



## Experimental Plasmas : Lecture notes

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# Plasmas expérimentaux UE 896 – M1 Physique

Rédigé par LEDIG Jordan

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## Overview and notation modalities

- Course on experimental plasmas (6hrs).
- Hands-on project (24hrs, or 3 times 8hours): 1 day Langmuir probe measurements on the Helicon reactor with Mr. De Poucques — 1 day on Laser Induced Fluorescence (LIF) on the magnetron reactor with Mr. El-Farsy — 1 day LIF data analysis and report redaction with Mr. Ledig.

You will have to write a report for all projects. It will contain a brief introduction on the plasma device, a part to expose results, and a part for discussions. To be simple, it should look like an article. The following 6hours course will give you the scientific background to understand both experiments, to help you for the comprehension and the redaction.

Note that this paper was meant to be my course support during class. But after several hours of  $\text{\LaTeX}$  compilations, plotting, and Tikz-ing (making the schemes) I decided to give a copy to you. It may contains some (many!!) grammatical and some (yes,some) mathematical errors, so let me know if you find, one. Or two.Or even more.

# 1 Langmuir probe measurements on the Helicon reactor

## 1.1 Experimental device

### 1.1.1 Generalities on the device

Please note that this section is not very complete because during the day of experiment L. De Poucques will explain it more precisely – and with the device in front of you. If you want some more details take a look at X. Glad's and T. Bieber's thesis [1, 2], chapter 2 for both.

The helicon is an experimental device which creates a plasma using a RF antenna ( $\nu_{\text{RF}} = 13.56$  MHz). The reactor is a superposition of two cylinder:

- the source chamber (radius 75 mm – height 300 mm) placed above. This chamber contains a Pyrex tube (radius 65 cm – height 300 mm), and around this tube is the RF antenna. The Pyrex tube protects the antenna from ion flux and from the temperature.
- the diffusion chamber (radius 130 mm – height 260 mm) placed bellow.

There are two coils around each chamber to create a magnetic field. We call  $B_0$  the field from the source chamber, and  $B_{\text{diff}}$  the field from the diffusion chamber. The turbomolecular pump generate a void of about  $10^{-7}$  Torr<sup>1</sup>.

### 1.1.2 Creation of the plasma

Obviously, the plasma is generated in the source chamber and diffuses along the magnetic field lines toward the diffusion chamber. It is the antenna that creates the plasma because it is crossed by an oscillating current at  $\nu_{\text{RF}}$ . At our pressure range (1 to 500 mTorr) ions do not respond to the RF oscillation, because:

$$\omega_{p,i} \ll \omega_{\text{RF}} \ll \omega_{p,e} \quad \text{where } \omega_p = \sqrt{\frac{ne^2}{m\varepsilon_0}} \text{ is the plasma pulsation} \quad (1)$$

We also call this "electron heating".

To ignite the plasma, electrons already present in the neutral gas responds to the strong RF electric field created by the hight potential variations (several kilovolts) near the antenna. Those electrons are accelerated and create new electrons thanks to collisions with gas atoms. Collision after collision, the plasma is generated, and there is the apparition of a sheath between the antenna and the bulk plasma to screen the RF oscillations inside the plasma (so that only the sheath reacts to the RF electric field). This collisional heating is the capacitive mode.

This mode is only used to ignite the plasma because the power absorbed by electrons decreases when  $n_e$  (electron density) increases. Indeed, taking Euler equation for homoge-

<sup>1</sup>The Torr (or mmHg) equivalents to 133.32 Pa. The name comes from Torricelli, inverter of the barometer.

neous plasma,

$$m_e n_e \frac{d\mathbf{u}_e}{dt} = -en_e \mathbf{E} - m_e n_e \nu_{\text{coll.}} \mathbf{u}_e \quad (2)$$

and now supposing that  $\mathbf{u}_e$  and  $\mathbf{E}$  varies with time like  $\mathbf{X} = |\mathbf{X}|e^{j\omega_{\text{RF}}t}$ , Euler becomes

$$j\omega_{\text{RF}} m_e n_e \mathbf{u}_e = -en_e \mathbf{E} - m_e n_e \mathbf{u}_e \nu_{\text{coll.}} \Leftrightarrow \mathbf{J}_e = -en_e \mathbf{u}_e = \omega_{p,e}^2 \frac{\varepsilon_0 \mathbf{E}}{\nu_{\text{coll.}} + j\omega_{\text{RF}}} \quad (3)$$

hence,

$$\langle \mathcal{P}_{\text{capa. mode}} \rangle_t = \langle \mathbf{J}_e \cdot \mathbf{E} \rangle_t = \frac{1}{2} \text{Re} \{ \mathbf{J}_e \cdot \mathbf{E}^* \} = \frac{1}{2} \frac{\omega_{p,e}^2 \nu_{\text{coll.}} \varepsilon_0 |E|^2}{\nu_{\text{coll.}}^2 + \omega_{\text{RF}}^2} \propto \frac{1}{n_e} \quad (4)$$

Once the plasma is ignited and that the electron density is sufficient, it can enters the inductive mode. Indeed the RF current generates a longitudinal varying magnetic field  $\mathbf{B}_{\text{RF}} = B \cdot \mathbf{e}_z$ ; and this varying magnetic field generates an azimuthal electric field  $\mathbf{E}_{\text{RF}} = E_{\text{RF}} \cdot \mathbf{e}_\theta$  according to Maxwell-Faraday equation,

$$\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \quad (5)$$

This electric field speeds up electrons and establish a current flow that counter the initial  $\mathbf{B}_{\text{RF}}$ : this is the induction principle. The generated plasma follows the static magnetic field and enters the diffusion chamber, where we are going to make some measurements on it.

Note that this quick introduction is very gross, and in reality is is more complex, but is beyond the scope of this course.

## 1.2 Langmuir probe measurements

### 1.2.1 The sheath formation

To understand properly Langmuir probe measurements, one must be clear with the sheath structure. After that, calculations are straightforward. Let us be a semi infinite 1D-plasma. A wall is placed at  $x = s$  and has an electric potential of  $\phi_W$ . The plasma is in the negative  $x$  region and  $\phi(x \rightarrow -\infty) = \phi_p$  is the plasma potential.

In the absence of wall, the plasma is quasi-neutral so that  $n_e \simeq n_i$ . According to Maxwell-Gauss law, this implies that there is no electric field in the plasma ( $\nabla \cdot \mathbf{E} = 0$ ). But since there is a wall, the potential must vary from the plasma potential, to the wall potential. This region where the quasi-neutrality is broken and where  $\phi \neq \phi_p$ , is called the sheath.

To understand what happens on the wall, assume we put it into the plasma at time  $t = 0$ . Since electrons are much faster than ions, they will reach the wall at first - and charge it negatively, i.e.  $\phi_W < 0$ . Thus, in the front of the wall there is then a lack of electrons (since they are on the wall) so  $n_e \leq n_i$ . Ions are then accelerated toward the wall, and electrons are repealed. At some point between the wall and the bulk plasma, we assume that the potential and its derivative (the electric field) are 0. We put this point to  $x = 0$ , so:  $\phi(0) = 0 = \phi'(0)$ . We also call this point the "sheath edge".

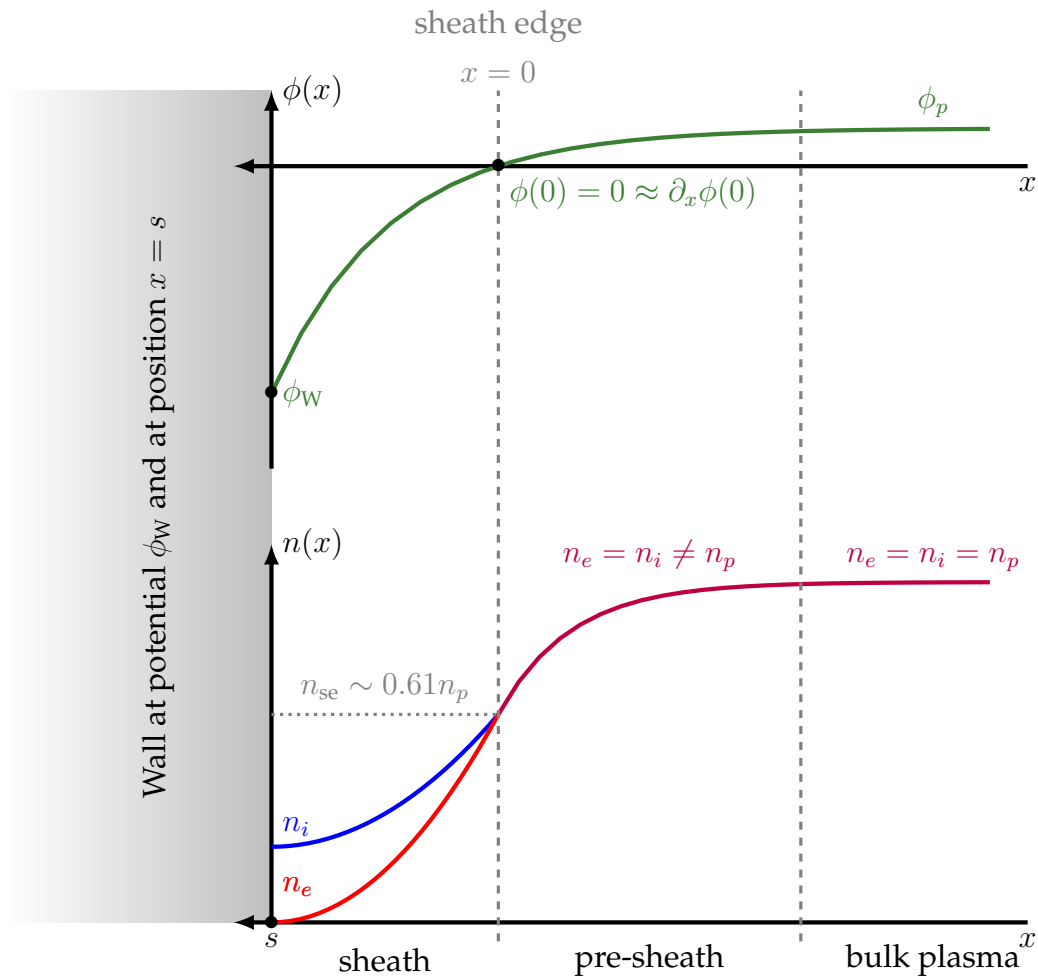


Figure 1: Scheme of our sheath (for the cold-ions model) in a positive space-charged configuration. Plots of densities and electric potential.

We assume the sheath to be collision-less, the ion flux is then conserved since sources and losses are balanced :

$$\frac{\partial n_i}{\partial t} + \nabla(n_i \mathbf{u}_i) = \nabla(n_i \mathbf{u}_i) = 0 \Leftrightarrow n_i(x)u_i(x) = \text{Cte in the sheath} \quad (6)$$

**Exercise.** Write down Poisson's equation, energy conservation and flux conservation for ions in the sheath. Show that we must have at sheath edge,  $u_i(0) = C_s$ . Write  $C_s$  in term of  $k_B$ ,  $T_e$  and  $m_i$ .

**Solution.** First of all, let us write the 3 equations. The conservation laws will be taken between the sheath edge and an arbitrary point  $x$  in the sheath, hence:

$$\left\{ \begin{array}{l} \frac{d^2 \phi}{dx^2} = \frac{e}{\epsilon_0} (n_e - n_i) \\ \frac{1}{2} m_i u_i^2(0) = \frac{1}{2} m_i u_i^2(x) + e \phi(x) \\ n_{se} u_i(0) = n_i(x) u_i(x) \end{array} \right. \quad (7)$$

Using ion energy conservation we can write  $u_i(x)$  in terms of potential. Putting this inside ion flux conservation gives an expression for ion density, that we can put inside Poisson:

$$u_i(x) = \sqrt{u_i^2(0) - \frac{2e\phi(x)}{m_i}} \Rightarrow n_i(x) = n_{se} \left(1 - \frac{2e\phi(x)}{m_i u_i^2(0)}\right)^{-1/2}. \quad (8)$$

In the sheath, the quantity  $n_e - n_i$  must be negative<sup>2</sup> (because there are less electrons than ions in the sheath). So, very close to sheath edge (i.e. at very small potential),

$$\begin{aligned} n_e - n_i < 0 &\Leftrightarrow n_{se} \exp\left[\frac{e\phi}{k_B T_e}\right] - n_{se} \left(1 - \frac{2e\phi(x)}{m_i u_i^2(0)}\right)^{-1/2} < 0 \\ &\sim n_{se} \left(1 + \frac{e\phi}{k_B T_e}\right) - n_{se} \left(1 + \frac{e\phi}{m_i u_i^2(0)}\right) < 0 \\ &\Leftrightarrow \frac{1}{k_B T_e} - \frac{1}{m_i u_i^2(0)} > 0 \quad \text{remember } \phi < 0 \text{ in the sheath} \end{aligned}$$

Which gives the well known Bohm criterion,

$$u_i(0) > C_s = \sqrt{\frac{k_B T_e}{m_i}} \quad (9)$$

$C_s$  is the Bohm speed, or the speed of sound for ions.

That means that to ensure stability of the sheath, ions must enter the sheath at the minimal speed of  $C_s$ . To calculate  $n_{se}$ , one must remember that electrons follow a Maxwell-Boltzmann distribution (i.e.  $n_e(x) = n_p \cdot \exp[e(\phi - \phi_p)/k_B T_e] = n_{se} \cdot \exp[e\phi/k_B T_e]$ ), and at sheath-edge  $n_i(0) = n_e(0)$ . Moreover, between sheath-edge and the bulk plasma, the energy conservation says that:

$$\frac{1}{2} m_i C_s^2 = \frac{1}{2} m_i u_i^2(-\infty) + e\phi_p \approx \phi_p \Leftrightarrow \frac{e\phi_p}{k_B T_e} = \frac{1}{2} \quad (10)$$

In other words, plasma potential equals the half of electron temperature. Finally,  $n_{se} = n_i(0) = n_e(0) = n_p \cdot \exp(-e\phi_p/k_B T_e) \approx 0.61 \cdot n_p$ .

In all the above considerations, the wall is not initially put at a certain potential. At steady-state or at equilibrium, electron and ion flux equals each other: we say that the wall floats. When the wall is at the floating potential,  $\phi_W = \phi_{fl}$ , the net current on the wall is then zero. Let us calculate the potential difference between the floating wall and the plasma. Mathematically speaking, floating condition involves that  $|J_e| = |J_i| \Leftrightarrow e\langle n_e v_e \rangle = e\langle n_i v_i \rangle \Leftrightarrow en_i - s)u_i(s) = en_e(s)\langle |\mathbf{u}_e| \rangle/4$ , indeed, electrons are Maxwellian – the flux equals the random flux on a surface – and ions are directed by the Bohm flux. Using the fact that ions flux is

<sup>2</sup>Remember that for all potential bellow plasma potential, electrons are considered as trapped by the potential well: they obey to a Maxwellian distribution. Their density is then given by the following formula:  $n_e(x) = n_A \cdot \exp[e(\phi(x) - \phi(x_A))/k_B T_e] \forall \phi \leq \phi_p$  and A a point between the wall and the bulk plasma.

conserved in the sheath, we have:

$$en_i(-s)u_i(-s) = en_{se}C_s = \frac{1}{4}en_p \exp \left[ e \frac{\phi_{fl} - \phi_p}{k_B T_e} \right] \sqrt{\frac{8k_B T_e}{\pi m_e}}, \quad (11)$$

or also,

$$\phi_p - \phi_{fl} = \frac{k_B T_e}{2e} \left( \ln \left[ \frac{m_i}{2\pi m_e} \right] + 1 \right). \quad (12)$$

### 1.2.2 The Child-Langmuir Law

From now on we drive our wall to a certain potential  $V$ . That causes the sheath structure to change, because the new potential drop oblige the plasma to widen or to shorten the sheath to achieve the shielding of the wall. In this section, we are going to calculate the value of  $s$ , the sheath length. With Poisson equation,

$$\frac{d^2 \phi}{dx^2} = -\frac{e}{\varepsilon_0} (n_i(x) - n_e(x)), \quad (13)$$

we won't get an analytical solution for  $s$ . We counter that problem with an approximation by assuming that the wall potential is very negative, so that  $e(V - \phi_p)/k_B T_e \ll 1$ . With that condition, electrons will be strongly repealed by the potential drop, so we omit their presence in the sheath:

$$\frac{d^2 \phi}{dx^2} = -\frac{e}{\varepsilon_0} (n_i(x) - n_e(x)) \sim -\frac{en_i(x)}{\varepsilon_0}. \quad (14)$$

Ion flux conservation in the sheath remains true and yields to:

$$en_i(x)u_i(x) = en_{se}C_s = J_i^0 \Leftrightarrow en_i(x) = \frac{J_i^0}{u_i(x)}. \quad (15)$$

Finally, ion energy conservation for strong potential drop between a point in the sheath and the sheath edge gives:

$$\frac{1}{2}m_i u_i^2(x) + e\phi(x) = \frac{1}{2}m_i C_s^2 \Leftrightarrow u_i(x) = \sqrt{-\frac{2e\phi(x)}{m_i}} \sqrt{1 - \frac{m_i C_s^2}{2e\phi(x)}} \sim \sqrt{-\frac{2e\phi(x)}{m_i}}. \quad (16)$$

Thus, the Poisson equation becomes:

$$\frac{d^2 \phi}{dx^2} = -\frac{J_i^0}{\varepsilon_0} \left( -\frac{2e\phi(x)}{m_i} \right)^{-1/2} = -\frac{J_i^0}{\varepsilon_0} \sqrt{\frac{m_i}{2e}} (-\phi(x))^{-1/2} \quad (17)$$

**Exercise.** Solve that integral. Tip: first multiply (17) with the first derivative of the potential. Integrate this new equation between sheath edge ( $x = 0$ ) and somewhere in the sheath ( $x$ ). Then, write the first derivative of  $\phi$  and solve.

**Solution.** The way of integrating this differential equation must be known because it can helps you in a lot of cases. After multiplying (17) we get:

$$\int_{x=0}^x \frac{d\phi}{dx} \cdot \frac{d^2\phi}{dx^2} dx = -\frac{J_i^0}{\varepsilon_0} \sqrt{\frac{m_i}{2e}} \int_{x=0}^x (-\phi(x))^{-1/2} \frac{d\phi}{dx} dx \quad (18)$$

The l.h.s. integrates straightforwardly and the r.h.s. becomes an integration over  $\phi$ , from  $\phi(0) = 0$  to  $\phi(x)$  (remember that  $\phi'(x=0) = 0$ ):

$$\begin{aligned} \frac{1}{2} \left[ \left( \frac{d\phi}{dx} \right)^2 \right]_0^x &= -\frac{J_i^0}{\varepsilon_0} \sqrt{\frac{m_i}{2e}} \int_0^\phi \frac{d\phi}{\sqrt{-\phi}} = \frac{2J_i^0}{\varepsilon_0} \sqrt{\frac{m_i}{2e}} \left[ \sqrt{-\phi} \right]_0^\phi \\ \Leftrightarrow \frac{1}{2} \left( \frac{d\phi}{dx} \right)^2 &= \frac{2J_i^0}{\varepsilon_0} \sqrt{\frac{m_i}{2e}} \sqrt{-\phi(x)} \\ \Leftrightarrow \frac{d\phi}{dx} &= -\sqrt{\frac{4J_i^0}{\varepsilon_0} \sqrt{\frac{m_i}{2e}}} (-\phi(x))^{1/4} \\ \Leftrightarrow \int_0^V -\frac{d\phi}{(-\phi)^{1/4}} &= \sqrt{\frac{4J_i^0}{\varepsilon_0} \sqrt{\frac{m_i}{2e}}} \int_0^s dx \\ \Leftrightarrow \frac{4}{3} (-V)^{3/4} &= s \sqrt{\frac{4J_i^0}{\varepsilon_0} \sqrt{\frac{m_i}{2e}}} \\ \Leftrightarrow J_i^0 &= \frac{4\varepsilon_0}{9} \sqrt{\frac{2e}{m_i}} \times \frac{|V|^{3/2}}{s^2} \end{aligned} \quad (19)$$

Which is the Child-Langmuir law. It shows that the ion current is proportional to the potential drop at a power of 3/2. After some algebra and rearrangements, the sheath length writes:

$$s = \frac{\sqrt{2}}{3} \lambda_{De} \left( 2 \frac{e|V|}{k_B T_e} \right)^{3/4} \quad \text{with electron Debye length } \lambda_{De} = \sqrt{\frac{\varepsilon_0 k_B T_e}{n_p e^2}} \quad (20)$$

The whole calculation is also true when considering a strong and positive wall potential: there will be no ions in the sheath, and the electron current will obey to equation (19) (of course, after changing  $m_i$  to  $m_e$ ).

### 1.2.3 The theory of Langmuir probes

Last section we put a voltage on our wall. From now on, we put a name to the wall and call it a probe. A Langmuir probe is nothing but a metal tip that we put into the plasma. It can have several shapes, planar, cylindrical, spherical,... Usually the probe potential swipes a range from -70 to 70 volts and the measured current is plot. The output  $I(V)$  curve is called the characteristics of the probe. This curve gives access to several plasma parameters, such as plasma potential, density and temperature.

The probe characteristics can be analytically calculated by computing the collected cur-

rent,

$$I(V) = I_e - I_i = \left\langle \iint_{\text{probe}} e(\Gamma_e - \Gamma_i) \cdot d\mathbf{S} \right\rangle, \quad (21)$$

note that here, the ion current is counted as negative, and the electron current is counted as positive (it is just a question of definition).

To do the derivation of  $I(V)$ , we must separate the two cases,  $V > \phi_p$  and  $V < \phi_p$ . Let us start with the second condition: the probe repels electrons and accelerate ions because it has a potential below the plasma potential. For that calculation, we call the  $z$  direction, the direction perpendicular to the probe surface. An electron from the bulk plasma must have a minimum velocity  $v_z$  to reach the probe:

$$\frac{1}{2}m_e v_{z,0}^2 - e\phi_p = 0 - eV \Leftrightarrow v_{z,0} = \sqrt{\frac{2e}{m_e}(\phi_p - V)} \quad (22)$$

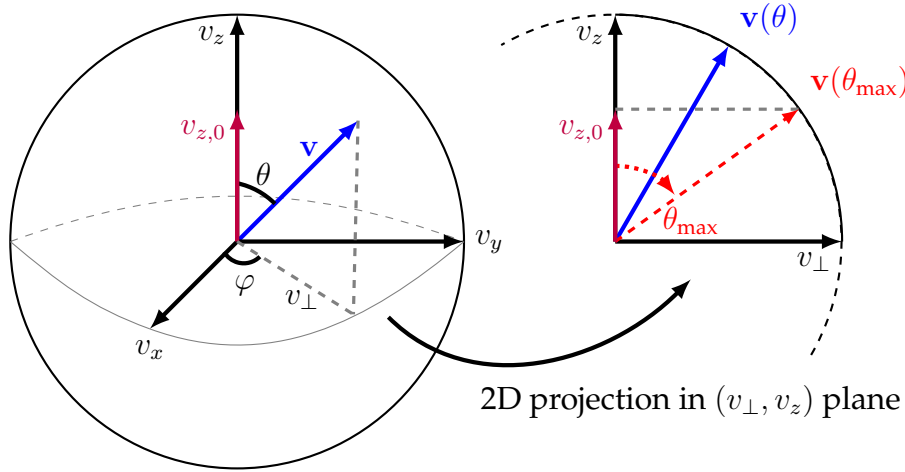


Figure 2: Visualisation of  $\theta_{\max}$  in velocity space. On the l.h.s. the sphere has a radius of  $v = |\mathbf{v}|$ , and we have represent a velocity vector in blue of length  $v$ . On the r.h.s. we show the projection at  $\varphi = \text{Cte}$ . As long as the angle  $\theta$  is below the value  $\theta_{\max}$  the electron of velocity  $v$  will touches the probe.

But as we know, in the plasma electrons have a random distribution of velocity, and the velocity vector has three component,  $\mathbf{v}$ . To ensure that an electron of velocity  $v = |\mathbf{v}|$  have a  $z$  component greater than  $v_{z,0}$  it must verifies (using the standard spherical coordinates) that:

$$\theta \leq \theta_{\max} = \arccos\left(\frac{v_{z,0}}{v}\right) \quad (23)$$

Let us now calculate the electron flux that verifies that limiting condition. For recall, the flux is estimated by averaging over the velocity space the quantity  $\Gamma = \langle n\mathbf{v} \rangle$ . In our case the only flux that will contribute to the current collection is its  $z$  component. Thus,

$$\Gamma_z = \iiint f_e(\mathbf{v}) v_z dv d\theta d\varphi \quad (24)$$

Using the fact that  $v_z = v \cos \theta$ , that  $v \in [v_{z,0}, +\infty[$ ,  $\theta \in [0, \theta_{\max}]$ ,  $\varphi \in [0, 2\pi]$ , and knowing that the electron velocity distribution function (which we assume to be Maxwellian) writes,

$$f_e(\mathbf{v}) d^3\mathbf{v} = n_p v^2 \left( \frac{m_e}{2\pi k_B T_e} \right)^{3/2} \cdot \exp \left[ -\frac{m_e v^2}{2k_B T_e} \right] dv \sin \theta d\theta d\varphi \quad (25)$$

And now place to the algebra:

$$\begin{aligned} \Gamma_z &= 2\pi n_p \left( \frac{m_e}{2\pi k_B T_e} \right)^{3/2} \int_{v_{z,0}}^{+\infty} v^3 e^{-mv^2/2k_B T_e} \int_0^{\theta_{\max}} \sin \theta \cos \theta d\theta dv \\ &= 2\pi n_p \left( \frac{m_e}{2\pi k_B T_e} \right)^{3/2} \int_{v_{z,0}}^{+\infty} v^3 \left[ \frac{\cos^2 \theta}{2} \right]_{\theta_{\max}}^0 e^{-mv^2/2k_B T_e} dv \\ &= \pi n_p \left( \frac{m_e}{2\pi k_B T_e} \right)^{3/2} \int_{v_{z,0}}^{+\infty} v^3 \left( 1 - \frac{v_{z,0}^2}{v^2} \right) e^{-mv^2/2k_B T_e} dv \\ &= \pi n_p \left( \frac{m_e}{2\pi k_B T_e} \right)^{3/2} \times 2 \left( \frac{k_B T_e}{m_e} \right)^2 e^{-mv_{z,0}^2/2k_B T_e} \\ \Gamma_z &= \Gamma_{\text{sat.}} \cdot \exp \left[ e \frac{V - \phi_p}{k_B T_e} \right] \end{aligned} \quad (26)$$

where  $\Gamma_{\text{sat.}}$  is the random flux on a surface,

$$\Gamma_{\text{sat.}} = \frac{1}{4} n_p \langle v \rangle = \frac{1}{4} n_p \sqrt{\frac{8k_B T_e}{\pi m_e}} \quad (27)$$

It is also the electron flux on a surface when they are no more trapped (i.e.  $V \geq \phi_p$ ). Finally, since the electron current is nothing but  $I_e = e\Gamma_z S$  (where  $S$  is the collection surface), we have:

$$I_e(V) = \begin{cases} \frac{1}{4} e n_p S \sqrt{\frac{8k_B T_e}{\pi m_e}} \cdot \exp \left[ e \frac{V - \phi_p}{k_B T_e} \right] & , \quad V \leq \phi_p \\ \frac{1}{4} e n_p S \sqrt{\frac{8k_B T_e}{\pi m_e}} & , \quad V \geq \phi_p \end{cases} \quad (28)$$

The quantity  $e\Gamma_{\text{sat.}}S$  is also called the "electron saturation current". For ions it is more easier: in our very simple case, we supposed that ions are cold, so that  $T_i \rightarrow 0$ . Thus, when they are attracted, they have a Bohm flux towards the probe, and when they are repelled we approximate the ion current to 0:

$$I_i(V) = \begin{cases} 0,61 \times e n_p S \sqrt{\frac{k_B T_e}{m_i}} & , \quad V \leq \phi_p \\ 0 & , \quad V \geq \phi_p \end{cases} \quad (29)$$

The collected current on the probe is the sum of both contributions (ion plus electron).

### 1.2.4 In practice and discussions

In what follows, we consider a cylindrical Langmuir probe of length  $L_p = 1$  cm and of radius  $r_p = 75 \mu\text{m}$ . The measurement takes a few seconds, and the probe makes several curves and then outputs the average characteristics. For an He plasma at 1.2 Pa the IV characteristics is plotted in figure 3. As we can see there are 3 distinct regions: the ion collection region, the electron collection one, and the transition region. I recommend the lecture of Merlino's paper [3] which is very clear and comprehensible by M1 students.

Graphically it is straight forward to spot the floating potential: it is the point where the curve cross the  $I = 0$  line. To find the plasma potential, one must remember that at the plasma potential, the electron current switches from an exponential behaviour to a saturation behaviour. The plasma potential is defined as the potential at which:

$$\left. \frac{dI}{dV} \right|_{V=\phi_p} = \max \left( \frac{dI}{dV} \right) \quad (30)$$

This criterion is quite simple, but experimentally there is a lot of incertitude in the determination of the plasma potential, as we can see in the figure 3-left.

Moreover, we notice that the electron- and ion- saturation currents do not saturate. This is due to the fact that the surface of collection  $S$  is not the probe surface: we need to take into account the sheath extension given by eq.(20)<sup>3</sup>! This adds a lot of problem to determine  $n_p$  and  $T_e$ , and there is no method that works every time. One method consist of using a self-consistent code that uses the transition region. First we assume that  $T_e$  equals a certain temperature. Second we calculate  $n_p$  using the current at plasma potential (which uses the input temperature). Then we calculate the effective collection area using the formula (20) and fit the ion current. Finally we compute the first derivative of  $I(V) - I_i(V)$  (because this quantity is the exponential electronic current) to get a new  $T_e$ . We start again until the values of  $n_p$  and  $T_e$  converges.

But in the experimental work you will use a software that calculates all plasma parameters by its own. So do not worry about computational problems, just keep in mind that it is not straight forward at all!

### 1.2.5 A bit further

**Energy distribution function.** Starting from the third line of calculation (26), we can introduce the electron kinetic energy expressed in volts as  $\mathcal{E} = mv^2/2e$ . Remembering that the energy distribution function for electrons is given by,

$$g_e(\mathcal{E}) = 2\pi \left( \frac{2e}{m_e} \right)^{3/2} \sqrt{\mathcal{E}} f_e(\mathbf{v}(\mathcal{E})) \quad (31)$$

---

<sup>3</sup>Note that you have to change  $V$  from equation (20) to  $V - \phi_p$ . The potential at which the collection surface is the probe surface is the plasma potential.

and using this integration rule,

$$\Psi = \int_a^b \psi(a, x) dx \Leftrightarrow \frac{\partial \Psi}{\partial a} = \int_0^b \frac{\partial \psi}{\partial a} dx - \psi(a, x) \quad (32)$$

we get something interesting, since

$$g_e(\mathcal{E} = V) = \frac{2m_e}{Se^2} \sqrt{\frac{m_e}{2eV}} \cdot \frac{d^2 I}{dV^2} \quad (33)$$

But again, even with a clean IV curve, the derivative and the second derivative are noisy, and it can be problematic to extract some physics from the second derivative (in our case as plotted in figure 3–right it is even impossible).

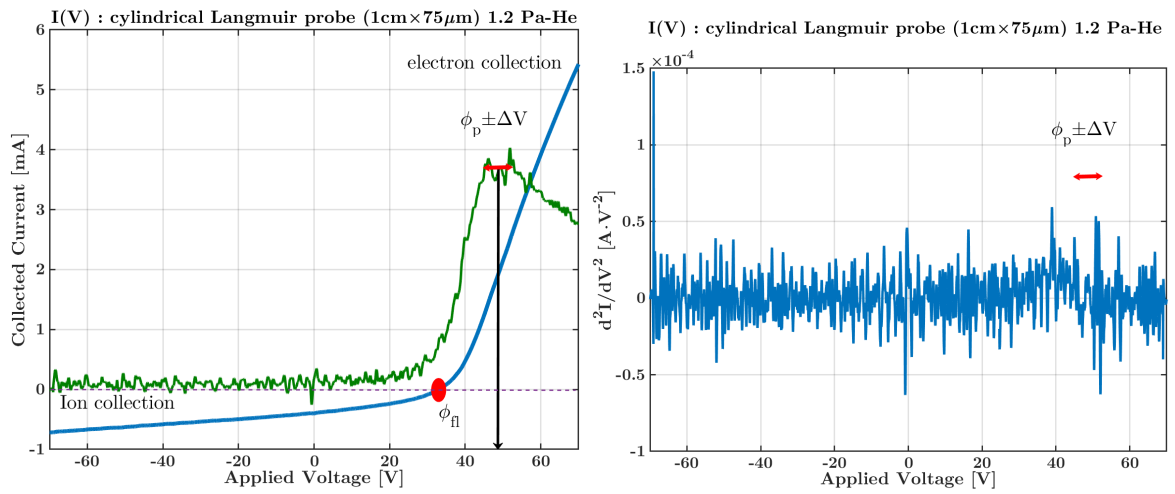


Figure 3: Left: Typical IV characteristics obtained on the ALINE experimental device (at null magnetic field). On blue the IV characteristics, and in green its first derivative – Right: Second derivative of the characteristic plotted left.

**In the presence of a strong magnetic field.** As you know, in the presence of a  $B$  field electrons have a so-called cyclotron motion around magnetic field lines. This causes the plasma to become an anisotropic medium. The consequences are that we do not know how to calculate the effective collection area, and the IV curve has a bump just after the plasma potential that we do not understand entirely. This bump may come from a plasma depletion [4] because the probe flush a lot of electrons in the magnetic flux tube.

## 2 Laser Induced fluorescence on the magnetron device

### 2.1 Experimental device

#### 2.1.1 Introduction

I recommend you to have a look on the first chapter of my Master's Thesis [5] and on the Thesis of M. Desecures [6] for large explanation on the subject. And anyway, during the manipulation day you will have the presentation with the device in front of you.

However, the reactor was designed to study the physical vapour deposition (PVD) by magnetron sputtering. To be simple: we turn on a plasma around a piece of metal called "target". By biasing the target to high negative voltage, ions from the plasma will be accelerated toward its surface. Incoming ions have a high kinetic energy, and through a complex mechanism, they eject some metal atoms: it's the sputtering of the target. The sputtered atoms transport themselves through the "sputtering cone". So if we want to achieve deposition, one can place a substrate in the front of the target so that sputtered atoms can deposit: it's the PVD. PVD has a lot of industrial applications<sup>4</sup>.

#### 2.1.2 The magnetron

Here we see that the sputter atom have to move from the target to the substrate. To have a good deposition rate, the pressure must be reduced to minimize the loss due to collisions between metallic atoms and the gas. But according to Paschen's law, reducing the pressure oblige us to increase the discharge voltage. At our pressure (about 10 mTorr) the ignition voltage would be at the order of several kilovolts! To counter that problem, we place a circular magnet, called magnetron, behind the target. Its magnetic field confines the electrons by increasing their mean free path and allows us to decrease the voltage around 100 volts. The plasma is denser at the positions where the drift velocity is higher, i.e. when  $\mathbf{E} \perp \mathbf{B}$  since,

$$\mathbf{v}_{\text{drift}} = \frac{\mathbf{E} \times \mathbf{B}}{B^2} \quad (34)$$

This inhomogeneity is visible on the target with the apparition of a racetrack as seen in figure 4. Ions from dense plasma are accelerated to the target, due to the electric field, and achieve its sputtering.

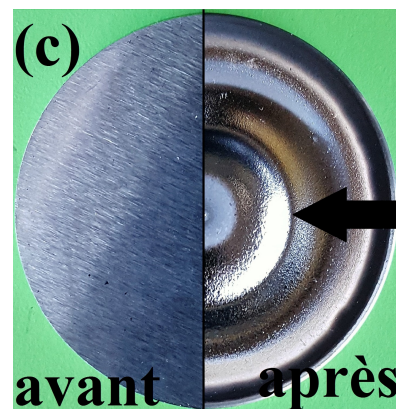


Figure 4: The same Tungstene (W) target before and after 10 hours of magnetron sputtering. The racetrack is pointed by the big fat and black arrow.

<sup>4</sup>This phrase doesn't sound convincing...

### 2.1.3 Experimental apparatus

We place a tunable diode (with spectral resolution of less than 0.005 pm) in the front of the probe and aim the center of the racetrack. About 20 percent of the laser intensity is redirected to a Fabry-Perot to measure the wavelength. The laser's trajectory goes along with the  $z$  axis (with  $e_z$  perpendicular to the target). Photons interact with the neutral metallic vapour. An optical fibre is placed perpendicular to the laser beam and is relied to an photomultiplier to convert the collected light intensity into voltage, and that voltage is plot on an oscilloscope.

## 2.2 The theory of Laser Induced Fluorescence

### 2.2.1 Using Doppler effect

LIF measurements are very complex to achieve, but the theory is quite intuitive. The tunable diode (TD) makes a periodic swipe of wavelength of few pm between  $\lambda_0$  and  $\lambda_0 + \delta\lambda$  (see fig.5). The range of wavelength is chosen so that it contains  $\lambda_{tr.}$ , the wavelength of transition from fundamental to a metastable state of the metallic atom specie. After absorption of the photon, the metastable atom releases almost instantly the photon and go back to it's fundamental. This photon is caught by the optical fibre. The conclusion of that, is that the oscilloscope displays the number of desexcitation over time.

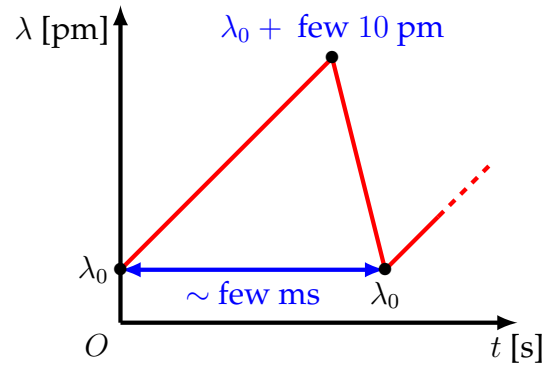


Figure 5: Laser wavelength over time.

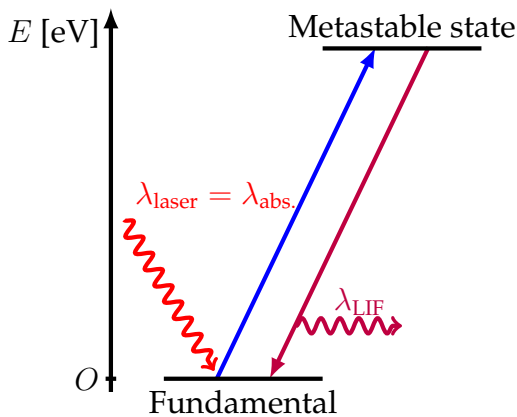


Figure 6: Scheme of the transition of a metallic atom.

But why shall we swipe a range of wavelength? A sputtered atom has a velocity  $v$ , with  $v_z$  the velocity perpendicular to the target (it is also the projection of  $v$  on the laser beam). When the atom catches a photon of wavelength  $\lambda_{laser}$  it actually "sees" another wavelength due to Doppler shift (because the atom is not at rest). So the absorption of the photon is actually done at a wavelength  $\lambda_{abs.} = \lambda_{tr.} + \Delta\lambda$ , where  $\Delta\lambda$  is the Doppler shift:

$$\Delta\lambda = \lambda_{laser} \cdot \frac{v_z}{c} \quad (35)$$

On the oscilloscope, we have the voltage  $V(t)$  that is proportional to the number of desexcitation for each time  $t$ . But thanks to Fraby-Perot, we know the relationship between time and wavelength,  $t \Leftrightarrow \lambda$ . Finally, the Doppler formula gives the velocity  $v_z$  for each  $\lambda$ . Thus, the oscilloscope displays  $V(v_z)$ : the number of metallic atom of velocity  $v_z$ . The LIF

signal is actually proportional<sup>5</sup> to the velocity distribution function of the metallic sputtered vapour!

### 2.2.2 Understanding and exploitation of a LIF curve

Now we know that the LIF measurements give us access to the velocity distribution function of the sputtered metallic atoms. But we do not now at which time  $t$  the probed velocity is  $v_z = 0$  (i.e.  $\lambda_{\text{abs.}} = \lambda_{\text{tr.}}$ ). In other terms we have the scale but not the origin. To solve that problem, let's take a look on a LIF curve, see fig. 7. The LIF profile is the convolution between 2 curves: a Gaussian and something else. In terms of distributions, a Gaussian corresponds to a "thermalized population", that means that those atoms are at equilibrium with the neutral gas because they had a lot of collisions (about 10). The mean velocity of those sputtered atoms is then 0: the maximum of the Gaussian defines the origin of our  $v_z$  axis. Hence, at Gaussian maximum  $v_z = 0 = \Delta\lambda$ .

The other distribution describes those atoms which did not had much collisions. The shape of the curve is called a modified-Thompson or a Stepanova; it is the distribution of energetic sputtered atoms at  $z = 0$ . For other distances  $z > 0$ , the Stepanova decreases in amplitude because energetic atoms do some collisions with the neutral gas, so they thermalize themselves and feed the thermalized population.

Mathematically, the distributions are given by:

$$f_{\text{TH}}(v_z) = N \sqrt{\frac{m}{2\pi k_B T}} \cdot \exp \left[ -\frac{mv_z^2}{2k_B T} \right] \quad (36)$$

for the thermalized population and by

$$f_{\text{EN}}(v_z) \propto \frac{\frac{1}{2}mv_z^2}{\left[ U_s + \frac{1}{2}mv_z^2 \right]^{-3+2b}} \left( 1 - \frac{U_s + \frac{1}{2}mv_z^2}{U_s + \frac{1}{2}mv_{z,\text{max}}^2} \right) \quad (37)$$

for the energetic population. Where  $m$ ,  $T$  and  $N$  are the mass, temperature and density of the sputtered vapour – and  $U_s$  is the metallic surface potential,  $b$  the binary interaction coefficient and finally  $v_{z,\text{max}}$  the maximum speed of the distribution.

Thanks to some numerical integrations, we are able to calculate the densities, mean velocity, mean energy, flux,... of each populations; which is very important if we want to

<sup>5</sup>The proportionality coefficient is obtained thanks to absorption spectroscopy.

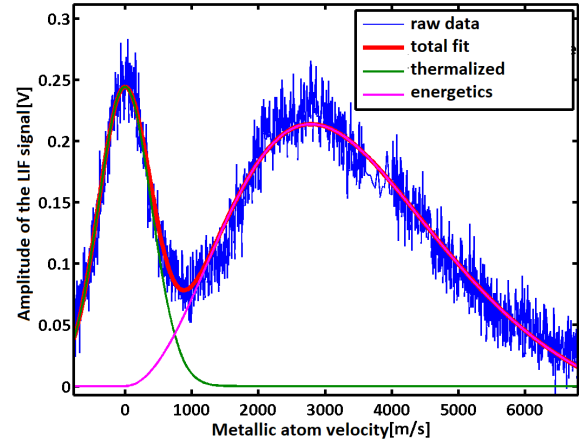


Figure 7: Example of the result of one LIF measurement and its fit curves.

characterize properly our PVD process (e.g. we can know exactly at which  $z$  the deposition will be optimal). Please, see appendix A (p.41) of ref. [5] for formulas. Again, during the experimental work, you will have access to a software that will provide you the numerical results. But you have to understand the meaning of the curve to give the good input parameters, and to discuss the results.

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