shows an angle of around 4°. In the latter case, with the higher ambient pressure, the jet expansion is closer to the ideal, which would be expected when considering only the method of Angelino.

**Fig. 89:** Measured map of the horizontal velocity in m/s in the space non-dimensionalized with TH. (a) at 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure. (b) at 0.85 mbar ambient pressure and 145.3 mbar stagnation pressure. Both cases are at the same pressure ratio. The jet in the lower pressure case lifts up with an angle of approximately 14°, whilst for the larger ambient pressure the angle is around 4°.

A more quantitative representation of the experimental results is shown in Fig. 90. The quantity chosen to compare the different flow structures is the KBL thickness, and it is calculated as the height of the flow from the surface when it reaches 95% of the local maximum velocity. The latter is calculated by considering the velocity profile along
directions perpendicular to the diverging surface. For every position on the spike surface, the local maximum velocity and the position of 95% of local maximum velocity along the perpendicular is evaluated. This quantity is represented by $t_{KBL}$ and is non-dimensionalized by the TH of the aerospike. In Fig. 90, the KBL thickness is plotted against the ambient pressure, and the different lines correspond to the KBL thickness at the exit, at 12.5 TH, and at 36 TH from the exit. The KBL thickness increases continuously from the exit of the aerospike to 36 TH from the exit. Furthermore, considering the furthest position from the throat, from the point of $7 \times 10^{-2}$ mbar, a 14-fold increase in the ambient pressure leads to a 3-fold decrease in the KBL thickness. The dependency between $t_{KBL}$ and $p_{amb}$ is discussed in Section 7.1.5.

![Fig. 90: Measured KBL thickness at different ambient pressures. The pressure ratio for all measurements is kept constant. The different lines show the KBL thickness at different distances from the exit of the aerospike. The KBL thickness is non-dimensionalized with TH. Considering the furthest position, from $7 \times 10^{-2}$ mbar a 14-fold increase in the ambient pressure leads to a 3-fold decrease in the KBL thickness.]

Near Field Measurements

Fig. 89 shows that the lifting of the jet is approximately linear downstream of the aerospike exit, therefore it is expected that the flow angle is defined by the flow expansion on the surface of the spike.
Two maps of the horizontal velocity on the spike surface are shown in Fig. 91. The maps show the cases with 0.11 mbar and 0.85 mbar ambient pressure as in Fig. 89. The spike surface is shown as the thick black curve, and the position of the corner that initiates the Prandtl-Meyer expansion is represented by the red dot. The light blue dots that are also included in the visuals represent the KBL boundary determined by the experimental results, and the dashed black line shows a linear growth rate superimposed on the spike surface in the perpendicular direction as described by Eq.(58).

![Fig. 91: Measured horizontal velocity map in m/s in the space non-dimensionalized with TH. (a) at 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure. (b) at 0.85 mbar ambient pressure and 145.3 mbar stagnation pressure. Both cases are at the same pressure ratio. The red dot represents the position of the corner that initiates the Prandtl-Meyer expansion. The thick black line is the spike surface. The light blue dots are the KBL boundary, and the dashed black line is a linear growth rate superimposed on the spike surface in the perpendicular direction. For both extremes, the KBL growth is linear versus the spike length.](image-url)
Fig. 91 shows that in both cases a linear growth rate predicts the thickness of the KBL from the experiments. This demonstrates that $t_{\text{KBL}}$ is linear versus the length of the spike, and the following relation is valid:

$$t_{\text{KBL}}(l) = c_{\text{GR}} l + t_0,$$  \hspace{1cm} (58)$$

where $l$ is the distance along the surface of the spike from the throat, $t_0$ is the KBL thickness at the throat, and $c_{\text{GR}}$ is the non-dimensional growth rate. In the first case with the ambient pressure equal to $p_{\text{amb}} = 0.11$ mbar the non-dimensional growth rate was $c_{\text{GR}} = 0.18$, whilst with $p_{\text{amb}} = 0.85$ mbar it was $c_{\text{GR}} = 0.1$.

### Homogeneity in the Third Dimension

**Fig. 92**: Measured map of the measured horizontal velocity in m/s in the y-z plane non-dimensionalized with TH. The horizontal position was 30 TH downstream of the throat. The operating condition was at 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure. The horizontal velocity is homogeneous in the width of the aerospike.
The three-dimensionality of the problem was considered by measuring not only in the x-z plane, but also in the y-z plane. Fig. 92 shows a map of the horizontal velocity in the latter plane, at a horizontal (along the x-axis) position 30 TH downstream of the throat. The operating condition was the same as in Fig. 89(a) and Fig. 91(a). Also in Fig. 92, the lift up of the jet in vertical direction is observed, and it is homogeneous in the spanwise dimension, along the width of the aerospike and over the whole measured domain. This justifies the assumption made for the analysis of the results in the x-z plane that the flow field is quasi-two-dimensional.

**Effect of the Roughness on the Jet Expansion**

The gas-surface interaction plays a major role in the generation of the KBL (Section 7.1.5). The surface roughness of the spike was changed, in order to influence the gas-surface interaction. The surface of the base case that was presented in the previous sections was milled with a CNC and had a roughness average of 0.49 μm. Other spikes were built with the same geometry, but with different techniques: electrical discharge machining (EDM) and rapid prototyping (RP). The EDM spike was then polished and a roughness average of 0.023 μm was obtained, whilst the spike manufactured with RP was not further processed and a roughness average of 14.157 μm was measured. The values of the roughness average were measured with a surface roughness measuring tester from Mitutoyo (model SJ-210). The flow field from the different profiles was then measured using the same experimental procedure as in the base case.

Fig. 93 shows the quantitative comparison of the velocity profile at the exit of the aerospike for the three different spikes. The difference is of only around 10 m/s, around 2% of the velocity value. The change in the roughness of the spike from 0.023 μm to 14.157 μm did not produce any significant effect (i.e. none larger than the measurement error) on the KBL growth. Therefore, the gas-surface interaction in the considered rarefied domain is not dependent on the surface roughness.
Fig. 93: Measured profile of the horizontal velocity at the exit of the aerospike: the height is non-dimensionalized with TH. The different points refer to measurements with different average roughness of the spike surface. The latter was 0.023 μm for the polished surface, 0.49 μm for the milled one, and 14.157 μm when using rapid prototyping. The difference lies in the error bars close to the surface, and at a height larger than 8 the polished case shows a slightly larger horizontal velocity, up to 2%.

**Correction of the Profile of the Aerospike**

To achieve a flow field at the exit of the aerospike closer to the ideal case with the expansion happening mainly following the x-direction, the geometry of the spike was modified. Since the KBL growth is linear versus the length of the spike (previous Section), the correction follows a linear growth and is subtracted in a direction normal to the surface. The correction is described by the following equations:

\[
\begin{align*}
    x_{LC} &= x_{\text{ideal}}(l) + n_z(l)c_{LC}l, \\
    z_{LC} &= z_{\text{ideal}}(l) + n_x(l)c_{LC}l,
\end{align*}
\]

where \(x_{\text{ideal}}, z_{\text{ideal}}, x_{LC}, z_{LC}\) are the coordinates of the profile of the ideal spike designed using the method of Angelino and of the linearly corrected one, respectively; \(c_{LC}\) is the degree of freedom of the correction; and \(n_z, n_x\) are the components of the vector normal to the surface of the ideal spike at a distance \(l\) from the throat. The results of the correction together with the ideal profile for a \(c_{LC}\) equal to 0.12 are shown in Fig. 94.
Fig. 94: The profiles of the spikes before and after the correction. The axes are non-dimensionalized with TH. The ideal profile, which has been designed using the method of characteristics, is represented by the solid blue line. The profile corrected with a linear function of the spike length is represented by the dashed red line.

Fig. 95: Measured horizontal velocity map in the space non-dimensionalized with TH. The operating condition is at 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure. The experiment was performed with the corrected shape of the spike, which is shown as the thick black line. For reference, the shape of the ideal contour is also plotted, with a dashed black line. The flow leaves the aerospike with a direction nearly parallel to the x-axis.

The flow field of the corrected spike geometry is illustrated in Fig. 95 at the operating condition with 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure. Together
with the map, the corrected spike surface (thick black line) and the ideal one (dashed black line) are also represented for reference. By comparing the map with Fig. 91(a), which was measured at the same operating condition, it is evident that the lift up of the jet versus the x-axis almost disappears. In this way the jet leaves the aerospike with a direction that is close to parallel to the x-axis.

In order to quantify the improvement in the direction of expansion of the jet, Fig. 96 shows the measurements of the velocity profiles at the exit of the aerospike and at 30 TH from the exit for both the corrected and uncorrected shapes. At the exit of the aerospike the correction has the effect of lowering the jet by around 3.5 TH, and at 30 TH from the exit by around 10 TH. Considering the angle between the mean flow direction and the x-axis, the correction straightened the jet from 14.4° to 4.8°.

**Fig. 96:** Measured profiles of the horizontal velocity, the height is non-dimensionalized with TH. The operating condition is equal for all the profiles to 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure. The solid blue profiles show the horizontal velocity for the ideal shape of the spike, whilst the dashed red lines show to the corrected one. The dots are measurements at the exit of the aerospike, and the triangles are measurements at 30 TH from the exit. The correction straightened the jet from 14.4° to 4.8°.
Temperature

Experiments to study the effects of stagnation temperature on a constant massflow were performed using the heater that was built in the aerospike system. For the tests illustrated in Fig. 97, the corrected contour of the spike was used. The figure shows three different profiles at the exit of the aerospike for stagnation temperatures of 294 K, 492 K, 703 K. As expected, the velocity increases and the profile becomes wider with the increase in the temperature. The width of the jet is defined by the portion of the flow that has a velocity equal to the ± 5% of the maximum. The difference in width between the standard case at 294 K and at 703 K is of a factor of 1.6, from 4 TH to 6.5 TH. The increase is due to an increase in the thermal velocity of the core of the jet, which has a random direction and is superimposed on the drift velocity.

![Graph showing measured velocity profiles at different temperatures](image)

Fig. 97: Measured velocity profiles at different temperatures. The profile at ambient pressure was 0.11 mbar. The massflow for the three cases was constant. The height is non-dimensionalized with the TH. The velocity profile is broader at larger temperatures, and the velocity of the flow increases.

7.1.4 Validation of the Numerical Model

In order to investigate the problem with finer details, in particular close to the surface of the diverging contour, DSMC simulations were performed.
Since the gas-surface interaction has a dominant influence on the KBL growth, the choice of the parameters of the CLL model was considered in the validation. The parameters of the CLL model are the accommodation coefficient for the normal kinetic energy $\alpha_n$ and the accommodation coefficient for the tangential momentum $\sigma_t$.\(^{215}\) The parameters are defined as:\(^{216,217}\)

$$
\alpha_n = \frac{e_i - e_r}{e_i - e_w}, \quad \sigma_t = \frac{\tau_i - \tau_r}{\tau_i},
$$

where $e$ refers to the average kinetic energy for the normal component of the velocity and $\tau$ to the momentum acting tangential to the surface; and the subscript $i, r, w$ mean incident component, reflected component, and the component that would be produced by a pure diffuse reflection at the temperature of the surface, respectively.

![Profile of the horizontal velocity versus the height for computations with different values of $\sigma_t$ at a constant $\alpha_n = 1$. The height is non-dimensionalized with TH. The closest value to experiments is with $\sigma_t = 0.875$.](image)

**Fig. 98:** Profile of the horizontal velocity versus the height for computations with different values of $\sigma_t$ at a constant $\alpha_n = 1$. The height is non-dimensionalized with TH. The closest value to experiments is with $\sigma_t = 0.875$.

The starting point of the parametric study was with full accommodation of both the normal kinetic energy and the tangential momentum coefficients $\alpha_n = 1, \sigma_t = 1$. First the accommodation coefficient for the tangential momentum was varied from 0.5 to 1. The
results are shown in Fig. 98, together with the experimental results at the end of the spike. The value of $\sigma_t = 0.875$ gives the closest profile to the experiments for vertical positions both below and above the location of the velocity maximum. The value of $\sigma_t$ is in accordance with Padilla et al.,\textsuperscript{218} who found the best agreement with $\sigma_t = 0.9$.

With the best value for $\sigma_t$, another set of simulations was performed by changing the accommodation coefficient for the kinetic energy from 0 to 1. This parametric study did not produce a significant change in the velocity profile, in particular close to the surface as shown in Fig. 99. For this reason, the value of the accommodation coefficient for the kinetic energy was left equal to $\alpha_n = 1$.

![Fig. 99: Profile of the horizontal velocity versus the height for computations with different values of $\alpha_n$ at a constant $\sigma_t = 0.875$. The height is non-dimensionalized with TH. The profile does not change significantly by varying $\alpha_n$, especially further from the surface.](image)

Comparing the results of the simulation using $\sigma_t = 0.875$ and $\alpha_n = 1$ with the measurements, the underprediction of the maximum horizontal velocity is 2%, and the underprediction of the thickness of the KBL is 5%.
7.1.5 *Kinetic Boundary Layer Formation*

The method of Angelino and the physical models based on the Euler equations predict that the behavior of the expanding flow in a nozzle is dominated by the pressure ratio between the stagnation pressure and the ambient pressure. These models are limited, because they assume that the flow is inviscid, in continuum, and in near equilibrium. The latter assumptions are not valid in reality under the operating conditions of this work, as shown by the difference between Fig. 89(a) and Fig. 89(b). The non-ideal behavior of the flow is due to the formation of the KBL, whose thickness increases with a decrease in the ambient pressure (shown in Fig. 90). The KBL is different from the viscous boundary layer, which is predicted by the Navier-Stokes equations, because in the viscous boundary layer the flow is in near equilibrium and in continuum, whilst in the KBL these assumptions are not valid. For this reason, reductions of Boltzmann equations such as the Navier-Stokes (or the Euler) equations are not suitable models in this work. The treatment has to be fully rarefied, as with the DSMC simulations, or the solution of special cases of the Boltzmann equations like the Kramers problem. In both approaches, the collisional operator and the gas-surface interaction model are crucial for the correctness of the results.

In order to classify the rarefaction of the flow and to confirm that the lifting of the jet observed in the experiments is due to the KBL, the non-equilibrium parameters are calculated, the influence of the gas-surface interaction on the KBL growth is shown, and a phenomenological model is proposed to describe the growth of the KBL.

**Non-equilibrium close to the Surface**

With the results of the simulation, it is possible to define with high resolution the non-equilibrium parameters. Two definitions were considered: the Bird parameter \( P \),\(^{219} \) and the gradient-length local Knudsen number \( Kn_{GLL} \) from Boyd *et al.*\(^ {148} \). \( Kn_{GLL} \) was defined in Eq.(20) and \( P \) is defined as follows:

\[
P = \frac{1}{\nu} \left| \frac{D (\ln \rho)}{Dt} \right|,
\]

\(^{(61)}\)
where $\nu$ is the collision frequency, and $\rho$ is the flow density. When the density is chosen as the flow property for $Kn_{GLL}$ and the flow is steady, the two parameters are directly related. As illustrated in Eq.(62), $P$ is then proportional to $Kn_{GLL,\rho}$ and to the Mach number $M$.

$$P = M \sqrt[8]{\frac{\pi \lambda}{8} K_{GLL,\rho}}.$$  \hspace{1cm} (62)

For both parameters a threshold is defined, above which the flow field is considered in non-equilibrium. From the definitions, the meaning of the non-equilibrium parameters is that the change of the flow quantity (for example the density) in space is larger than the maximum that the flow can handle to restore equilibrium through collisions. The estimation of the break down of the approaches that assume continuum and near equilibrium depends on the numerical value that is set as threshold. Considering $Kn_{GLL,\rho}$, values larger than 0.05 indicate rarefaction and non-equilibrium, whereas for $P$ the threshold value is 0.02.

![Graph](image)

**Fig. 100**: Non-equilibrium parameters on the surface of the spike with 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure. The distance from the throat is non-dimensionalized with TH. $Kn_{GLL,\rho}$ is represented by the continuous blue line, and $P$ by the dashed green line. Both parameters predict non-equilibrium flow over the whole spike surface, and both show the same trend.

The values of $P$ and $Kn_{GLL,\rho}$ on the diverging surface are plotted in **Fig. 100** for the base case with 0.11 mbar of ambient pressure and 18.4 mbar of supplied stagnation pressure.
Both $\textit{Kn}_{\text{GLP}}$ and $P$ start and remain above the threshold of non-equilibrium, and the values at the throat are $\textit{Kn}_{\text{GLP}} = 0.1$ and $P = 0.035$, which are approximately twice the threshold for both the definitions. The values increase up to $\textit{Kn}_{\text{GLP}} = 6.4$ and $P = 1.28$ at the exit of the aerospike. Both definitions of non-equilibrium predict a similar profile, meaning that the Mach number at the surface is not a strong function of the position along the spike.

**Fig. 101**: (a) $\textit{Kn}_{\text{GLP}}$ with 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure. (b) $\textit{Kn}_{\text{GLP}}$ with 0.85 mbar ambient pressure and 145.3 mbar stagnation pressure. The axes are non-dimensionalized with the TH. The dashed curves represent the KBL from experimental results. The location of equilibrium as predicted by the parameters is approximately the core of the flow, which lifts up from the surface. Close to the surface the flow is in non-equilibrium.

Since the trend between the two parameters is similar, in **Fig. 101** only $\textit{Kn}_{\text{GLP}}$ is considered. **Fig. 101** compares qualitatively the two extreme cases, i.e. with an ambient pressure of 0.11 mbar and 0.85 mbar and supplied stagnation pressure of 18.4 mbar and 145.3 mbar, respectively. The color scale in both **Fig. 101(a)** and **Fig. 101(b)** is the same, and is chosen such that the dark blue color means that there are enough collisions to reach an equilibrium Maxwellian state despite the density gradient ($\textit{Kn}_{\text{GLP}} \leq 0.05$), and any other color means that the collisions are not sufficient ($\textit{Kn}_{\text{GLP}} > 0.05$). Together with the
map of $Kn_{GLL}$, the contour of the spike is drawn as a solid black line, and the linear growth approximating the KBL from the experimental results is represented by the dashed line (as in Fig. 91). The plots confirm that the region, which has a deficit in velocity is in non-equilibrium and corresponds to what was identified as KBL. The non-equilibrium condition of the flow is more severe downstream and towards the surface of the aerospike. Also in Fig. 101(b), with a stagnation pressure of 145.3 mbar, the flow within the KBL is mostly in non-equilibrium, and the equilibrium region around the core flow is wider. This shows that the trend towards larger pressures decreases the KBL and the non-equilibrium region, as expected from the measurements.

**Gas-Surface Interaction in the Kinetic Boundary Layer**

The computations were used to study the influence of the gas-surface interaction on the dynamics of the lifting of the jet from the surface. To show that the lifting of the jet observed in Fig. 89(a) results from the diffusive nature of the interaction between jet and surface, two simulations are compared: one with the validated setup, and one where a fully specular boundary condition for the spike was implemented using the Maxwell model.

![Fig. 102: Computational results for the horizontal velocity with an ambient pressure of 0.11 mbar and a stagnation pressure of 18.4 mbar, the axes are non-dimensionalized with the TH.](image)
(a) The spike surface is modeled as the validated case. (b) The spike surface is modeled with specular reflections. In the fully specular case the flow follows the surface with no lifting, as observed at larger ambient pressures.

The comparison for $p_0 = 18.4$ mbar and $p_{amb} = 0.11$ mbar is depicted in Fig. 102. In the fully specular case (Fig. 102(b)), the jet follows the surface with no lifting; instead, the lifting is observed in the result validated in Section 7.1.4 (Fig. 102(a)). The lifting of the jet occurs because the particles that impinge onto the surface are reflected in an almost random direction (the accommodation coefficients are close to 1), and undergo collisions with other particles approaching the jet, transferring momentum in a vertical direction and creating on the surface a layer of particles with randomized motion. The motion of particles in this layer is shown in Fig. 103(a), where the trajectories of the particles that hit the surface during one time step are plotted. The particles need the thickness of the KBL in order to move again with a direction close to that of the main stream. By way of comparison, Fig. 103(b) shows the trajectory of the particles that hit the surface at the same physical time as in Fig. 103(a), but for the specular case. In the case of the specular boundary condition, the layer of randomized motion is shorter. This is due to the direction that the particles take after reflection, which is nearly parallel to the surface. Therefore fewer collisions are needed to redirect the reflected particles with a direction close to parallel to the surface. A larger ambient pressure instead has the effect that more particles are present in the main stream, therefore more collisions happens in the same space, and more momentum is transferred per unit length on the reflected particles. In both the cases, the KBL thickness is shorter than the real case at 0.11 mbar of ambient pressure.

A similar methodology was applied by Santos, who distinguished molecules of the freestream, molecules that collided with the surface, and molecules indirectly affected by the body. As in this work, he found a smaller $\sigma$, (towards specular reflection) led to less molecules affected by the body, which is due to the less diffusive nature of the particle direction once they leave the surface.
Fig. 103: Trajectory of selected particles after reflection off the surface with an ambient pressure of 0.11 mbar, the axes are non-dimensionalized with the TH. (a) With validated CLL model for the gas-surface interaction. (b) With specular boundary condition for the gas-surface interaction.

From the plots illustrated in Fig. 102 and Fig. 103 it is clear that the gas-surface interaction plays an important role in the formation of the KBL. In reality the gas-surface interaction can be modified either physically, by changing the topography of the surface, or chemically by changing the constituents of the surface and their structure. In this work the first type of modification has been investigated in Section 7.1.3, and it has been shown that modifying the surface from an average surface roughness of 0.023 μm to 14.157 μm does not produce any significant change in the flow lifting. This means that the gas-surface interaction is dominated by the adsorption of the gas atom or molecule from the surface, rather than by the topography in tens of microns scale. The adsorption is dependent on the material at the surface,\textsuperscript{220} therefore a change in the material of the aerospike might lead to a change in the gas-surface interaction, and indirectly to other accommodation factors for the CLL model.
Phenomenological Model for Kinetic Boundary Layer Growth

The obtained experimental and numerical results have shown that the KBL grows monotonically with decreasing pressure (Fig. 90) and grows linearly with the distance along the surface under a steady operating condition (Fig. 91).

![Graph](image)

**Fig. 104:** Mean free path on the surface of the contour of the spike versus the position from the throat. The mean free path and the distance from the throat are non-dimensionalized with the TH. (a) Comparison between ideal case, the validated model for the basic spike geometry and for the corrected one, and the model with specular boundary condition. (b) Comparison at different operating conditions, with the same pressure ratio.

The mean free path of the flow at the surface \( \lambda_{KBL} \) is plotted against the distance along the surface of the spike in Fig. 104 for different cases. Fig. 104(a) compares different models at the same operating condition of \( p_0 = 18.4 \) mbar and \( p_{\text{amb}} = 0.11 \) mbar. The ideal treatment (calculated with isentropic relations) gives the lowest curve. The case with specular gas-surface interaction shows that it is the closest scenario to the ideal calculations (see also Fig. 102), and the furthest from reality. The slopes of the curves with the validated boundary conditions are instead larger due to the growth of the KBL, and are also approximately linear. Furthermore, the slope of the curve of the ideal geometry and the corrected one are similar, meaning that the slope remains approximately constant for linearly corrected spike geometries. Fig. 104(b) shows the influence of the ambient pressure for the validated case, keeping the pressure ratio constant. The curves still show an approximately linear trend,
and the larger the ambient pressure, the smaller the slope of the curve. The linearity condition of the curves lead to:

$$\lambda_{KBL}(l, p_{KBL}, T_{KBL}) = \psi(p_{KBL}, T_{KBL})l + \lambda_0(p_0, T_0),$$

(63)

where the proportionality factor $\psi$ is a measure of the expansion of the flow along the spike contour and depends on the mean static pressure $p_{KBL}$ and on the mean static temperature $T_{KBL}$ in the KBL; and $\lambda_0$ is the mean free path at the throat of the aerospike, which is at sonic condition and mainly depends on the stagnation flow properties.

Recalling Eq.(58), which was found experimentally, and combining it with Eq.(63) gives a relation between the mean free path in the KBL and the KBL thickness:

$$t_{KBL}(p_{KBL}, T_{KBL}) = \frac{c_{GR}(p_{KBL}, T_{KBL})}{\psi(p_{KBL}, T_{KBL})}(\lambda_{KBL}(l, p_{KBL}, T_{KBL}) - \lambda_0(p_0, T_0)) + t_0.$$ 

(64)

Indeed, Eq.(64) is coherent with the solution to the Kramer’s problem, and the solution of the Kramers problem can be applied also to diverging surfaces based on ideal geometries (constructed for ideal flow expansion) and on their linear corrections. The Kramer’s problem was considered, since it has an analytical solution. The analytical solution states that the thickness of the KBL is in the order of some mean free paths. Further, the proportionality between the KBL thickness and the mean free path is expected because the mean free path relates the statistical number of collisions to the distance travelled by the particle. Considering again Fig. 103(a), the particles that hit the surface are reemitted with a close to random direction, and the direction approaches the x-axis only after a certain number of collisions. Therefore, these collisions happen along a distance from the surface proportional to the mean free path. The complementary way to see this scenario is to consider the particles that are hit by the ones that underwent a collision with the surface. From this point of view, the emission of the particles with a close to random direction implies that additional momentum in z direction is given to those particles. A certain number of particles are needed to absorb the momentum component perpendicular to the surface, and the distance travelled to collide with the needed particles is proportional to the mean free path.
The physical meaning of the constant \( c_{GR}/\psi \) is the statistical number of collisions that are needed to redirect a reflected particle with the core of the jet. Eq. (64) is verified with experiments in Fig. 105, where the dependence of the mean free path on \( t_{KBL} \) at the exit of the aerospike is shown. The experimental data are at different operating conditions with the same pressure ratio. The blue dots represent the experimental values for the ideal contour designed with the method of Angelino. The red triangles represent instead \( t_{KBL} \) for the corrected contour of the spike. On the two sets of data points a linear fit was superimposed, where the slope is equal for both. The linear correction of the spike contour has the effect of lowering the line, but without changing the slope. And, the linear proportionality of Eq. (64) approximately predicts the experiments. The linear correction straightens the jet (as shown in Fig. 96) even though the KBL thickness increases, because the spike divergence is larger than the ideal case, and at the exit the spike contour is lower (see Fig. 94).

![Fig. 105: Measured KBL thickness as a function of the mean free path non-dimensionalized with TH. The experimental data points are all taken at the same pressure ratio, and at the exit of the aerospike. The ideal contour and the linearly corrected one are compared in the same plot. For both profiles, a linear function with the same slope approximately predicts the experiments.](image)

Further questions arise regarding the dependence of the proportionality factors with the condition of the surrounding. The experimental values of the KBL thickness non-dimensionalized by the mean free path in the KBL \( t_{KBL}/\lambda_{KBL} \) are plotted in Fig. 106 against
the ambient pressure, which is assumed to approximate the mean static pressure in the KBL at the exit of the aerospike. The data points are all taken at the exit of the aerospike and with the same pressure ratio. The plot shows that $t_{KBL}/\lambda_{KBL}$ is linearly dependent on the pressure, meaning physically that more collisions are necessary at larger pressures in order to redirect a reflected particle with the core flow.

\[ t_{KBL}/\lambda_{KBL} = \text{linear function of pressure} \]

**Fig. 106:** The ratio $t_{KBL}/\lambda_{KBL}$ from the measurements dependent on the pressure. The experimental data points are all taken at the same pressure ratio, and at the exit of the aerospike. The ratio is approximately linear with the ambient pressure.

From Fig. 106 it is concluded that $t_{KBL}/\lambda_{KBL}$ is a linear function of the pressure in the KBL:

\[ \frac{t_{KBL}}{\lambda_{KBL}} = \frac{c_{GR} \lambda_{KBL} - \lambda_0}{\psi \lambda_{KBL}} + \frac{t_0}{\lambda_{KBL}} = p_{KBL} \chi + \delta, \]  

where $\chi$ and $\delta$ are not dependent on the pressure.

The dependence of the thickness of the KBL on the temperature, which was observed in Section 7.1.3, is quantified in Fig. 107. The experiments show that the relation between the KBL growth and the temperature is approximately linear. Indeed, a term of direct proportionality between $t_{KBL}$ and $T_{KBL}$ is found combining Eq.(64) with Eq.(65), and applying the definition of the mean free path, which states $\lambda \propto T/p$. A larger temperature leads to gas atoms or molecules that move faster, and the space between the particles becomes larger, therefore $t_{KBL}$ also increases.
7.1.6 Impact of the Kinetic Boundary Layer on the Thrust

The thrust is estimated in this Section as a measure of total momentum that the aerospike delivers to the gas curtain. As it was described in the previous sections, the growth of the KBL leads to the lifting of the jet. The lifting of the jet observed in the more rarefied cases in particular is detrimental, because it causes unwanted particle trajectories and less momentum on the surface in the horizontal direction. For this reason, together also with the horizontal momentum of the ejected particles, the thrust is affected by the formation of the KBL close to the surface.

In Fig. 108, the base case with 0.11 mbar ambient pressure and 18.4 mbar stagnation pressure is considered for the ideal and the corrected contour. In the plot, the pressure on the surface was multiplied by the $x$ component of the local normal vector of the surface of the spike $n_x$ and it is plotted versus the distance from the throat normalized by the TH. The surface under the curve is the thrust in axial direction normalized by the width of the aerospike. Over the whole surface, the contour that was corrected is exposed to a larger component of the pressure in the horizontal direction, and a smaller one in the vertical direction. The correction optimizes the main core propagation direction and has a direct effect on the thrust.

Fig. 107: Measured KBL thickness non-dimensionalized with the TH versus the stagnation temperature. The trend is approximately linear.
7.1 Aerospike Performance

Fig. 108: The pressure on the surface of the spike multiplied by the x component of the local normal vector of the surface of the spike $n_x$ is plotted versus the distance from the throat normalized by the TH. The blue line represents the ideal contour, and the red dashed line the corrected one. The corrected contour is exposed to a larger component of the pressure in axial direction over the whole surface.

The particle trajectory results from a momentum balance with the surface. Therefore, if the particles expand with a direction closer to the axial direction as it was shown experimentally in Section 7.1.3, the thrust is also expected to become larger.

In order to quantify the improvement of the thrust, the local horizontal pressure component was integrated and the contribution of the throat cross-section was included. The thrust calculated assuming the ideal behavior of the flow is 1.49 mN/mm. However, the real thrust is 0.97 mN/mm for the uncorrected case (spike geometry directly from the method of Angelino), where the value of the thrust is normalized by the width of the aerospike. Therefore the KBL formation decreases the thrust by 35%. The linear correction, instead, improves the thrust and the specific impulse from the base case by 20%. A further improvement is achieved by heating the flow. At a stagnation temperature of 703 K the thrust was 1.87 mN/mm, improving by a factor of almost 2 compared to the uncorrected case at 293 K.
7.2 Effects on Matter Expansion

The aerospike system with corrected spike geometry was fixed to the box for Si sample exposure, in order to study the effects of the known quasi-parallel high-momentum gas stream on the debris expansion. The surface where the aerospike system was attached was perpendicular to the plasma expansion; therefore also the high-momentum gas stream is quasi-perpendicular to the main direction of expansion of debris. The condition studied is representative for operation, as in the study of the fragments the EUV energy monitor was used to control the shutter placed just before the sample, and was opened only at a given EUV energy per pulse threshold.

7.2.1 Experimental Results

Fig. 109 shows an example of a sample exposed to plasma at 90° from the laser axis, with the end of the spike at 5 mm from the aperture of the sample box, $p_{amb} = 0.11 \text{ mbar}$, $p_0 = 24.6 \text{ mbar}$, $T_0 = 350°$, and an exposure to 6 Mshot. The sample shows that the fragments are successfully deflected by the aerospike system. As the deflection of the fragments can be distinguished from the one of the faster particles at the known operating condition of the gas curtain system, Fig. 109 was used for the calibration of the model exposed in the following Section. The darker region marked with the dashed circle shows where the debris (see Chapter 5) with the highest velocity arrived on the surface, and the continuous circle indicates the region where the majority of the fragments landed. The continuous line is a slightly elongated ellipse, because the fragments have a component of the velocity perpendicular to the normal vector of the surface where the circular aperture was placed. In the frame of reference of the fragments, the latter would see a projection of the circular aperture on a plane with a normal vector parallel to their trajectory: an ellipse.
7.2 Effects on Matter Expansion

![Image with a sample exposed to plasma at a given operating condition of the mitigation system]

**Fig. 109:** Sample exposed to plasma at a given operating condition of the mitigation system. The mitigation system deflects the fragments of a distance indicated by the white arrow.

The velocity profile of the mitigation system at the measurement location was measured with the experimental setup exposed in Section 3.3, and is plotted in **Fig. 110**.

![Graph showing the velocity profile measured at the operating condition of the aerospike system used for the sample exposure shown in Fig. 109]

**Fig. 110:** Velocity profile measured at the operating condition of the aerospike system used for the sample exposure shown in Fig. 109.
7.2.2 Model for Fragment Deflection

A model was developed, in order to estimate the deflection of the fragments with different operating parameters of the aerospike, and tailor the operating conditions for shielding of a wanted surface. The model is based on the measurements of the velocity profile of the gas curtain and on the measured fragment deflection on the samples. In the model a distribution of spherical fragments with various sizes and velocities was considered, and the trajectory from IS to the sample was calculated with the measured gas curtain velocity.

The flow seen by the fragments is in molecular flow regime. Indeed, considering the classical definition of Knudsen number with a fragment diameter of 5 µm as characteristic dimension (one tenth of the larger droplet diameter that was considered) it results in $Kn \approx 400$. Therefore, the drag force on the fragments was modeled after the empirical relation for the rarefied drag coefficient $C_D$ by Henderson.\textsuperscript{221} The algorithms solves the following equations for the calculation of the magnitude of the drag force $F_D$:

$$
C_D = \frac{0.9 + 0.34/M^2 + 1.86(M_\infty/Re_\infty)^{1/2} \left[ 2 + 2/S^2 + 1.058/S_\infty(T/T_w)^{1/2} - 1/S^4 \right]}{1 + 1.86(M_\infty/Re_\infty)^{1/2}},
$$

$$
S_\infty = M_\infty \sqrt{\frac{\nu}{2}},
$$

$$
F_D = \frac{1}{2} \rho \pi (D_f^2) v_{rel}^2.
$$

$M_\infty$ is the Mach number and $Re_\infty$ is the Reynolds number of the flow seen by the droplet with the fragment diameter as characteristic dimension, $D_f$ is the diameter of the fragment, and $v_{rel}$ is the relative velocity as seen in the fragment frame of reference.
7.2 Effects on Matter Expansion

**Fig. 111:** Map of the deflection on the sample for the gas curtain operating condition used for the experiment shown in Fig. 109. The map shows the dependence on the diameter and on the velocity of the fragments. The white line is the mean deflection measured on the sample.

**Fig. 111** shows the result of the model, a map of the deflection depending on the dimension and on the velocity of the fragments for the operating condition of the mitigation system equal to those used for the experiment shown in **Fig. 109**. The model result is linked to the experiment by the white line, which represent the mean deflection measured from the experiments. Therefore, the line indicates the combinations of velocity and dimension of the fragments that underwent the observed deflection. It is interesting to point out that there is a non-dimensional number that characterizes the deflection, which is the Weber number $We$. Originally, the Weber number relates the kinetic energy of the droplet relative to a flow or a surface to the surface tension force of the droplet. In this case, the proportionality to the kinetic energy is still needed, but the cross-sectional area of the droplet substitutes the surface tension force. **Fig. 112** shows the map of $We$ for fragments with different diameters and velocities together with the white line representing the mean deflection from the measured sample. Indeed, $We$ is constant along the constant deflection line and equal to around $3’200$. Therefore, $We = 3’200$ is used for the characterization of the fragments at the operating condition of the exposure.
Fig. 112: Logarithmic map of the Weber number calculated for the fragments with a given velocity and dimension. The constant Weber number contour indicates also a constant deflection, exemplified by the white line.

The fragments that arrived on the sample splashed and solidified, therefore the diameter seen on the sample is larger than the one of the fragments. To estimate the splashes diameter with the fragment dimension and their velocity, a further model was implemented. The latter was based on the analytical expression for the droplet maximum spread factor of Pasandideh-Fard et al. The expression was tested with experimental results and showed an error of less than 15%. The maximum spread factor \( \xi_{\text{max}} \) is the ratio of the maximum diameter reached in the splash \( D_{\text{max}} \) and the droplet diameter that in this case is the diameter of the fragment coming from the main droplet target \( D_{fr} \). The model at the limit for large \( We \) (\( We \gg 12 \)) reduces to:

\[
\xi_{\text{max}} = \frac{D_{\text{max}}}{D_{fr}} = 0.5 Re^{0.25}.
\]

This model applied here is intended to give an estimation of the process, and in particular the dependence on the velocity; a more refined treatment should take into account the transient cooling of the splash during deformation on the surface and its solidification.
7.2 Effects on Matter Expansion

Fig. 113: Estimation of the diameter of the splash for different diameters of the fragments and velocities.

Fig. 114: SEM image of a splash of a fragment seen on the side. The contour of the splash is highlighted by the white line. The mean height is smaller than one tenth of the splash diameter.

The result of the model for the prediction of the diameter of the splash alone is shown in Fig. 113. The larger the kinetic energy and the diameter of the fragment are, the larger the splash diameter is expected to be. To verify that the model prediction of the splash diameter was in the right order of magnitude a sample was sectioned after exposure and analyzed under the SEM to image the height of the splashes and their radius. An example is
shown in Fig. 114, where the mean height of the splash is smaller than one tenth of the diameter, also considering the fact that the side picture shows the border of the splash, which is normally higher than the middle due to partial recoil. The dimensions result in spread factor between 2.5 and 4, and using the model to an impact velocity between 20 m/s and 100 m/s.

The result of the overall model is shown in Fig. 115, where the diameter of the splash of the fragments on the impact surface (the sample) is directly plotted. Following the white line, which represent the mean deflection experimentally observed, and with a splash diameter between 5 μm and 50 μm, the fragment velocity ranges from 150 m/s to 400 m/s. This range of velocities is more plausible than what given by the spread estimation for tin laser produced plasmas. The underprediction of the velocity with the study of the spreading could be due to the recoil before solidification of the tin fragment after the maximum extension on the sample surface, which leads to an underprediction of the maximum spread factor.

![Fig. 115: Map of the deflection of the sample as a function of the diameter of the splash directly observed on the impact surface. The white line represents the observed mean deflection.](image-url)
7.2.3 Fragments Deflection Depending on Gas Curtain Operating Condition

The mean fragment Weber number deduced by the measured deflection on the surface of the sample with the known velocity profile of the mitigation system is used to estimate the deflection at other measured operating conditions. In particular, the operating conditions are at different distances along the direction of the flow injected by the aerospike in the vacuum vessel, at different stagnation pressures and stagnation temperatures of the flow. The base operating condition of the mitigation system was at $p_{0,b} = 18.4$ mbar, $T_{0,b} = 300$ K.

Deflection versus Distance along the Flow Expansion

The deflection of the fragments along the flow expansion is important for the use of the mitigation system with large collector optics or large apertures. In order to protect the largest portion of the surface/hole, the mitigation system should deflect the fragments along the whole length of the surface/hole.

Fig. 116 shows the trajectory of the fragments at different distances from the end of the spike. The dimensions are normalized with the throat height of the aerospike. The mitigation system has a double effect: it shifts the fragments, and it imposes a component of the velocity perpendicular to the initial velocity (seen by the slope of the curve). The two different effects are partially separated in Fig. 117, and Fig. 118. The first figure represents the deflection at the level of the aerospike exit and of the aperture, hence directly the shift of the accelerated fragments. Instead, the second figure is the deflection 50 mm (71 TH) from the aperture. The deflection is larger because it takes advantage of the component of the velocity perpendicular to the initial velocity.
Fig. 116: Trajectory of the fragments with $We = 3'200$ at different distances from the end of the spike. The two axes have different scales. The coordinates are normalized with the throat height (TH) of the aerospike.

The deflection at the aperture level increases monotonically with the distance along the spike, this is visible also from the trajectory of the fragments. The increase is due to the thicker layer of the gas curtain. The overall momentum of the flow remains approximately the same, but it begins to act on the fragments at larger distances from the aperture, therefore the velocity component of the fragments along the flow has more time to influence the trajectory. Instead, the deflection of the fragments below the aperture is given mainly by the overall momentum of the flow, and reaches a plateau at around 15 TH from the exit. In the first 15 TH the fragments are too fast to appreciate the whole momentum of the gas curtain.
Fig. 117: Deflection of the fragments along the direction of the gas curtain from the end of the spike. The deflection was calculated as detected directly at the level of the exit of the aerospike, which corresponds to the entrance plate in the experiments of sample exposure.

Deflection versus Stagnation Pressure

The stagnation pressure was changed in the flow field measurements keeping the pressure ratio between the stagnation pressure and the ambient pressure constant and equal to the
design ratio $p_0/p_{\text{amb}} = 158$. On one hand the mass flow increases proportionally to the stagnation pressure, on the other hand also the profile of the flow field strongly changes due to the kinetic boundary layer as exposed in Section 7.1. The combined effect on the trajectory of the fragments is shown in Fig. 119. Comparing for example the profile at 4.4 $p_{0,b}$ with the other profiles, the initial trajectory (from the top) at the larger stagnation pressure is less affected, since the flow is confined closer to the surface. But, as the fragments approach the surface the deflection becomes larger due to the concentrated high-momentum flow.

As in the previous Section, the deflections at the level of the spike and further away are represented separately; see Fig. 120 and Fig. 121, respectively. The deflection at both the positions increases monotonically with the stagnation pressure. In particular at the sample height (Fig. 121), the deflection increases of approximately one order of magnitude with an increase in the stagnation pressure of around 6 times.
7.2 Effects on Matter Expansion

Fig. 120: Deflection of the fragments at different stagnation pressures with a constant pressure ratio. The deflection was calculated as detected directly at the level of the exit of the aerospike, which corresponds to the entrance plate in the experiments of sample exposure.

Fig. 121: Deflection of the fragments at different stagnation pressures with a constant pressure ratio. The deflection was calculated at the position where the samples were exposed, at 50 mm (71 TH) from the aperture.

**Deflection versus Stagnation Temperature**

The stagnation temperature was changed keeping the mass flow constant. In this way, the momentum of the flow was changed only by the velocity, which depends on the square
root of the stagnation temperature. The effect on the flow field is not only an increase in the velocity, but also an increase in the static temperature. The latter is directly related to the thermal velocity, and as it increases, the flow expands more in directions perpendicular to the main expansion leading to a wider profile in vertical direction (see the profiles in Fig. 97). Therefore, the increase of stagnation temperature affects the fragment trajectory depositing more momentum on the one hand, and by acting further away from the surface on the other hand. Both the effects increase the deflection as shown in Fig. 122. The increase of the width of the gas curtain profile with the stagnation temperature has to be considered together with the position of the falling tin droplets with respect to the aerospike location.

![Figure 122: Trajectory of the fragments with We=3'200 at stagnation pressures at the same mass flow. The two axes have different scales. The coordinates are normalized with the throat height (TH) of the aerospike.](image)

The deflections at the level of the aerospike exit, which is the same as the entrance plate, and at the sample level are plotted in Fig. 123, and Fig. 124, respectively. As with the increase of stagnation pressure, also increasing the stagnation temperature increases monotonically the deflection. It is interesting to notice that the increase of the deflection of the fragments for a given increase in temperature $x$ is approximately equal to the increase of deflection that would be produced increasing the pressure of $\sqrt{x}$. The relation is coherent.
with the dependence on the momentum, directly dependent on the stagnation pressure, as it is proportional to the mass flow, and to the velocity, hence to the square root of the temperature.

![Graph](image1.png)

**Fig. 123:** Deflection of the fragments at different stagnation temperatures with constant mass flow. The deflection was calculated as detected directly at the level of the exit of the aerospike, which corresponds to the entrance plate in the experiments of sample exposure.

![Graph](image2.png)

**Fig. 124:** Deflection of the fragments at different stagnation temperatures with constant mass flow. The deflection was calculated at the position where the samples were exposed, at 50 mm (71 TH) from the aperture.
7.3 Considerations on Radiation Source Design

The different types of debris that have detrimental effect on components of an EUV source and on the next stages are mainly neutrals (see Section 5.4.2), fragments (see Section 6.1.3), and ions (see Section 5.1.2). Depending on the pressure of the ambient gas, ions might recombine and become neutrals with high kinetic energy (see Section 5.4.3). With argon gas, a pressure larger than $6 \times 10^{-2}$ mbar and a distance larger than 150 mm are enough to stop slow neutral particles (with velocities of around 5 km/s). The previous Section shows that the mitigation system developed in this work successfully deflects the fragments of a distance, which depends on the operating condition. Neutrals with high kinetic energy from ion recombinations are the only specie that needs a larger pressure in the vacuum vessel than what was considered in this work. The extrapolation of the mitigation effectiveness for neutrals with high kinetic energy with argon ambient gas indicates a pressure in excess of 1.6 mbar to stop 99.73% (3σ).

![EUV transmission at different angles and path lengths](image)

**Fig. 125: EUV transmission at different angles and path lengths.**

The gas curtain was tested computationally also with neutrals with high kinetic energy, but the effect around the nominal pressure of 0.1 mbar was negligible because of the dominant velocity of the recombined ions with respect to the gas curtain particles (one order of magnitude difference). On the other hand, the ambient pressure cannot be increased as
needed for debris mitigation. Indeed, at the pressure range of millibars, EUV is significantly absorbed, as shown in Fig. 125, in particular with path lengths typical of a high-brightness EUV source, larger than 400 mm.\textsuperscript{61}

A strategy for the design of a high-brightness source, on the basis of the measurements performed in this work, follows. The collector mirrors should be placed in an enclosure as already shown by Giovannini \textit{et al.},\textsuperscript{61} and Ellwi \textit{et al.},\textsuperscript{63} in order to locally protect them from fragments and slow neutrals with the developed mitigation system placed at the aperture, on the outside. The ambient pressure should be larger than \((9/d_{s})\) mbar, where \(d_{s}\) is the distance between the IS and the aperture of the enclosure and it is in millimeters. If the condition is satisfied the slow neutrals should be deflected before arriving at the aperture. \(d_{s}\) should be larger than the expansion of the gas curtain in the direction of the plasma, in order not to compromise the jet stability. The extrapolation of the velocity profile of the aerospike during operation at base condition with the corrected spike geometry lead to a impact of the aerospike on the ambient gas up to around 50 TH from the base of the aerospike at a distance from the end of the spike of 30 TH (see Fig. 96). The choice of the angular position of the collector mirror with respect to the laser axis is a tradeoff between EUV emission and debris load. The position around 90° would allow the minimization of the recombined ions with high kinetic energy (see Fig. 47), without entering the region of large fragment load (at angles larger than 105°). And, the EUV energy per pulse centered at 13.5 nm in the EUV +/-2% BW would still be 0.3 mJ/sr at 2 x \(10^{11}\) W/cm\(^2\) for both 30 μm and 50 μm droplet diameter, and of around 0.25 mJ/sr at 5 x \(10^{10}\) W/cm\(^2\) with 30 μm droplet diameter (see Fig. 34).

The aperture of the enclosure should be narrower in the direction of the flow of the gas curtain, in order to facilitate the deflection of the fragments along the way between the said aperture and the ML collector. The pressure in the enclosure should be as low as possible in order to maximize the EUV transmission, eventually with the circulation of other gases with a lower EUV absorption. The presence of the high momentum flow at the aperture is expected to partially prevent the inflow of ambient gas particles. The low pressure in the enclosure should be used to place the collector mirror at a distance in the order of hundreds of millimeter, e.g. at a \(p_{0} = 4p_{0}\) and a first collector mirror of 25.4 mm in diameter, the distance should be around 360 mm from the IS to avoid fragments. In order to grant no
debris to later stages, a module between the source and the following stage should be built in order to host a second aerospike system. The module should be at a larger pressure, and at the inlet as well as at the exit differential pumping systems should be installed (patent-pending)\textsuperscript{62}.

### 7.4 Summary

A mitigation system based on inertial collisions was designed and optimized to deliver a high-momentum gas curtain in the rarefied condition of the vacuum chamber. The acceleration of the gas curtain was performed with an aerospike. The rarefaction of the flow on the surface of the spike during acceleration caused the growth of a kinetic boundary layer and the elevation of the jet in a direction perpendicular to the surface. A joint experimental, numerical and phenomenological characterization of the KBL over diverging surfaces (generic case of an aerospike) has been presented. The flow field close to the surface and downstream of its end was measured, and the thickness of the KBL was derived along the surface at different ambient pressures from $p_{\text{amb}} = 7 \times 10^{-2}$ mbar to $p_{\text{amb}} = 1.5$ mbar. Over the whole pressure range, the condition of the flow close to the surface, where the KBL forms, is in non-equilibrium. The physics used for the description of the standard viscous boundary layer is therefore not applicable.

A phenomenological description of the thickness of the KBL is produced by coupling the results from the Kramers problem with the observation from the experiments and the computations. The thickness of the KBL is proportional to the mean free path, and to the length along the surface of the spike from the throat. Furthermore, it was found that the constant of proportionality between the KBL thickness and the mean free path is a linear function of the pressure; meaning that the mean number of collisions experienced by the particles in the KBL increases linearly with the pressure. Computationally, the importance of the gas-surface interaction for the growth of the KBL was shown. The average roughness of the surface was changed experimentally from 0.023 $\mu$m to 14.157 $\mu$m, whilst the change in the horizontal velocity profile and in the thickness of the KBL was negligible (less than 2% of the horizontal velocity and within the measurements error). The average roughness
of the spike surface could not be used to affect the gas-surface interaction and the flow field. Instead, the linear growth rate of the KBL was used to correct the spike geometry and the effects were studied experimentally. The result was an optimization of the flow field, decreasing the jet height of 10 TH after 30 TH from the exit. The correction lowered the jet from 14.4° to 4.8° along the x-axis, demonstrating how to tune the main direction of propagation of the jet, which is important for all the considered applications. The ratio between the KBL thickness and the mean free path did not vary significantly between the ideal and the corrected spike geometry, and can be used as an invariant in the optimization of the geometry of the spike through linear corrections.

The thrust as a measure of the overall momentum that the aerospike delivers to the gas curtain is quantified. For the base case, the thrust was 0.97 mN/mm, which was 35% lower than the one predicted by ideal calculations. Instead, an improvement of 20% was achieved by correcting the geometry of the spike, and by heating the gas flow to a stagnation temperature of 703 K the thrust was approximately doubled with respect to the real uncorrected case in cold gas testing.

The effect of the gas curtain on the expanding debris from droplet-based LPP was characterized with the EUV-controlled exposure of Si samples. The aerospike system could successfully deflect the fragments ejected by the droplet target. The measurements were used to estimate the velocity distribution of the fragments, leading to a mean velocity between 150 m/s and 400 m/s depending on the fragment diameter. The deflection of the fragments using the gas curtain was found to be invariant with respect to velocity and diameter for the same We number. The deflection of fragments with different operating conditions of the aerospike system is estimated. The deflection at two different positions is estimated, namely at the outlet of the aerospike and 71 TH downstream, in the enclosure. The deflection of the fragments in the enclosure increases with the distance from the throat for the first 20 TH after the exit, and then remains approximately constant. Increasing stagnation pressure and temperature leads to monotonic increase of the deflection. The deflection increases of approximately one order of magnitude with an increase in the stagnation pressure of around 6 times, and with a double stagnation temperature the deflection increases of around 50%.
A strategy for source design is to place the collector mirrors in an enclosure, protect it with an aerospike at the inlet, have a pressure in the vessel of at least \((9/d_s)\) mbar, and leave a distance between aerospike and droplet targets of at least 50 TH. The aperture at 90° from the laser axis is a compromise between debris and EUV emission. In the enclosure, the pressure should be as low as possible, and at the exit of the source a second aerospike system should be integrated in the differential pumping system to further protect the next stages.
7.4 Summary
Chapter 8

CONCLUSIONS AND RECOMMENDATIONS FOR FUTURE STUDIES

A series of experiments and of analytical and computational modeling was performed to characterize the droplet-based laser-produced plasma in terms of EUV emission and matter expansion. The experimental measurements of EUV emission, matter in form of ionic populations, and fragments are believed to be the first angular-resolved data around droplet targets at laser intensities representative for an EUV source.

In this final Chapter the main conclusions of the present study are presented, whereas the summary of the results is given at the end of each Chapter.

8.1 Conclusions

8.1.1 EUV Emission

EUV emission from the droplet-based LPP was measured for the first time in three dimensions up to 120° from the laser axis, demonstrating that EUV radiation can be collected also behind the droplets. Three-dimensional analytical models are developed in
8.1 Conclusions

this work, tested, and successfully used to relate the angular distribution of EUV to the density and temperature profile of the plasma. The models extrapolate EUV emission up to 150°. The angular distribution was shown to differ significantly from slab targets mainly because of the formation of sidelobes. The plasma expansion at experimental conditions is found to be anisotropic. The condition for anisotropic plasma expansion is \( \phi_{\text{las}}/2 \tau_{\text{las}} c_s < 1 \).

The developed analytical models are tools that are also applicable to other cases. In general, to link radiation emission distribution with density distribution.

Experimental parametric studies on EUV angular distribution led to the following conclusions. The use of larger droplets leads to a more forward-peaked angular distribution of the EUV emission with larger emission close to the laser axis. Instead, with lower irradiance the emission close to the laser axis decreases, but shows a flatter EUV emission at larger angles, and the conversion efficiency increase. The difference between the distributions comes mainly from the amount of mass ablated during irradiation. A novel procedure is developed to estimate the angular distribution of EUV emission with a given laser irradiance, laser wavelength and droplet diameter. EUV emission and plasma dynamics during EUV emission is not significantly affected by the presence of ambient gas up to millibars.

8.1.2 Matter Expansion

The contribution of different ion charge states in a LPP was discriminated for the first time around a droplet target in the presence of ambient gas. The expansion of Sn\(^{2+}\) ions is forward peaked and faster than Sn\(^+\) ions, which are quasi-isotropic.

Models were developed to highlight the dominant processes during matter expansion and calibrated with experiments. The ablation process needs a negligible amount of the laser energy (around 5%) for droplets aligned with the laser axis. The maximum absorption of laser energy is at around \( 1 L \), from the target surface, closer to the surface than the EDR.

The region from where ions are accelerated is centered with the absorption region, therefore a shorter laser wavelength would lead to a more isotropic ion charge state, as opposed to a longer wavelength. During acceleration the mean ion charge ranges from 5 to 2.
The ambient pressure has multiple consequences on the ionic expansion: partial deflection through inertial collisions, loss of charge through CEX collisions, and loss of charge through recombinations. The latter mechanism is expected to be accelerated by electrons donated by ambient gas atoms, which are ionized through high-energy photons from the plasma, electrons and ionizing collisions with ions.

The ablation process is affected by offset between droplet and laser axis, and by the \( \frac{D_{dl}}{\phi_{las}} \) ratio. The main consequences are a different direction of initial expansion of the vapor and plasma, and a loss of laser energy. Provided that \( \tau_{las} > \sqrt{2} \Delta l/c_s \) the plasma should still absorb the main part of the laser energy. The main direction of EUV emission and ion expansion are correlated, which demonstrates that opacity effects are dominant in EUV emission distribution. Electrons are more mobile than ions and the correlation is weaker.

The load of fragments from the free-falling droplets changes dramatically with the angular position from the laser axis. The large load increase from around 105° is in agreement with angle of recoil momentum calculated from the ablation process with offset. Fragments ejection increases with ambient pressure mainly due to the larger evaporation rate of tin vapor on the droplet surface.

### 8.1.3 Mitigation

A debris mitigation system was developed, tested, and optimized. Deflection of debris through high momentum gas was the chosen mitigation strategy. An aerospike nozzle operating up to rarefied condition was designed and its flow investigated experimentally, computationally and analytically. Kinetic boundary layer growth caused the gas jet to lift from the surface, expanding along an unwanted direction. Optimization of the geometry led to an order of magnitude reduction of the jet-lifting angle, making the aerospike-based debris mitigation system applicable to the EUV source without disturbing the falling droplets.

The debris mitigation system once tested in the EUV source proved to successfully deflect the droplet fragments. Knowing the flow characteristics and the fragment deflection, the velocity of the fragments was estimated to be in the order of hundreds of meter per second. Larger stagnation pressure and temperature increase the deflection capability.
The debris mitigation strategy for an EUV source for metrology and inspection could be optimized. The collector mirrors should be located in an enclosure with an aerospike system at the aperture, and at an angle with the laser axis of around 90°. The argon ambient pressure in the vacuum vessel should be of at least $6 \times 10^{-2}$ mbar in order to stop the slow neutral particles along 150 mm. The distance between the aerospike system and the droplets should be at least 50 TH. The pressure in the enclosure should be as small as possible. A second aerospike system should be located at the exit of the EUV source, close to the intermediate focus, to further protect the next stages.

8.1.4 Experimental Facility

A further accomplishment of this work is the coordination of the overall design and the definition of the interfaces between the main subsystems of the ALPS 2 laboratory; and the design and commissioning of its main components as the vacuum system, the debris mitigation and collector systems, the beam delivery system, and the laboratory infrastructures. The use of this facility enabled the majority of the experimental study of this work, of parallel projects, and is going to be used for future projects.

8.2 Recommendations for Future Studies

Measurements performed in this work show that emission of radiation and matter expansion depend on the droplet position relative to the focal spot and on their size. The capabilities of the ALPS 2 laboratory to synchronize droplet position and size measurements at the focal spot with the laser pulses should be used to develop the understanding of how EUV emission and matter expansion change with the offset. Further parameters to be studied with the offset are the laser irradiance and the laser pulse duration, in order to maximize the conversion efficiency and the source stability, and minimize the debris load.

Experiments on matter expansion in presence of ambient gas could be repeated with different fuel/gases combinations, in order to further analyze the influence of CEX collisions and ambient gas ionization on the recombination of the ions of the target
material. In particular, in using the ESA with background gases with high ionization potentials (for example He) and tin droplet targets would be interesting to investigate whether Sn$^{3+}$ ions are visible due to the higher ionization potential of the background gas compared with the one of Ar used in this work.

A droplet charging and steering system$^{228}$ could be integrated in the dispenser, in order to fine adjust the droplet position, and to accelerate the droplets during fall from the nozzle to the irradiation site. Furthermore, the effect of an already charged droplet on matter expansion could be investigated, in particular the influence on the fragment distribution.

The tailored spectral emission distribution in space around the droplet targets centered at other wavelengths could be tested experimentally with the energy monitor. Depending on the wavelength, another combination of multilayer mirror and filer should be used in the energy monitor, and if no multilayer combination is available a band-pass filter or a system of filters should be used. The constants in the analytical models could be adjusted with experimental data to predict the emission distribution. The longer the measured radiation wavelength the larger the emission dominant region is expected to be, arriving at the limit to surround the main plasma, which would lead to an almost isotropic emission.

Different droplet target materials could be studied for the fragments formation. For example, the critical tension for spallation on and the critical thermodynamic temperature for explosive boiling are material dependent. Therefore, materials with higher critical tension and critical thermodynamic temperature could lead to a reduction of the fragmentation.

In-situ and pulse-to-pulse debris measurements could be implemented to directly measure the debris load depending on various laser parameters and on droplet position and size at different angles from the laser axis. A possible detection mechanism could be based on transmission measurements on a transparent sample that is exposed to the plasma. A constant laser beam would shine through the sample, and its transmission would be monitored by a photodiode. The decrease of intensity measured by the photodiode would be directly related to transmission loss, and with the target optical constants to thickness of the coating. The same system could be used to confirm the debris mitigation system effectiveness with time-resolved measurements during the source run-time. Parametric
8.2 Recommendations for Future Studies

studies varying source and debris mitigation system operating conditions using such a measurement system would fully characterize the debris characteristics.

The loss of reflectivity of EUV radiation due to debris should be quantified directly with ML samples and with the capabilities of ALPS 2. The reflectivity measured before and after sample exposure together with EUV generation and droplet position per pulse would give a direct estimation of the collector lifetime. The debris mitigation effectiveness in protecting the ML collector could be directly tested. Furthermore, the angle of incidence of the incoming debris on the ML sample surface could be varied in order to simulate different position on an elliptical collector, and grazing incidence collector could also be tested. The latter parametric study would be interesting to quantify the angle of incidence has an influence on debris sputtering and deposition. The suggested parametric studies on collector lifetime should be an input in the estimation of the operating costs of the EUV source. Together with the costs for replacing and cleaning the ML collector it should lead to the choice of the best strategy.

Charge-exchange collisions between target material and ambient gas could be implemented in the BGHP code to consider this mechanism in the loss of charge of the ions. Furthermore, the presence of more electrons due to ambient gas ionization could be implemented in the modeling together with an ionization-recombination model, e.g. CR model. The implementations would allow a direct comparison between the time-resolved debris measurements and the simulations. A further extension of the BGHP code to three dimensions would allow simulations at offset conditions.
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APPENDIX

List of Publications

Journal Articles

Published


Submitted
A. Z. Giovannini, I. Barendregt, T. Haslinde, C. Hubbs and R. S. Abhari, Self-confined plasma in a magneto-plasma compressor and the influence of an externally imposed magnetic field, Plasma Sources Science and Technology

A. Z. Giovannini and R. S. Abhari, Rarefied flow expansion along diverging surfaces, Physics of Fluids
Patents


Conference Publications


