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Metastable-Helium-4 Source For Cold Atom Experiments

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Contents

Acknowledgements	iii
Introduction	1
Chapter 1. Theory of free-jet sources	3
1.1. Mean free path	3
1.2. Effusive beams	4
1.3. Supersonic beams	5
Chapter 2. Interaction between light and atom	11
2.1. Radiation pressure force	11
2.2. Deceleration of an atomic beam	16
2.3. Optical molasses	18
Chapter 3. Experimental setup	23
3.1. Source chamber	23
3.2. Beam chamber	28
3.3. Laser setup	32
3.4. Collimation	36
Chapter 4. Results	39
4.1. Intensity measurements	39
4.2. Time-of-flight	42
4.3. Collimation	46
Chapter 5. Summary and Outlook	49
Appendix A. Helium atom	51
A.1. Spectroscopic properties of ^4He	51

A.2. De-excitation processes	51
Appendix B. Curriculum Vitae	55
Bibliography	59

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Introduction

At the beginning of the 20th century scientists were occupied with a description of quantum nature of matter. In particular interest was the behaviour of bosons (with an even multiple of $\hbar/2$ as total spin) and fermions (with an odd multiple of $\hbar/2$ as total spin) at energies near the quantum degeneracy. In the mid-twenties of the last century Bose and Einstein predicted [1, 2] a macroscopic population of the ground state by bosons. This phase transition occurs when the interatomic distance becomes comparable to the de Broglie wavelength of atoms and is known as a Bose-Einstein condensate (BEC).

The invention of a laser in 1960 [3] enabled a rapid development of controlled manipulation of neutral atoms with light. In 1975 Hänsch and Schawlow [4] proposed a laser cooling scheme for atoms, whose first realisation was done by Chu et al. in 1985 [5]. Further development of the cooling techniques allowed Wieman and Cornell with ^{87}Rb [6] along with Ketterle with ^{23}Na [7] to realise BEC in 1995. For this achievement all of them were honoured with a Nobel Prize in 2001. Since then, BEC has been observed in other species: ^1H [8], ^4He [9, 10], ^7Li [11], ^{40}Ca [12], ^{41}K [13], ^{52}Cr [14], ^{84}Sr [15], ^{85}Rb [16], ^{131}Cs [17], ^{170}Yb [18], ^{174}Yb [19] and ^{176}Yb [20]. Each of the atomic species needed a unique treatment in order to achieve BEC. Most interestingly, the helium condensate was the first BEC to have been realised in an excited, metastable, state.

This thesis presents a first step on the way towards a metastable helium BEC. Metastable helium has a low mass, a simple electronic level structure and 19.8 eV of internal energy. Due to this large internal energy, helium can be efficiently detected using a microchannel plate detector, which also provides position information of the particle. These features make the metastable helium a perfect candidate in studying fundamental quantum mechanics. The scope of this thesis was to build and characterise a metastable

helium source. This kind of experiment requires a high-flux source with a low peak velocity. In addition, the atomic beam was collimated and deflected, in order to increase its intensity.

The thesis is structured as follows:

- Chapter 1 describes the theory of flow in the supersonic beam. In addition, this chapter contains a comparison between Maxwell-Boltzmann, effusive and supersonic velocity distributions.
- The most important basics of the theory of interaction between light and a neutral atom are presented in Chapter 2. The model illustrates the situation of a photon impinging on an idealised two-level atom. This leads to the emergence of dissipative and conservative forces, which act on an atom interacting with a light field.
- In Chapter 3 the experimental setup is presented. In the first part a continuous discharge source is compared to electron impact and pulsed discharge sources, as well as a description of the vacuum chamber is given. The second part depicts the laser arrangement which is used for the collimation and characterisation of the atomic beam entering the next part of the experiment. This section expounds also the polarisation spectroscopy — a method which is used to stabilise the lasers' frequency in the experiment.
- The source was characterised in a series of experiments whose results are presented in Chapter 4. They consist of flux measurements of the metastable helium beam, time-of-flight measurements which enabled us to determine the peak velocity of atoms in the stream and the beam collimation outcome.
- At the end, Chapter 5 summarises the results of the thesis and presents an outlook of the experiment.

CHAPTER 1

Theory of free-jet sources

In this chapter a theory of supersonic molecular beams is presented. In the past few decades molecular beams have been an increasingly important tool in the atomic and molecular physics. Particularly the supersonic expansion is extremely useful in providing beams with a well-defined kinetic energy in a given direction. This in particular permits to study some exotic atomic or molecular species.

Let us assume that we have a reservoir with a small orifice of diameter D , containing a gas at a pressure p_0 and a temperature T_0 . The ambient pressure is p_b . The important quantities governing the behaviour of the gas in the box are [21]:

- Mean free path λ_0 ;
- Size of the orifice D ;
- Size of the box;
- Lengthscale of density fluctuations.

It is assumed that the box is much larger than λ_0 and that the density fluctuations are negligible. Therefore the expansion of the gas from the reservoir is influenced by λ_0 and D .

The expansion takes place in two physically distinct regimes: effusive, if $\lambda_0 \gg D$ and supersonic, if $\lambda_0 \ll D$. An equivalent and experimentally more feasible description [22] is done with $\alpha = p_0/p_b$ and $G = ((\gamma + 1)/2)^{\gamma/(\gamma-1)}$ (where $\gamma = C_p/C_V$ is the heat capacity ratio), which measures if the expansion is supersonic or not. In the limit $\alpha \leq G$ one works in the effusive regime and otherwise in the supersonic. In the case of helium ($\gamma = 5/3$) $G = 2.05$.

1.1. Mean free path

The average distance travelled by an atom between two successive collisions is called a mean free path. The mean free path λ_0 , is equal to the average speed $\bar{v}$ divided by the

collision frequency f

$$(1.1) \quad \lambda_0 = \frac{\bar{v}}{f}.$$

It can be shown [23] that the collision frequency is given by

$$(1.2) \quad f = \sqrt{2}n\pi d^2\bar{v},$$

where n is the number density of particles in the unity volume ($V = 1 \text{ m}^3$) and d the diameter of the particle. Taking into account that

$$(1.3) \quad n = \frac{p_0}{k_B T_0},$$

where k_B is Boltzmann's constant, then from Eqs. (1.1) and (1.2) it follows that

$$(1.4) \quad \lambda_0 = \frac{k_B T_0}{\sqrt{2}p_0\pi d^2}.$$

Typical mean free paths for ^4He in normal conditions and in a precooled reservoir are given in Table 1.3.

1.2. Effusive beams

As the characteristic parameters approach the limit $\lambda_0 \gg D$ the number of collisions experienced by the gas approaches zero and an effusive beam is generated. The beam leaks through a very small hole straight into the vacuum without any increase in velocity.

The velocity distribution of a gas confined in a volume is given by the Maxwell-Boltzmann distribution

$$(1.5) \quad P_M(v) \propto v^2 e^{-\frac{mv^2}{2k_B T}},$$

where m is the mass of the atom and k_B is Boltzmann's constant. The peak and average velocities are given by

$$(1.6) \quad v_{\text{pM}} = \sqrt{\frac{2k_B T}{m}} \quad \text{and} \quad \bar{v}_M = \int_0^\infty v P_M(v) dv = 1.13 \sqrt{\frac{2k_B T}{m}},$$

respectively. One might anticipate that Eq. (1.5) is also valid in the molecular beam. Notwithstanding, as it was proven empirically, this is not the case. The distribution of

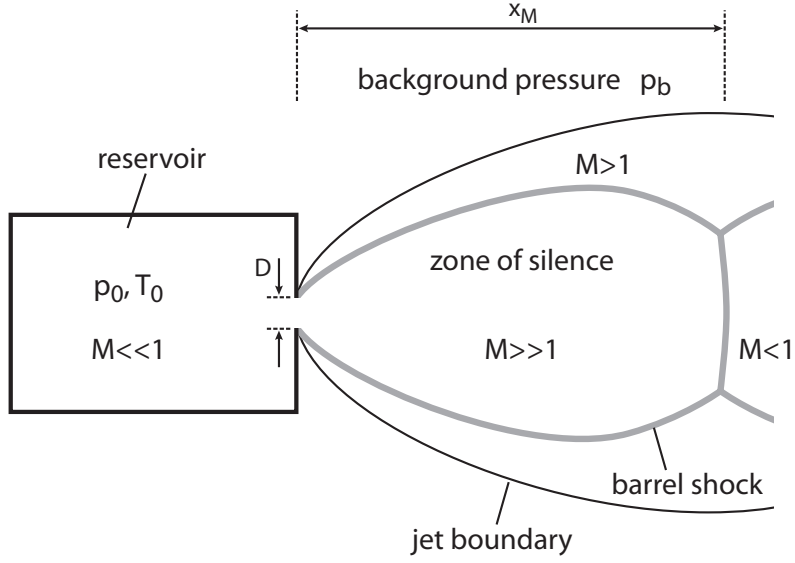


FIGURE 1.1. Depiction of a freely expanding jet from the reservoir which is held at pressure p_0 and temperature T_0 into a vacuum chamber maintained at the pressure p_b .

atomic velocities in the effusive beam can be derived and is given as [24]

$$(1.7) \quad P_E(v) \propto v^3 e^{-\frac{mv^2}{2k_B T_0}},$$

The most probable and average velocities are then equal to

$$(1.8) \quad v_{pE} = 1.22 \sqrt{\frac{2k_B T}{m}} \quad \text{and} \quad \bar{v}_E = 1.33 \sqrt{\frac{2k_B T}{m}}.$$

The normalised velocity distributions from Eqs. (1.5) and (1.7) are shown in Fig. 1.2.

1.3. Supersonic beams

In the limit of $\lambda_0 \ll D$ atoms undergo many collisions as they pass through the orifice. The expanding gas forms a supersonic beam whose expansion is shown in Fig. 1.1 and it can be described by a hydrodynamic-flow model. There are two distinct features of the supersonic beams [22]. First, unlike in effusive beams, the velocity of atoms increases after traversing the nozzle, which can be described by a Mach number $M = \bar{v}/c$ – where $\bar{v}$ is the mean gas velocity, and c the speed of sound in the medium, which for an ideal

gas is given by [24]

$$(1.9) \quad c = \sqrt{\frac{\gamma k_B T}{m}},$$

Second, there exist shock waves, as seen in Fig. 1.1, which result from the fact that the flowing atoms cannot sense the boundary conditions (hence the zone of silence) since this information propagates with the speed of sound whereas the atoms move faster than c . The zone of silence extends to [22]

$$(1.10) \quad x_M = 0.67D \sqrt{\frac{p_0}{p_b}},$$

where flow enclosed by the barrel shock has constant entropy (isentropic) and does not depend on the ambient pressure p_b , because the supersonic flow in this region is not aware of any of any external condition. The volume between the barrel shock and the jet boundary is nonisentropic and experiences very complicated nonlinear flow dynamics, including viscosity and heat conductance. Thus the favourable region from which one could extract the molecular beam is the zone of silence, where the nonlinear dynamics does not play crucial role in the molecular flow.

1.3.1. Idealised continuum model. The analysis presented in Refs. [22] and [24] neglects the effects of viscosity and of heat transfer in the gas. The flow in the supersonic beam may be regarded as an adiabatic and isentropic expansion. The adiabatic assumption leads to the conservation of energy in the system, which is written as

$$(1.11) \quad h_0 = h + \frac{1}{2}\mu\bar{v}^2 = \text{constant},$$

where h_0 is the enthalpy in the reservoir, h the enthalpy in the expanding free-jet and μ is the average molar weight. Due to the thermal equilibrium in the reservoir, the mean velocity of the atoms within the container is zero. One can notice that as the gas expands it cools and the velocity increases. For the ideal gas

$$(1.12) \quad dh = C_p dT,$$

then from Eq. (1.11) it follows

$$(1.13) \quad \bar{v}^2 = \frac{2}{\mu}(h_0 - h) = \frac{2}{\mu} \int_T^{T_0} C_p dT,$$

therefore

$$(1.14) \quad \bar{v} = \sqrt{\frac{2C_p(T_0 - T)}{\mu}}.$$

For an ideal monoatomic gas, $C_p = 5/2R$, and provided that the gas is cooled considerably, such that $T \ll T_0$, one obtains the terminal velocity

$$(1.15) \quad v_{\max} = \sqrt{\frac{5RT_0}{\mu}}.$$

Equation (1.15) yields a significant result, namely since $v_{\max}$ depends on the reciprocal of the average molecular weight of the beam, one can decelerate light species or accelerate heavy species by diluting it in heavy or light gas, respectively.

Adiabatic expansion from initial conditions (p_0, T_0, ρ_0) to final conditions (p_1, T_1, ρ_1) , where ρ_i ($i = 0, 1$) is the gas density, gives

$$(1.16) \quad \frac{T_1}{T_0} = \left(\frac{p_1}{p_0}\right)^{(\gamma-1)/\gamma}; \quad \frac{\rho_1}{\rho_0} = \left(\frac{p_1}{p_0}\right)^{1/\gamma}; \quad \frac{\rho_1}{\rho_0} = \left(\frac{T_1}{T_0}\right)^{1/(\gamma-1)}.$$

Equations (1.16) directly imply that for a large pressure difference the temperature of the gas will be significantly reduced. Assuming that C_p is constant and using Eqs. (1.9), (1.11) and (1.16) one can express the temperature, pressure and density as

$$(1.17) \quad \frac{T(x)}{T_0} = W^{-1}; \quad \frac{p(x)}{p_0} = W^{-\gamma/(\gamma-1)}; \quad \frac{\rho(x)}{\rho_0} = W^{-1/(\gamma-1)},$$

where

$$(1.18) \quad W = 1 + \frac{\gamma-1}{2}M(x)^2,$$

It has been shown that the Mach number on the axis of a supersonic beam can be calculated from the formula [22]

$$(1.19) \quad M(x) = A \left(\frac{x-x_0}{D}\right)^{\gamma-1} - \frac{\frac{1}{2} \left(\frac{\gamma+1}{\gamma-1}\right)}{A \left(\frac{x-x_0}{D}\right)^{\gamma-1}}, \quad \text{for} \quad \left(\frac{x}{D}\right) > \left(\frac{x}{D}\right)_{\min},$$

where D is the orifice diameter and parameters A , x_0/D and $(x/D)_{\min}$ depend on γ and are given in Table 1.1. Table 1.2 presents calculated Mach numbers along the axis of the expanding supersonic ^4He beam ($\gamma = 5/3$) versus different orifice diameters.

TABLE 1.1. Necessary parameters for calculation of the centerline Mach number for three given values of γ .

γ	A	x_0/D	$(x/D)_{\min}$
1.67	3.26	0.075	2.5
1.4	3.65	0.40	6
1.29	3.96	0.85	4

TABLE 1.2. Calculated Mach numbers on the axis of the expanding supersonic ^4He beam for $x = x_{\min}$.

D [mm]	$x_{\min}$ [mm]	M
0.3	0.75	1.7
0.5	1.25	3.1
0.75	1.875	4.4

The velocity distribution of particles in the supersonic beam is given as

$$(1.20) \quad P_S(v) \propto v^3 e^{-\frac{m(v-\bar{v})^2}{2k_B T}},$$

where T and $\bar{v}$ are given by Eqs. (1.17) and (1.14), respectively. Figure 1.2 shows the normalised velocity distributions: the Maxwell-Boltzmann distribution given by the formula (1.5), the distribution for the effusive beam (see Eq. (1.7)) and the one for the supersonic free-jet (see Eq. (1.20)) for three different values of the Mach number. The velocity spreads are plotted for $T_0 = 300$ K. An apparent inference from Fig. 1.2 is that the supersonic expansion not only increases the peak velocity, but also narrows the velocity spread as the Mach number increases. This is an advantage over the two other distributions if one wants to manipulate the molecular beam using velocity selective techniques, without sacrificing much of the intensity.

Table 1.3 shows the mean free path λ_0 , the terminal velocity $v_{\max}$, and the speed of sound c in the supersonic beam, calculated for ^4He in normal and precooled conditions.

TABLE 1.3. Characteristic parameters of the expanding free-jet calculated for ^4He , whose effective diameter is $d = 0.218 \text{ nm}$ [25], with reservoir kept in normal conditions and precooled.

T_0 [K]	p_0 [mbar]	λ_0 [μm]	v_{max} [m/s]	c [m/s]
300	1013	0.2	1765	1020
30	10	1.9	558	323

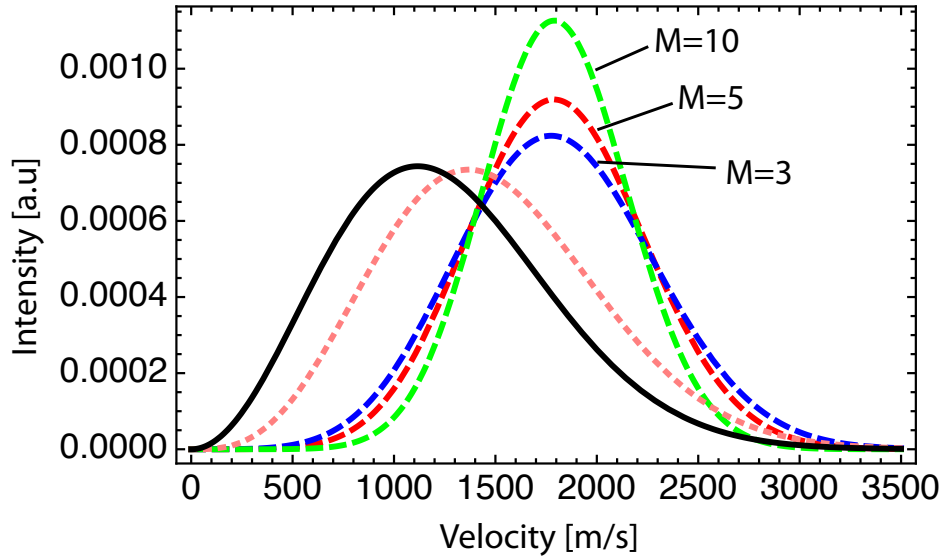


FIGURE 1.2. Comparison of the three normalised velocity distributions at room temperature $T_0 = 300 \text{ K}$: standard Maxwell-Boltzmann distribution (solid), a velocity distribution for an effusive beam (dotted) and the velocity distribution in the supersonic beam (dashed) for three different Mach numbers $M = 3, 5$ and 10 .

1.3.2. Atomic beam extraction. Some experiments may be conveniently performed directly in the free-jet zone of silence, but it is often advantageous to skim the supersonic free jet. This allows the supersonic beam to be collimated and transmitted to the next vacuum chamber which is operated at much lower pressure.

Skimming of the supersonic beam is done by placing a skimmer, with an orifice of diameter D_s , at a distance $x_s < x_M$. The shape of the skimmer is important (proved experimentally, cf. Sec. 4.1.2) as gas particles scatter off the wall and may collide with

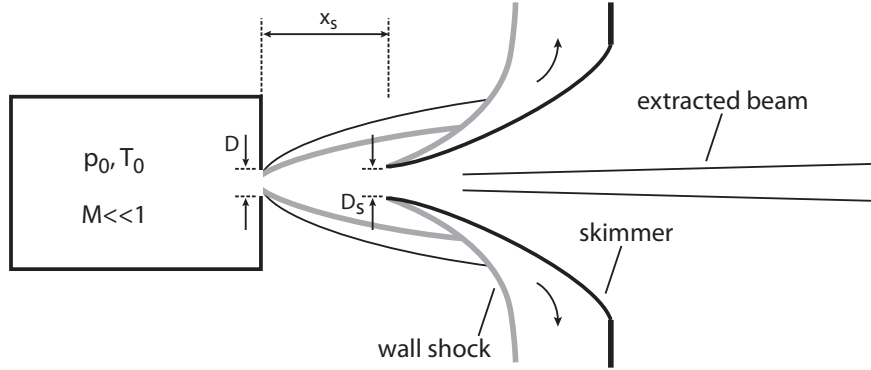


FIGURE 1.3. Schematic of the extraction of the supersonic beam. This is done by placing a conical-shaped skimmer in the zone of silence, which hampers interference between the extracted jet and the wall shock.

the atoms in the free jet. This inhibits the transfer efficiency of the beam into the second chamber without serious attenuation. Utilisation of a conically-shaped skimmer (see Fig. 1.3) allows positioning the orifice at some distance from the back wall of the first chamber, to hinder the unwanted collisions.

The extracted beam is much more collimated than the free jet, because the skimmer selects particles with small transversal velocities. Yet in order to prevent the beam from further transverse expansion one needs to employ laser collimation which is theoretically described in the following chapter.

CHAPTER 2

Interaction between light and atom

There are two fundamentally different forces acting on an atom moving in a radiation field. The first one is a radiation pressure force, which is a dissipative force as it transfers the energy from an atom+laser system to the reservoir and leads to a damping force which is used to cool atoms. On the other side, the second one is a dipole (or reactive) force which conserves the energy in an atom+laser system and arises from the light induced ac-Stark shift in the atomic levels. This force is responsible for trapping atoms in an optical dipole trap.

2.1. Radiation pressure force

The first experimental observation of resonant light transferring its momentum to an atom was done by Frisch [26] in 1933. It was achieved 16 years after Einstein wrote

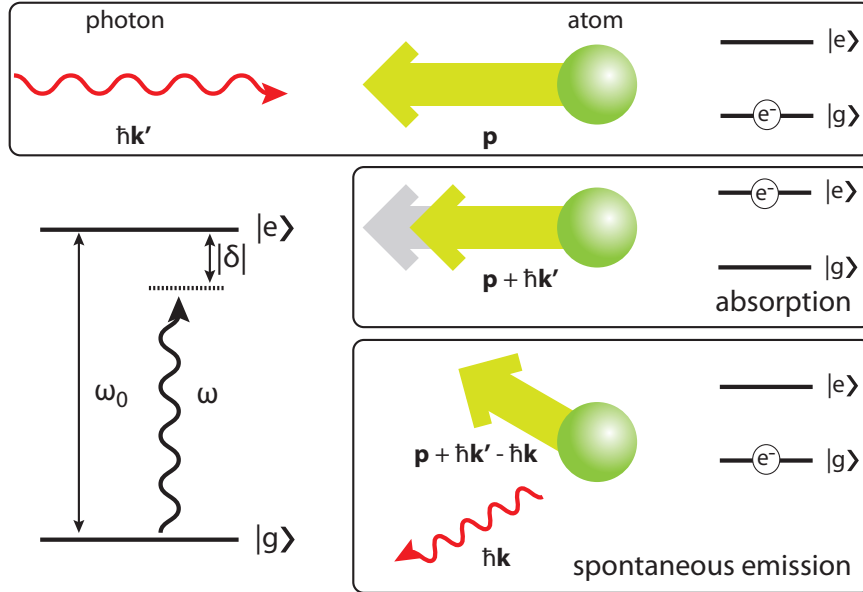


FIGURE 2.1. A scattering event of a photon with momentum $\hbar \mathbf{k}'$ and a frequency ω , which is detuned by δ with respect to an atomic transition frequency ω_0 .

a theoretical paper about absorption and emission possibilities of light quanta by an atom [27].

Let us consider a photon and a two-level atom carrying a momentum $\hbar\mathbf{k}'$ and $\mathbf{p}$, respectively, moving in opposite directions (see Fig. 2.1). The photon is detuned by δ from an atomic transition ω_0 . The atom, by absorbing the photon, obtains the momentum kick $\hbar\mathbf{k}'$. That is followed by a second momentum kick $\hbar\mathbf{k}$ resulting from a spontaneously emitted photon. The direction of the spontaneous emission is random, thus after N scattering events the momenta of the spontaneously emitted photons cancel out and the atom notices the effective momentum change from the impinging photons

$$(2.1) \quad \Delta\mathbf{p} = -\sum_N \hbar\mathbf{k} + \sum_N \hbar\mathbf{k}' = N\hbar\mathbf{k}'.$$

However, if the atom would emit the photon resulting from a stimulated emission, then the momentum change of the atom would be $\Delta\mathbf{p} = 0$, due to the fact that the stimulated photon is emitted in the same direction as the stimulating photon. This phenomenon allowed in the mid-seventies Hänsch and Schawlow to propose a laser cooling scheme of atoms [4].

2.1.1. The Rabi model. The Hamiltonian of an electron bound to an atom interacting with an electric field can be written as [28]

$$(2.2) \quad \hat{H} = \hat{H}_0 + \hat{H}_1,$$

where

$$(2.3) \quad \hat{H}_0 = \frac{\hat{\mathbf{p}}^2}{2m'} + V(r)$$

is the Hamiltonian of an electron in the absence of external fields, and

$$(2.4) \quad \hat{H}_1 = -\hat{\mathbf{d}} \cdot \mathbf{E}(t),$$

where $\mathbf{E}(t) = \mathbf{E}_0 \cos(\omega t)$, is the Hamiltonian of an electric field. We assume that the field is abruptly turned on at $t = 0$ and until then the atom is in the ground state $|\phi(0)\rangle = |g\rangle$. Furthermore, we consider a weak field, such that the population changes are very little,

and treat it as a perturbation. For $t > 0$, we expand the state vector $|\psi(t)\rangle$ as

$$(2.5) \quad |\psi(t)\rangle = c_g(t)e^{-\frac{i}{\hbar}E_g t}|g\rangle + c_e(t)e^{-\frac{i}{\hbar}E_e t}|e\rangle,$$

where c_g and c_e are time-dependent probability amplitudes ($|c_g|^2 + |c_e|^2 = 1$). We can write the time-dependent Schrödinger equation

$$(2.6) \quad i\hbar \frac{\partial |\psi(t)\rangle}{\partial t} = (\hat{H}_0 + \hat{H}_1)|\psi(t)\rangle.$$

Using Eqs. (2.5) and (2.6) and the fact that $\hat{H}_0|i\rangle = E_i|i\rangle$, where $i = \{e, g\}$, we end up in the set of coupled differential equations

$$(2.7a) \quad \dot{c}_g = -\frac{i\mathcal{V}}{\hbar}c_e \cos(\omega t)e^{-\frac{i}{\hbar}\omega_0 t},$$

$$(2.7b) \quad \dot{c}_e = -\frac{i\mathcal{V}}{\hbar}c_g \cos(\omega t)e^{\frac{i}{\hbar}\omega_0 t},$$

where $\mathcal{V} = \langle e| -\hat{\mathbf{d}} \cdot \mathbf{E}_0 |g\rangle$. Solving Eqs. (2.7) and making a rotating wave approximation¹ (RWA) leads to a solution

$$(2.8a) \quad c_g = \left(\cos \frac{\Omega_R t}{2} - i \frac{\delta}{\Omega_R} \sin \frac{\Omega_R t}{2} \right) e^{\frac{i}{2}\delta t},$$

$$(2.8b) \quad c_e = -i \frac{\mathcal{V}}{\hbar \Omega_R} \sin \frac{\Omega_R t}{2} e^{-\frac{i}{2}\delta t},$$

where $\delta = \omega - \omega_0$ is the detuning and $\Omega_R = \sqrt{\delta^2 + (\mathcal{V}/\hbar)^2}$ is the Rabi frequency. The Rabi frequency characterises the coupling strength between an atom and field. Its relevance can be seen in the probability of finding an atom in the excited state

$$(2.9) \quad P_e(t) = |c_e(t)|^2 = \left(\frac{\mathcal{V}}{\hbar \Omega_R} \right)^2 \sin^2 \frac{\Omega_R t}{2}.$$

The population of the excited state oscillates at the Rabi frequency Ω_R . In Fig. 2.2 the probability of finding an atom in the excited state (Eq. (2.9)) is plotted for three values of detuning. For $\delta \neq 0$ the oscillation frequency increases, but the probability amplitude decreases.

¹If the radiation frequency ω is close to the atomic transition ω_0 , then the term $\omega - \omega_0$ changes much rapidly and dominate $\omega + \omega_0$. Hence, one may neglect the latter term and make the so called rotating wave approximation.

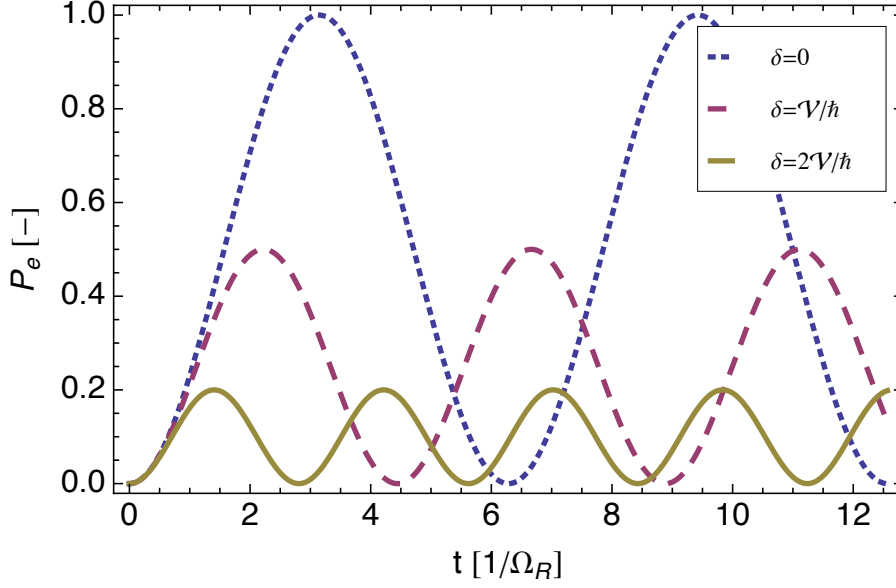


FIGURE 2.2. Rabi oscillations for three values of detuning. For $\delta \neq 0$ the oscillation frequency increases, but the probability amplitude decreases.

2.1.2. The Optical Bloch Equations. Further analysis of the evolution of populations including spontaneous emission requires employment of the density matrix formalism. (This derivation follows [29].) For a pure state

$$(2.10) \quad \rho = \begin{pmatrix} \rho_{ee} & \rho_{eg} \\ \rho_{ge} & \rho_{gg} \end{pmatrix} = \begin{pmatrix} c_e c_e^* & c_e c_g^* \\ c_g c_e^* & c_g c_g^* \end{pmatrix},$$

whose time evolution is described by the Liouville equation

$$(2.11) \quad i\hbar \frac{d\rho}{dt} = [\hat{H}, \rho].$$

Inserting Eqs. (2.8) and (2.10) into Eq. (2.11) leads to

$$(2.12a) \quad \dot{\rho}_{gg} = \Gamma \rho_{ee} + i \frac{\mathcal{V}}{2\hbar} (\tilde{\rho}_{eg} - \tilde{\rho}_{ge}),$$

$$(2.12b) \quad \dot{\rho}_{ee} = -\Gamma \rho_{ee} + i \frac{\mathcal{V}}{2\hbar} (\tilde{\rho}_{ge} - \tilde{\rho}_{eg}),$$

$$(2.12c) \quad \dot{\rho}_{ge} = -\left(\frac{\Gamma}{2} + i\delta\right) \tilde{\rho}_{ge} + i \frac{\mathcal{V}}{2\hbar} (\rho_{ee} - \rho_{gg}),$$

$$(2.12d) \quad \dot{\rho}_{eg} = -\left(\frac{\Gamma}{2} - i\delta\right) \tilde{\rho}_{eg} + i \frac{\mathcal{V}}{2\hbar} (\rho_{gg} - \rho_{ee}),$$

where all terms containing Γ are contributions of spontaneous emission, $\tilde{\rho}_{ge} = \rho_{ge} \exp(-i\delta t)$, $\Gamma = 1/\tau$ is the decay rate of the excited state and τ is its natural lifetime. Equations

(2.12) are called Optical Bloch Equations (OBE), in analogy to the Bloch equations for nuclear magnetic resonance. Since we are interested in the steady-state solutions of OBE, we set the time derivatives to zero. Knowing that population is conserved ($\rho_{ee} + \rho_{gg} = 1$), $\rho_{eg} = \rho_{ge}^*$ and introducing a new variable $w = \rho_{gg} - \rho_{ee}$, we obtain

$$(2.13a) \quad 0 = -\left(\frac{\Gamma}{2} - i\delta\right) \rho_{eg} + i\frac{w\mathcal{V}}{2\hbar},$$

$$(2.13b) \quad 0 = -\Gamma w - i\frac{\mathcal{V}}{\hbar}(\rho_{eg}^* - \rho_{eg}) + \Gamma.$$

Then the population of the excited state follows

$$(2.14) \quad \rho_{ee} = \frac{s}{2(1+s)},$$

where s is the saturation parameter given by

$$(2.15) \quad s = \frac{s_0}{1 + (2\delta/\Gamma)^2},$$

where $s_0 = I/I_s$ is the on-resonance saturation parameter with light field intensity I and saturation intensity I_s given by

$$(2.16) \quad I_s = \frac{\pi\hbar c}{3\lambda^3\tau}.$$

The force which light exerts on the atom is directly related to the scattering rate of impinging photons which one may write as

$$(2.17) \quad \gamma_p = \Gamma\rho_{ee} = \frac{s_0\Gamma/2}{1 + s_0 + (2\delta/\Gamma)^2}.$$

Figure 2.3 illustrates the scattering rate γ_p versus detuning of the laser light for a few values of the saturation parameter s_0 . For large values of s_0 , γ_p saturates at a value of 0.5Γ , since an increase in s_0 yields higher intensity of light, what raises the rate of stimulated emission, where $\Delta p = 0$. As we will see in the next paragraph, the saturation of γ_p puts a constraint on the maximum force which can be exerted on an atom. Hence the force is given by

$$(2.18) \quad F_{sp} = \gamma_p\hbar k = \frac{\hbar k s_0 \Gamma/2}{1 + s_0 + (2\delta/\Gamma)^2},$$

whose shape is Lorentzian as in Eq. (2.17). The force F_{sp} saturates at large values of s_0 , where it becomes $F_{\max} = \hbar k\Gamma/2$. Taking into account that in a two-level atom one cannot

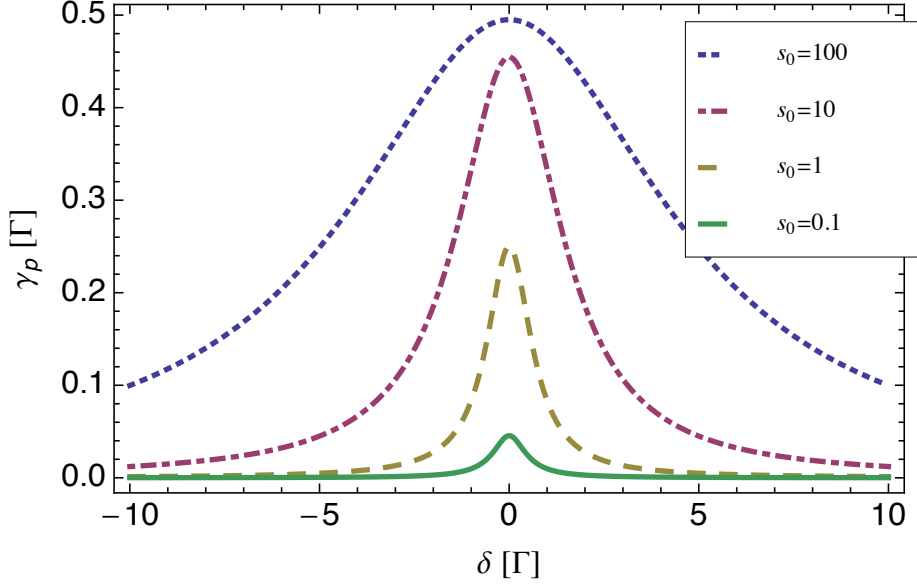


FIGURE 2.3. Scattering rate γ_p as a function of detuning δ for a few values of the saturation parameter s_0 . For high intensities, $s_0 \gg 1$, the linewidth becomes $\Gamma' = \Gamma\sqrt{1+s_0}$.

obtain the population of the excited state ρ_{ee} greater than $1/2$, the expression for $F_{\max}$ can be seen as intuitive. Equation (2.18) can only be applied to a two-level atom. The two-level atom can be realised only on a closed transition, that is if an atom evolves only between two states with no deexcitation to the external one. This is precisely what we have in metastable helium. If one pumps He^* to the 2^3P_2 state, then, due to the selection rules, it can decay only to the 2^3S_1 level.

2.2. Deceleration of an atomic beam

In order to slow down an atomic beam one has to utilise the dissipative force. The force given by Eq. (2.18) acting on a thermal atomic beam has to be corrected by the Doppler shifted laser frequency seen in the moving atoms' reference frame

$$(2.19) \quad \omega_D = -\mathbf{k} \cdot \mathbf{v},^2$$

²One has to note that atoms and light, moving in an opposite direction, always produce a positive Doppler shift.

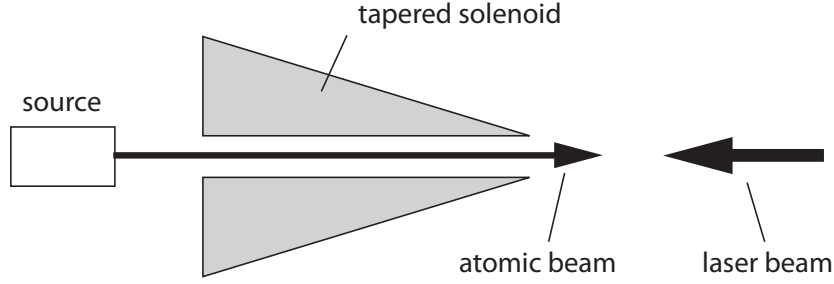


FIGURE 2.4. A schematic diagram of the Zeeman slower. A solenoid with windings inclining in a way that the energy shift given by Eq. (2.21) compensates for the decreasing Doppler shift of the light frequency.

leading to

$$(2.20) \quad F = \gamma_p \hbar k = \frac{\hbar k s_0 \Gamma / 2}{1 + s_0 + (2\delta_{\text{eff}}/\Gamma)^2},$$

where $\delta_{\text{eff}} = \delta + \omega_D$. For a fixed laser detuning and intensity a maximum deceleration³ requires $\delta + \omega_D \ll \Gamma$, so that the laser light is almost resonant in the atoms' frame of reference. One can infer from Eq. (2.19) that several scattering cycles suffice to significantly reduce the scattering rate, thus hindering the deceleration. Nonetheless, the Δv (recoil velocity) of a few m/s is still significantly less than the thermal velocity of atoms. For helium, one photon of 1083.33 nm wavelength (for spectroscopic data of helium atom see Sec. A.1) carries a momentum which can change the velocity of helium by $\Delta v = 9.2$ cm/s.

In order to slow down the atoms by a few hundreds of m/s, it is necessary to maintain $\delta + \omega_D \ll \Gamma$, what can be achieved by altering either δ or ω . The two most common experimental approaches of keeping the atoms in resonance are either sweeping the laser frequency or putting the atoms in a spatially varying dc magnetic field. The latter method (often being referred to as a Zeeman slowing technique) is commonly used in the atomic physics community, where it is often used to capture atom species in various traps. Figure 2.4 shows a schematic diagram of a Zeeman slower, in addition the next paragraph describes its principle of operation. The first experimental deceleration of an atomic beam, using a Zeeman slower, was done by Phillips and Metcalf [30]. It purely relies on the

³A remark should be made that deceleration of an atomic beam is not the same as cooling, which requires a compression of the velocity distribution in the phase space [29].

Zeeman effect to compensate the Doppler varying shift, which produces an energy shift

$$(2.21) \quad \Delta E_Z = -\vec{\mu} \cdot \mathbf{B},$$

where μ is the magnetic moment of a given state. In order to be able to simulate properly an incline of the required magnetic field, one needs to incorporate $\delta_Z = \mp \Delta_\mu B / \hbar$ in the δ_{eff} of Eq. 2.20, where $\mp$ is the sign for σ^+ and σ^- polarised light, respectively, and $\Delta_\mu = \mu^e - \mu^g$ is the difference between components of magnetic moments along the quantisation axis of the excited state and the ground state. The magnetic field, which is used to compensate for the decreasing Doppler shift, is created by a tapered solenoid (see Fig. 2.4). The decrease in the magnetic field is such that the slowed atoms, at each point along the solenoid, experience the maximal decelerating force F_{max} , hence under the assumption of constant force acting on atoms one expects the velocity to be changing as follows

$$(2.22) \quad v(z) = v_0 \sqrt{1 - \frac{2a_s z}{v_0^2}},$$

where v_0 is the initial velocity of atoms coming from the source, and $a_s = F_{max}/m$ is the deceleration. This leads to a required change in the magnetic field given by

$$(2.23) \quad B(z) = B_0 \sqrt{1 - \frac{2a_s z}{v_0^2}}$$

where $B_0 = \hbar k v_0 / \Delta_\mu$.

2.3. Optical molasses

An optical molasses (OM) is commonly used to transversely cool atomic beams, as well as to trap, with a help of a quadrupole magnetic field, atoms in a magneto-optical trap (MOT). The first experimental realisation of OM was done by Chu et al. [5]. The notion OM originates from the fact that an atom moving slowly along two beams, which are directed opposite to one another, experiences a net force proportional to its velocity, as shown in Fig. 2.5,

$$(2.24) \quad \mathbf{F}_{OM} \cong -\xi \mathbf{v}, \quad \text{where } \xi = \frac{8\hbar k^2 \delta s_0 / \Gamma}{(1 + s_0 + (2\delta/\Gamma)^2)^2}.$$

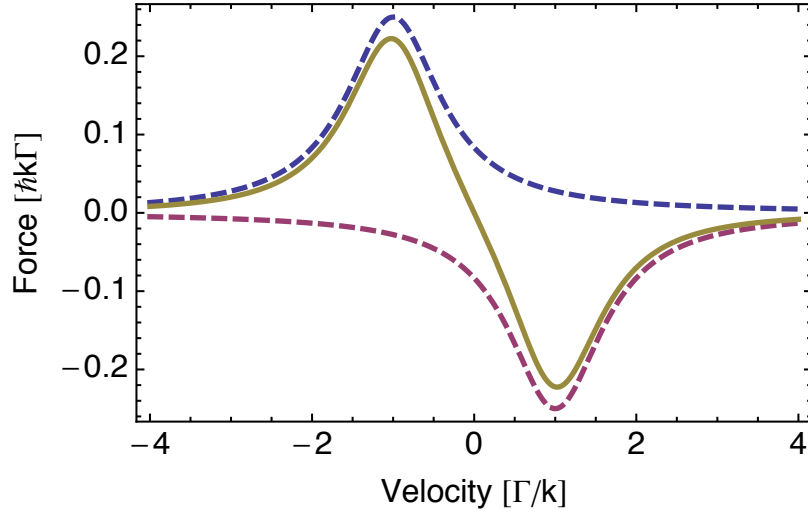


FIGURE 2.5. Force acting on an atom moving along two counterpropagating laser beams in 1D as a function of its velocity. The dashed lines show contribution of each beam and the solid one is the resultant force. Plotted for $s_0 = 10$ and $\delta = -\Gamma$.

For $\delta < 0$, this force opposes the velocity what leads to a viscous damping of the atomic motion. Equation (2.24) is valid only for a sufficiently low intensity of the laser beams ($I \ll I_0$), so that one could neglect stimulated emission. For an atom moving along the two beams, the damping force F_{OM} leads to a loss rate of kinetic energy of [31]

$$(2.25) \quad \left(\frac{dE}{dt} \right)_{\text{cool}} = \mathbf{F}_{OM} \cdot \mathbf{v} = -\xi v^2.$$

Classically one could think that the cooling process would lead atoms to achieve $v = 0$, yet due to the randomness of the absorption and the spontaneous emission of photons, one has to take into account a heating process of the sample. The F_{OM} given by Eq. (2.24) is only an average force, and its fluctuations produce heating. One can imagine an atom at rest, which is equally likely to absorb a photon from a wave travelling to the right or to the left. Absorption of a photon causes a random kick in momentum space, with a step size $\hbar k$. Another random step occurs after spontaneous emission of a photon. Thus the atom follows two random-walk steps in the absorption-emission cycle. Having a photon scattering rate γ_p one can calculate, for a truly one-dimensional case, the kinetic

energy increase [31]

$$(2.26) \quad \left(\frac{dE}{dt} \right)_{\text{heat}} = \frac{(\hbar k)^2 \gamma_p}{m}.$$

The competition between cooling and heating leads to a nonzero velocity in a steady state. In the equilibrium the cooling rate Eq. (2.25) is equal to the heating rate Eq.

(2.26)

$$(2.27) \quad \left(\frac{dE}{dt} \right)_{\text{cool}} + \left(\frac{dE}{dt} \right)_{\text{heat}} = 0,$$

this leads to a steady-state kinetic energy

$$(2.28) \quad \frac{mv_D^2}{2} = \frac{\hbar\Gamma}{8} \left(\frac{2|\delta|}{\Gamma} + \frac{\Gamma}{2|\delta|} \right).$$

This equation can be expressed in terms of temperature, taking into account that the minimum kinetic energy is obtained for $\delta = -\Gamma/2$ and $s_0 = 1$, as

$$(2.29) \quad T_D = \frac{\hbar\Gamma}{2k_B},$$

where T_D is the temperature limit of the so called Doppler cooling, which for helium $2^3S_1 \rightarrow 2^3P_2$ transition equals $T_D = 38.95$ mK, which corresponds to the velocity $v_D = 28.44$ cm/s.

OM in this work are used to increase the intensity of the atomic beam by collimation with curved wavefront. Utilisation of spherical waves instead of plane waves allows for capturing broader transverse velocities [32]. For the purpose of collimating a helium beam, which rapidly expands after passing through the skimmer, one needs to drive with the laser light on a closed transition between two helium sublevels. This can be realised with exciting the transition $2^3S_1 - 2^3P_2$ utilising σ^+ polarised light which induces transitions between the magnetic sublevels differing by $\Delta m_F = +1$ (see Fig. 2.6). After a few absorption-emission cycles the atom, following the selection rules, can decay only to the $(2^3S_1, m_F = +1)$ magnetic sublevel. Due to optical pumping the population oscillates between the $(2^3S_1, m_F = +1)$ and $(2^3P_2, m_F = +2)$ substates. And an atom experiences the radiation pressure. This optical pumping scheme is not possible with neither 2^3P_0 nor 2^3P_1 sublevel because of the fact that after a few cycles the atom will end up in a magnetic sublevel which cannot be excited by the σ^+ light.

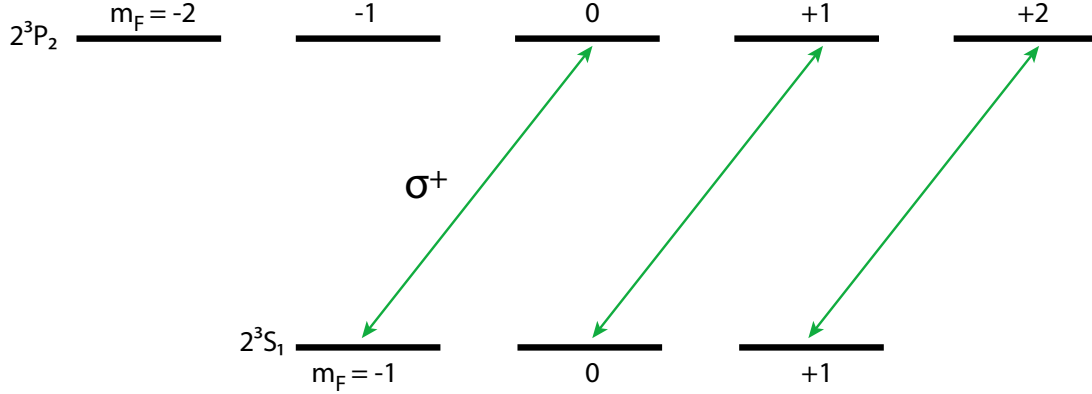


FIGURE 2.6. A $2^3S_1 - 2^3P_2$ transition in combination with σ^+ light, which excites only transitions with $\Delta m_F = +1$, resembles a two-level system. More precisely after a few absorption-emission cycles, atoms are cycling between $(2^3S_1, m_F = +1)$ and $(2^3P_2, m_F = +2)$.

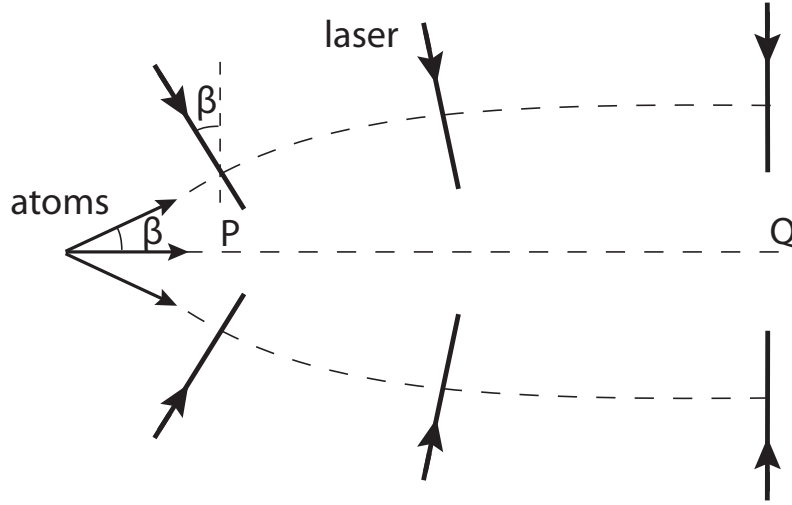


FIGURE 2.7. Collimation by curved wavefronts. Dashed lines are the atoms' trajectories, thick arrows is the laser light and β is a capture angle.

Figure 2.7 shows schematically collimation with curved wavefronts. For light launched at an angle $\beta = PQ/R$, where PQ is the length of the collimation region and R is the radius of the trajectory curvature, effectively only atoms moving perpendicularly to the laser beam follow its curvature. Therefore, in order to capture as much atoms as possible, one needs to have β as large as possible. However, β cannot be arbitrarily large, its maximal value is defined by the minimal radius of atoms' trajectory $R_{\min}$, since the

resonant radiation pressure force (Eq. (2.18)) must be equal or greater than the centrifugal force. It follows

$$(2.30) \quad \frac{mv^2}{R_{\min}} \leq \frac{\hbar k s_0 \Gamma / 2}{1 + s_0},$$

so for $s_0 = 1$, one obtains

$$(2.31) \quad R_{\min} \geq \frac{4mv^2}{\hbar k \Gamma}.$$

For example, for ^4He with velocity $v = 1300 \text{ m/s}$ interacting with $\lambda = 1083 \text{ nm}$ light the radius cannot be smaller than $R_{\min} = 8 \text{ m}$.

CHAPTER 3

Experimental setup

In this chapter an experimental setup, which was realised in our laboratory, is described in detail. The chapter consists of four main parts: a description of (1) the source chamber – a vacuum chamber where the metastable helium is produced; (2) the beam chamber – a vacuum chamber next to the source chamber where the helium beam is characterised and collimated; part (3) describes the laser setup and (4) depicts the laser collimation of atoms.

3.1. Source chamber

We want to study BEC of metastable helium atoms, therefore we need a source. There are a few approaches to realise a source of metastable atoms. Sources are based on various mechanisms: electron impact [33, 34, 35, 36, 37], pulsed discharge [38] and dc discharge sources [39, 40, 41, 42, 43]. In the Sections 3.1.1-3.1.3 these three kinds of sources are described.

We need a reliable source with high atom flux, narrow velocity distribution of helium metastable atoms with peak velocity as low as possible. Initial cooling of the atomic beam is vital for experiments with cold atoms, for the reason that its peak velocity at room temperature is typically in the range of 1700–2000 m/s. These velocities are way too high for slowing the atoms with a Zeeman slower of reasonable size. The slower the atomic beam coming from the source the shorter the slower one has to build, what considerably improves the loading rate into the magneto-optical trap (MOT).

3.1.1. Electron impact sources. In electron impact sources, an electron beam collides with a beam of neutral atoms and excites atoms to the metastable state. Since the cross section for metastable excitation is small ($\leq 10^{-17} \text{ cm}^2$ [44]), one needs a larger

overlap between the electron and atom beams, and for this reason a coaxial alignment of the two beams is preferred.

Early designs [33, 34] of electron impact sources resulted in a broad velocity distribution of the atomic beam due to the wide momentum spectrum of electrons. The population transfer is more noticeable for lighter species, such as helium. However, further developments (such as utilisation of a supersonic beam expansion and a pulsed discharge [37]) allowed for narrowing of the velocity distribution. A peak velocity of 2000 m/s was reported for a source operating at room temperature.

Such kind of sources are very complex in design and do not guarantee a high flux (typically $10^{10} - 10^{15} \text{ atoms s}^{-1} \text{ sr}^{-1}$ [33, 36, 37]). Furthermore no one has built this kind of source with precooled atoms, hence we would not know if we would get slow enough helium atoms.

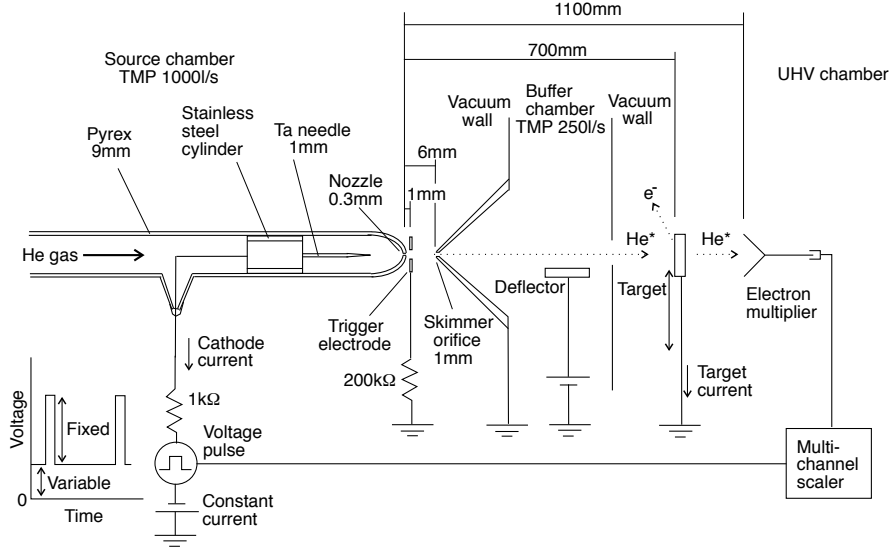


FIGURE 3.1. A schematic drawing of a pulsed discharge source used by Yamauchi et al. [38]. Metastable helium is produced in a discharge, which takes place between a tantalum needle and a nozzle.

3.1.2. Pulsed discharge sources. A pulsed discharge source is far less complicated than an electron impact source, yet sacrificing some of the intensity. Possible ways of achieving the pulsed operation are either usage of a pulsed nozzle [45] or application of a pulsed voltage to the anode [38]. A flux of $5.7 \times 10^{14} \text{ atoms s}^{-1} \text{ sr}^{-1}$ [38] was reported,

with a pulse duty of 0.05 ms in 2 ms. A schematic drawing of such a source is shown in Fig. 3.1. In principle, these sources do not differ too much from sources using a dc discharge, yet utilisation of the pulsed nozzle enables usage of the vacuum pumps with lower pumping speed without sacrifice of the intensity.

The discharge takes place between a sharp needle, which serves as an anode, and a grounded nozzle – a cathode. Atoms get ionised and accelerated towards the nozzle. Positively charged ions colliding with the cathode eject secondary electrons from the surface, which are accelerated towards the anode. Then they may collide with neutral atoms, and can ionise the atoms only if the electrons carried a sufficient amount of energy, otherwise the electrons can recombine with the atoms, which in turn become excited. If the secondary electrons yield at the cathode surface and the collisional probability of producing positive ions are sufficient, then the discharge will continue to “run” and current will flow. Otherwise, the discharge will “go out” and will not draw any current [46]. At this stage atoms in every possible state are produced, in addition, due to the high gas density between the needle and the nozzle many atoms undergo relaxation processes including Penning ionisation (see Section A.2). Thus most of the helium atoms in the metastable state are produced just after the nozzle, where the density is lower and atoms are subjected to the expansion in the free jet (see Sec. 1.3). The efficiency of this process¹ is estimated to be $\sim 10^{-5}$ which is still a low value, nonetheless much better compared to the electron impact sources².

3.1.3. DC discharge sources. There are a few types of dc discharge sources such as effusive hollow cathode sources [47, 43], hot cathode effusive sources [48] and supersonic cold cathode discharge sources [39, 40, 41, 42]. The flux magnitude is diverse $10^{12} - 10^{15} \text{ atoms s}^{-1} \text{ sr}^{-1}$, yet for the realisation of a helium BEC the intensity $\sim 10^{14} \text{ atoms s}^{-1} \text{ sr}^{-1}$ [41, 49] suffices.

¹A quotient of the number metastable atoms to the rest of the atomic beam.

²Although one should remember that in the electron impact sources, in the collinear design, one can make an arbitrarily long overlap between the atomic beam and the electron beam. So in the end the electron impact sources are more efficient.

The last type of DC discharge source was the choice of ours, since its design is straightforward, the flux is sufficient for creation of a helium BEC and the peak velocity of the helium beam is reported in the literature to be widely tunable from 300 m/s at the temperature of liquid helium, even though with a very low intensity of $10^{12} \text{ atoms s}^{-1} \text{ sr}^{-1}$, to 2000 m/s at room temperature. (Moreover one of our colleagues, Michael Keller, had an experience [50] of working with such a source at Stony Brook University, therefore we did not have to spend too much time on the design of the source.)

A schematic diagram of the source is depicted in Fig. 3.2. Our design is based on the one of Kawanaka et al. [40], who improved the design from Fahey et al. [39]. The novelty of Kawanaka's approach was application of liquid nitrogen to cool the gas.

The source is operated in a reverse flow mode. Helium enters the source vacuum chamber through a Pfeiffer Vacuum EVN 116 gas dosing valve and traverses between the glass tube and the steel tube until it reaches a teflon piece. There, due to the spiral grooving of the teflon spacer, the atoms' path is lengthened in order to improve the thermal contact with the copper head, which is attached to the cold finger cryostat (Vericold VT4-500). The temperature of the cryostat (without a load) at the lowest stage is about 4 K. Atoms enter the glass tube through an orifice where they undergo a process of excitation to the metastable state which is described in Sec. 3.1.2 (only with anode and cathode swapped). All atoms which do not get through to the source chamber flow backwards and get pumped out by a Pfeiffer Vacuum XtraDry 150-2 pump. The high voltage (typically 0.5 – 4 kV) is delivered by an Iseg HPn 60-506 power supply. It is applied to the tungsten needle³ through a 230 k Ω high power ballast resistor, which limits the current drawn by the discharge and makes the discharge itself more stable.

The source chamber is evacuated using a Pfeiffer Vacuum HiPace 400 turbomolecular pump with a pumping speed of 355 l/s (for N₂) and a forevacuum is created with a dry, oil-free piston pump Pfeiffer Vacuum XtraDry 150-2 with a pumping speed of 7.5 m³/h. Owing to the frequent access to the source and various selfmade components placed

³In a pressure region of a few mbar dielectric strength is very low thus to avoid a parasitic discharge at the other end of the source one needs to either insulate every metallic part or use KF glass vacuum components.

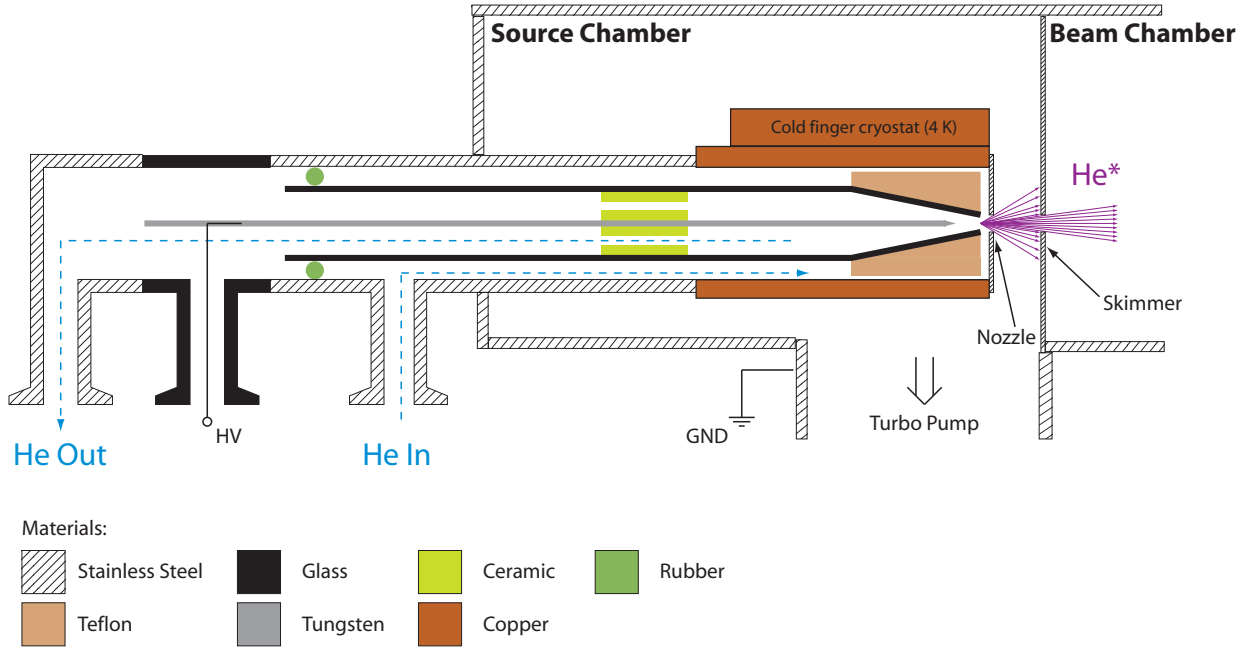


FIGURE 3.2. A schematic diagram (not to scale) of the source used in our experiment. It consists of a glass tube (10 mm in diameter, 1 mm thick wall and 1 mm orifice diameter), inside of which a 1 mm diameter tungsten needle is centred using ceramic spacers. The tube is surrounded by a stainless steel jacket, which outside the vacuum chamber is connected to KF glass vacuum elements and from inside it is attached to a copper head. The copper head is fixed to the lowest part of the cold finger cryostat. The needle is clamped to a translation stage (not shown in the drawing). A rubber O-ring is placed between the steel jacket and the glass tube, which prevents helium from flowing directly to the outlet. A teflon spacer, with a spiral grooving on outer walls, is fitted tightly in the copper part. The vacuum system is made of stainless steel and consists mostly of ConFlat (CF) flanges. Klein Flanges (KF) are used to connect helium in- and outlet points, high voltage feedthrough, as well as the glass and steel jacket.

inside the chamber, the source is only operated in a high vacuum (HV) regime. During the operation, the pressure in the source chamber reaches 10^{-3} mbar. Pressure is measured using a Stanford Research Systems (SRS) Pirani PG 105 gauge, which allows pressure measurements down to 10^{-4} mbar. The helium in- and outlet pressures are determined using a Pfeiffer Vacuum APR 260 piezo gauge (0.1 – 1100 mbar), which in contrast to the Pirani⁴ and ion gauges⁵ is independent of the gas type and therefore measures the absolute pressure.

Due to the connection between the head of the cryostat and the steel jacket, the initial (with helium inlet shut-off) temperature of the cryostat's head is 7 K. The temperature is measured using a Lakeshore CernOxTM CX-1050-Cu thermometer, and depends on the helium inlet pressure and the voltage applied to the needle.

3.2. Beam chamber

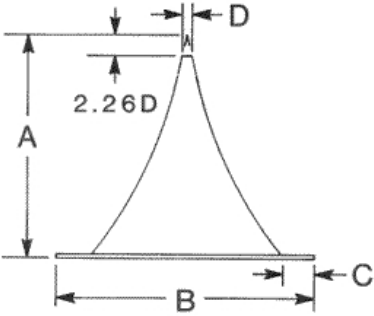
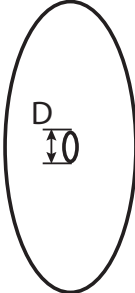
The beam chamber (see Fig. 3.2) is connected to the source chamber by a 0.5 mm diameter opening in the skimmer⁶, and is pumped out using the same pump setup as in the source chamber (see Sec. 3.1.3, turbomolecular pump Pfeiffer Vacuum HiPace 400 and the backing pump Pfeiffer Vacuum XtraDry 150-2.) In contrast to the source chamber, the pressure in the beam chamber reaches 10^{-8} mbar, and the vacuum is measured with a SRS NR-F-UHV (dual filament ThO₂/Ir) Bayard-Alpert ionisation gauge, which measures pressures down to $2 \cdot 10^{-11}$ mbar.

⁴Pirani gauge measures the thermal conductivity of the gas, thus the output is always calibrated to one kind of gas, normally N₂, therefore to get an actual pressure in the chamber, one needs to multiply the reading by a gas correction factor, which for helium is 1.1.

⁵An ion gauge is utilised in the beam chamber. It measures a current of ionised atoms, thus the pressure reading depends on the type of gas due to different sensitivities (cross section for ionisation at the given electron energy, number of molecules per unit volume etc.). The ion gauge is calibrated to N₂, hence the actual pressure for helium will be given after division of the pressure reading by 0.18.

⁶To optimise the metastable helium flux a few skimmers were tested, even though their orifice diameters were equal.

TABLE 3.1. A list of skimmers that are available in the experiment and their dimensions.

	Curved skimmers		Flat skimmer
			
Orifice diameter (D) [mm]		0.5	
Length (A) [mm]	6.99	19.1	NA
Wall thickness [μm]	50 – 80 (10 μm at the orifice)		500
Material	nickel		stainless steel

In the first part of the work, the beam chamber was employed to optimise the beam of metastable helium. The critical properties of the beam are its peak velocity and the intensity, which were measured by means of time-of-flight and flux measurements, respectively.

3.2.1. Flux determination. The flux is determined using a setup (see Fig. 3.3) consisting of a Faraday cup, which is placed about 10.8 cm behind the skimmer, and an ion deflection plate which is placed just in between the skimmer and the Faraday cup. All building elements, excluding homemade mask and aluminium cylinder, were stainless steel and ceramic eV Parts manufactured by Kimball Physics.

The metastable helium beam passes through a 5.5 mm diameter mask and impinges on the stainless steel plate releasing electrons with $\eta = 69\%$ probability [51] from the plate. These electrons are then collected with the cylinder in front of the plate to which a $+V_2 = +1000\text{ V}$ potential is applied. The current flowing from ground to the plate is determined using a Keithley 6485 picoammeter.

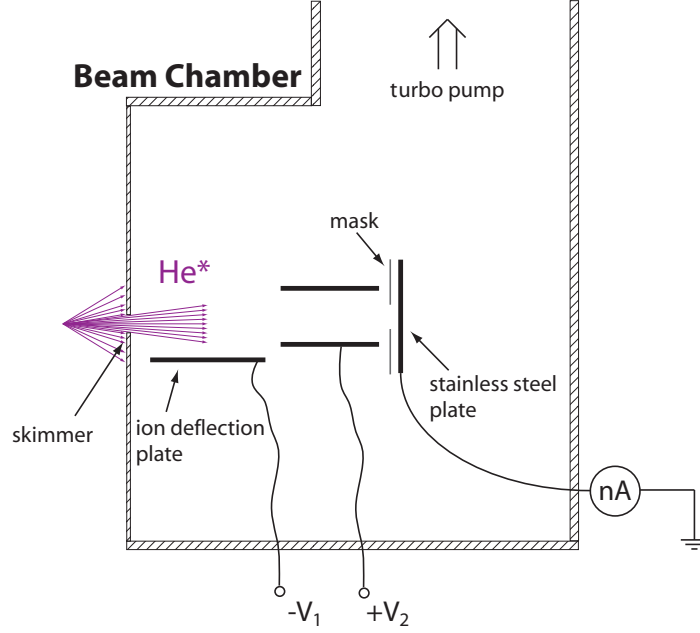


FIGURE 3.3. Experimental setup for the flux measurements. Atoms exiting the skimmer traverse a mask and hit the stainless steel plate, ejecting electrons. The electrons are collected by the cylinder which is connected to a $+V_2$ potential. The flux of atoms is determined with a picoammeter nA . All ions and free electrons are deflected by an ion deflection plate which is connected to $-V_1$ potential.

A voltage of $V_1 = -1500 \text{ V}$ is applied to the ion deflection plate, which diverts not only helium ions but also free electrons.

3.2.2. Time of flight. The uncollimated metastable helium beam exiting the skimmer is chopped with a mechanical chopper (see Fig. 3.4) – a round plate (10 cm in diameter), with two rectangular slits ($0.5 \times 2.7 \text{ mm}^2$) which is attached to a Maxon dc 2326.945-12.111-100 motor – which is revolving at 45.8 Hz. Several atoms whose transverse velocity is sufficiently low traverse a distance $L = 1.26 \text{ m}$ and hit a microchannel plate (MCP), whose principle of operation is explained in the following paragraph, where they eject electrons. The electrons are detected on the steel plate and a LeCroy WaveRunner 104 Mxi oscilloscope measures the temporal structure of the current. Two SRS PS350 high voltage power supplies provide -1500 V to the front of the MCP plate and -400 V to the back side.

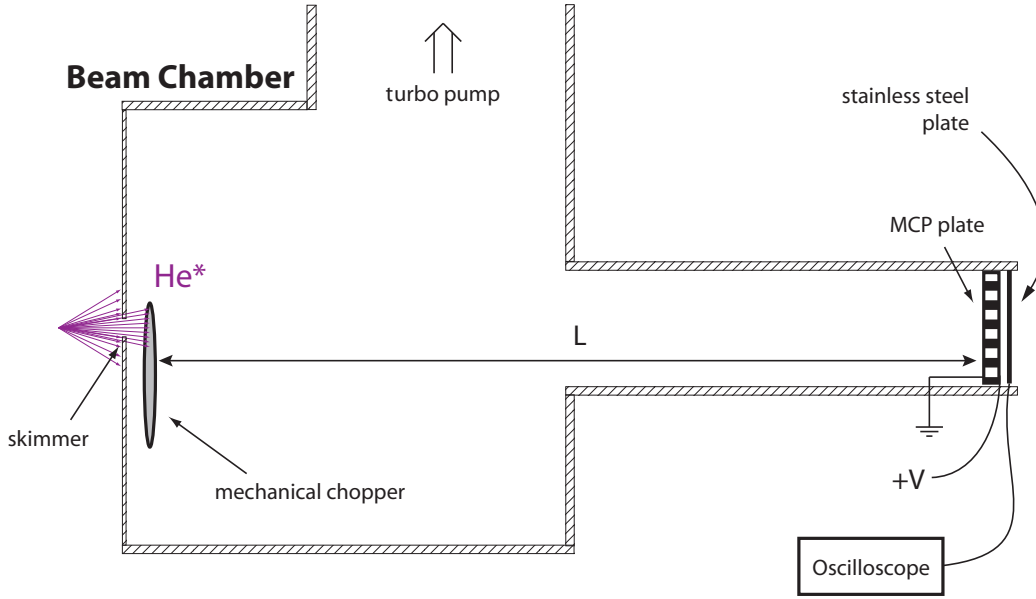


FIGURE 3.4. Experimental setup for the time-of-flight measurements. The helium beam is chopped with a mechanical chooper just after exiting the skimmer, and then travels a distance L to the MCP plate. Due to the high energy of metastable helium, electrons are ejected from the MCP and amplified. Electrons hit the stainless steel plate and a TOF spectrum is observed in the oscilloscope.

3.2.2.1. *Microchannel plate.* A microchannel plate is formed by stacking glass tubes (see Fig. 3.5), which are tens of microns in diameter, and then are cut into millimetre thick slices. Both sides of the plate are metalised in order to apply an electric potential.

The principle of operation of the MCP is shown in Fig. 3.6 [52]. An energetic particle hits the channel and ejects an electron which is accelerated in an electric field and collides with the channel wall (because of a small angle between the channel and the perpendicular axis of the MCP). Such a collision ejects secondary electrons whose number is proportional to the acquired energy of the electron.

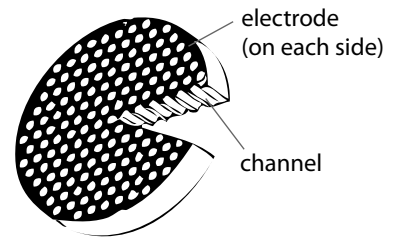


FIGURE 3.5. A cutaway view of the MCP.

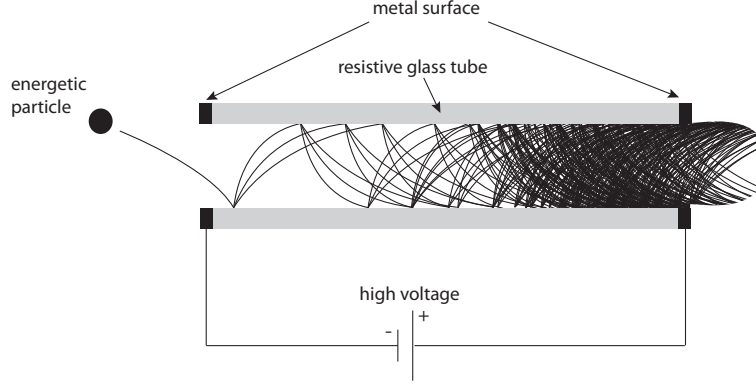


FIGURE 3.6. Principle of operation of a microchannel plate. A highly energetic particle hits a glass surface of the microchannel and ejects electrons, which are accelerated by the potential difference, between the metal surfaces, towards the other end. Those ejected electrons extract other secondary electrons from the surface, whose number increases exponentially.

The secondary electrons undergo the same process as the first electron, thus creating an electron avalanche. In this way a single particle can be amplified by a factor of $10^3 - 10^4$, making its detection far more feasible.

Our microchannel plates were manufactured by Photonis. They have 32.7 ± 0.05 mm diameter, the channels are $10 \mu\text{m}$ in diameter and are inclined by 13° to the perpendicular axis of the disc. The MCPs have 60% of open area, i.e. ratio of the holes' area to the overall area, and can be operated at the maximum voltage difference of 1.2 kV between both sides. The gain specified for 1 kV is 10^3 .

3.3. Laser setup

The laser system shown in Fig. 3.7 consists of two lasers: a custom-built, frequency doubled diode laser Toptica TA/DL-SHG 110 whose output is at 389 nm, linewidth of < 500 kHz and 100 mW output power. The second laser is Koheras fibre laser (model ADJUSTIK-Y10-10) at 1083.3 nm, linewidth < 70 kHz and 10 mW output power. All fibres, including the laser output fibre, used in our experiment are single mode, polarisation maintaining and terminated with FC/APC.

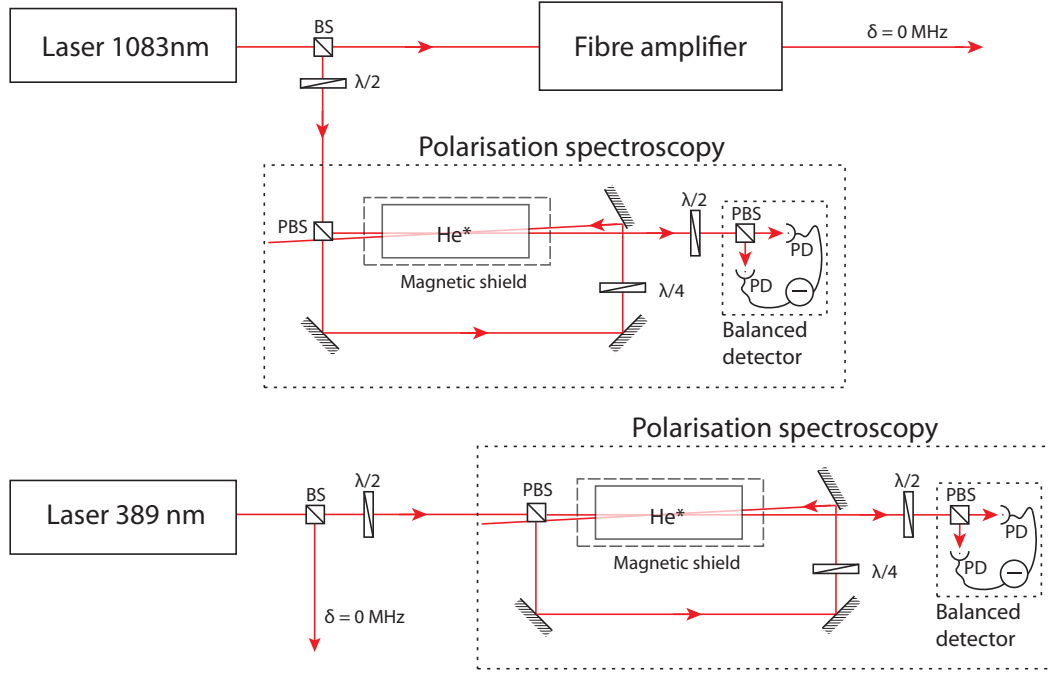


FIGURE 3.7. A schematic diagram of the optical setup. It consists of two lasers, one at 1083 nm and the other one at 389 nm, which are frequency stabilised on the helium atomic transition using a radio-frequency discharge cell (marked as He^*). The cell is placed inside a μ -metal magnetic shield. Where BS: beam splitter, PBS: polarising beam splitter, $\lambda/2$: half-wave plate, $\lambda/4$: quarter-wave plate, PD: photodiode.

3.3.1. Polarisation spectroscopy. Both lasers are frequency stabilised on the corresponding atomic transition using polarisation spectroscopy. It is one of the Doppler-free spectroscopic techniques and can directly detect the atomic dispersion. The technique was originally developed by Wieman et al. [53], although we utilise a slightly modified scheme [54, 55]. The basic scheme of Wieman features two counterpropagating beams in an atomic vapour cell – a circularly polarised pump and a linearly polarised probe beam. The pump creates an anisotropy in the medium that yields a rotation of the polarisation of the probe after traversing the sample. The detection is made with nearly crossed polarisers. The residual background light cannot be eliminated what degrades the S/N ratio. In Yoshikawa’s et al. [55] scheme a balanced detection between two orthogonal polarisation components is used with much improved S/N ratio.

Figure 3.7 shows polarisation spectroscopy scheme. A strong, σ^+ polarised pump beam optically pumps the metastables to the 2^3P_2 , $m_F = -2$ state. When a weak, linearly polarised probe beam traverses the cell, its circular components σ^+ and σ^- experience an optical birefringence — the phase velocities of σ^+ and σ^- differ. The absorption coefficients α^+ and α^- are also different for both polarisations what results in a circular dichroism. The probe is detected by a balanced⁷ detector.

The detector signal can be calculated using a simple model [56], which neglects the absorption in the glass windows and atmospheric-induced birefringence, and assumes an ideal polariser in the setup. During the derivation a convention is used that every $+$ or $-$ superscript relates to σ^+ or σ^- polarisation, respectively. Suppose that we have the same setup as in Fig. 3.7 except for the half-wave plate. A linearly x -polarised probe beam, propagating along the z direction can be written as

$$(3.1) \quad \mathbf{E} = \hat{\mathbf{x}}E_0e^{i(\omega t - kz)},$$

traversing through a pumped medium of a length L . Because of the anisotropy caused by the circularly polarised pump beam, the two circular components of the probe beam are absorbed unevenly, thus from the Lambert-Beer law follows

$$(3.2a) \quad \mathbf{E}^+ = (\hat{\mathbf{x}} + i\hat{\mathbf{y}})\frac{E_0}{2}e^{-\alpha^+L/2}e^{i(\omega t - k^+z)},$$

$$(3.2b) \quad \mathbf{E}^- = (\hat{\mathbf{x}} - i\hat{\mathbf{y}})\frac{E_0}{2}e^{-\alpha^-L/2}e^{i(\omega t - k^-z)},$$

and acquire a phase difference due to the different refractive indices $\Delta n = n^+ - n^- \neq 0$

$$(3.3) \quad \Delta\phi = (k^+ - k^-)L = (n^+ - n^-)\frac{\omega L}{c}.$$

Then inserting (3.2) and (3.3) into (3.1) yields

$$(3.4) \quad \begin{aligned} \mathbf{E} &= \mathbf{E}^+ + \mathbf{E}^- = \\ &= \frac{E_0}{2}e^{i\omega t}e^{-i\bar{n}L\omega/c - \bar{\alpha}L/2} [(\hat{\mathbf{x}} + i\hat{\mathbf{y}})e^{-i\omega L\Delta n/2c - L\Delta\alpha/4} + (\hat{\mathbf{x}} - i\hat{\mathbf{y}})e^{i\omega L\Delta n/2c + L\Delta\alpha/4}], \end{aligned}$$

⁷In the off-resonant case the half-wave plate is rotated in such a way that the signal difference vanishes.

where $\Delta\alpha = \alpha^+ - \alpha^-$, $\bar{n} = 1/2(n^+ + n^-)$ and $\bar{\alpha} = 1/2(\alpha^+ + \alpha^-)$. If the transmission axis of the PBS is tilted by a small angle θ with respect to the y -axis, then the transmitted amplitude becomes

$$(3.5) \quad E_t = E_x \sin \theta + E_y \cos \theta.$$

Furthermore, in reality the differences $\Delta\alpha$ and Δn are very small, that is $\Delta\alpha L \ll 1$ and $\Delta n L \ll 1$. Taking this into account one can express the transmitted amplitude as

$$(3.6) \quad E_t = E_0 e^{i\omega t} e^{-i\omega \bar{n} L/c - \bar{\alpha} L/2} \left[\sin \theta + \cos \theta \left(\frac{\omega}{2c} L \Delta n - \frac{i}{4} L \Delta \alpha \right) \right].$$

Since the detector signal is proportional to $I_t = c\epsilon_0 E_t E_t^*$, where c is the speed of light in vacuum and ϵ_0 denotes the electric permittivity of vacuum, the intensity reaching one of the detectors follows

$$(3.7) \quad I_t(\theta) = I_0 e^{-\bar{\alpha} L} \left\{ \sin^2 \theta + \frac{\omega}{2c} L \Delta n \sin(2\theta) + \left[\left(\frac{\omega}{2c} L \Delta n \right)^2 + \left(\frac{L}{4} \Delta \alpha \right)^2 \right] \cos^2 \theta \right\},$$

where $I_0 = \epsilon_0 c |E_0|^2$. A balanced detection corresponds to $\theta = \pm 45^\circ$, hence the output of the balanced detection is

$$(3.8) \quad \Delta I_t = I_t(45^\circ) - I_t(-45^\circ) = 2I_0 e^{-\bar{\alpha} L} \frac{\omega}{c} L \Delta n,$$

which shows a complete dispersion profile with no background.

Because of the dispersion signal at the output of the detector, there is no necessity of frequency modulation and usage of a lock-in amplifier, therefore it makes a stabilisation of the laser frequency far more feasible. Moreover, its sensitivity is 2 – 3 orders of magnitude larger than that of saturation spectroscopy. Therefore polarisation spectroscopy is preferred over frequency modulated saturation spectroscopy. Figure 3.8 shows both spectroscopy signals. Both signals are pressure and power broadened and the linewidth of the spectroscopy signal is about 40 MHz. It can be seen that the S/N ratio is in favour of polarisation spectroscopy.

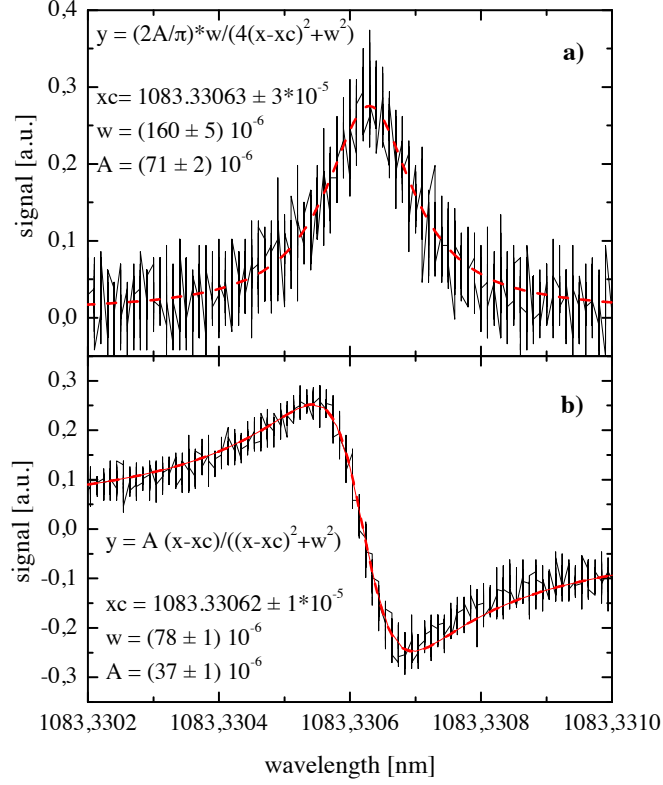


FIGURE 3.8. A comparison between saturation spectroscopy (a), and polarisation spectroscopy (b) on the $2^3S_1 - 2^3P_2$ helium transition. Both plots contain fitted curves (dashed) from which the Doppler-free linewidths are inferred. The linewidths are 40.9 ± 1.3 MHz and 39.8 ± 1.0 MHz for (a) and (b), respectively.

To lock the laser we use an 11 cm long and 2.5 cm in diameter glass cell containing helium vapour. Helium in the glass cell is excited to the metastable state using a radio-frequency (rf) discharge, as mentioned in Sec. 3.1.2. The cell is placed in a μ -metal magnetic shield, in order to attenuate stray magnetic fields which disturb the spectroscopy, as well as to separate the rf field from the rest of the experiment.

3.4. Collimation

Figure 3.9 presents an implementation of the collimation with curved wavefronts in 2D. Optical access to the atoms is provided by four antireflection coated windows of 17 cm

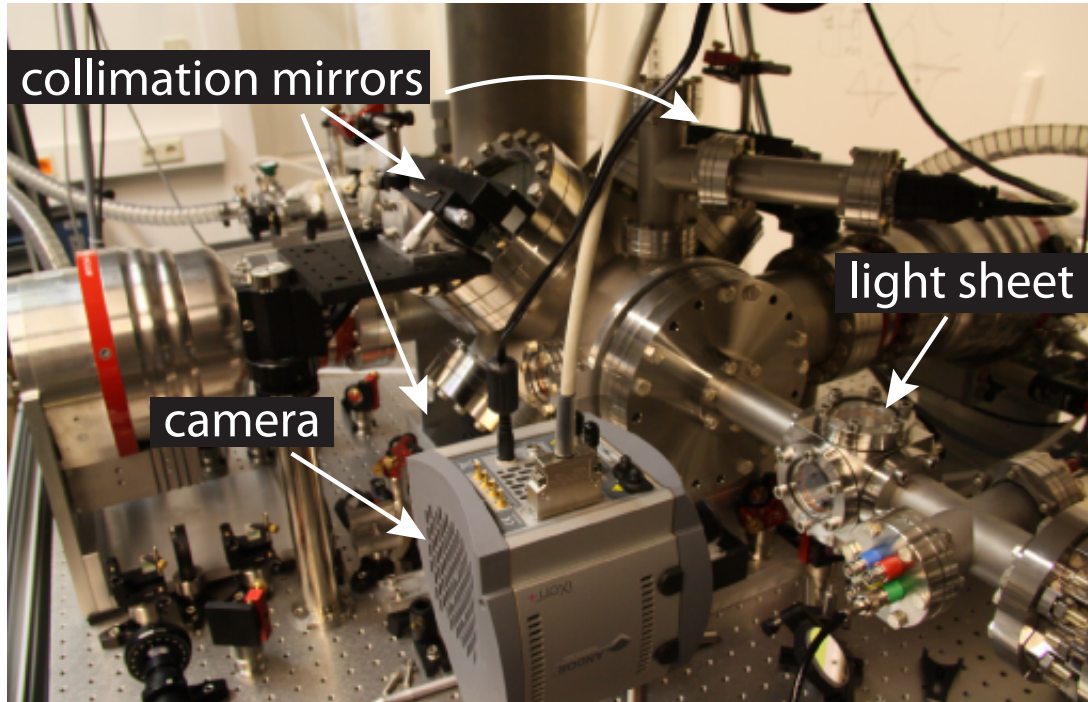


FIGURE 3.9. Picture showing an actual collimation and detection setup. Four large area mirrors, which are used for reflecting the collimation beams, are centred on the vacuum windows outside the vacuum chamber. The collimation beams are injected from the side of the two upper mirrors.

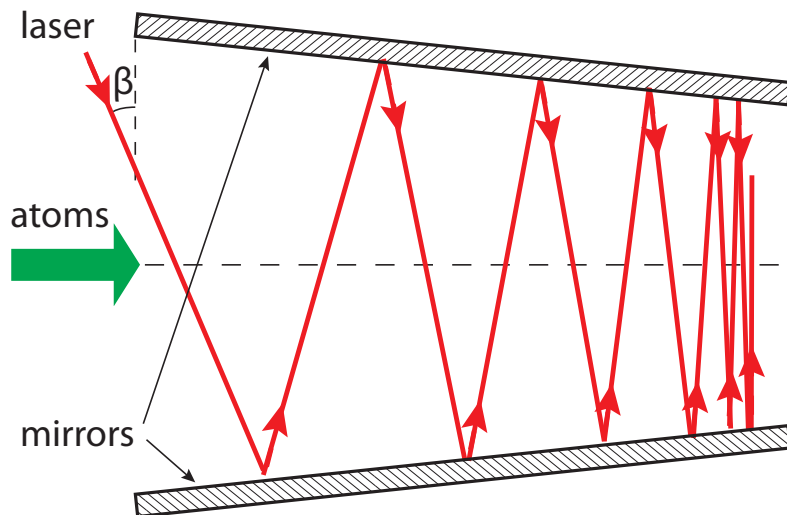


FIGURE 3.10. Scheme of the laser beam alignment between two collimation mirrors. The same alignment is used for the perpendicular pair of mirrors. All angles are greatly exaggerated.

in diameter. The windows are located right at the beginning of the beam chamber, so that the laser beam interacts with atoms right behind the skimmer. Outside the vacuum system four mutually perpendicular mirrors ($2 \times 18 \text{ cm}^2$) are placed, whose surface is pairwise slightly inclined with respect to the atomic beam axis (see Fig. 3.9). Two 1083 nm laser beams (6 mm in diameter) enter the vacuum system through the two upper windows from the side of the collimation mirrors. Each beam carries 25 mW of power and enters the vacuum chamber at a capture angle $\beta \approx 20 \text{ mrad}$ and interacts with atoms on the 13 cm-path which corresponds to the collimation radius $R = 7 \text{ m}$. There are eleven reflections on each mirror.

The (un-)collimated beam was sliced at 45° (approximately 32 cm behind the collimation region) by the $\sim 0.6 \text{ mm}$ thick light sheet of resonant 389 nm light. Scattered photons were collected using an Andor IXON 885 EMCCD camera with a Fujinon HF25HA-1B (28.6 mm in diameter) objective placed $14.5 \pm 0.2 \text{ cm}$ away, perpendicularly to the vacuum chamber.

CHAPTER 4

Results

4.1. Intensity measurements

The intensity of the metastable helium beam coming out of the source chamber depends on many parameters. Their influence on the performance of the source was investigated and the results are presented in this section. The intensity of the source was maximised with respect to the orifice diameter of the nozzle (Sec. 4.1.1) and the shape of the skimmer (Sec. 4.1.2). Moreover, the flux was measured as a function of the inlet pressure, the position of the needle and the voltage applied to the cathode (Sec. 4.1.3). The reader should notice that the final intensity of the source is presented in Sec. 4.1.3 and this should not be compared to other sections, since many parameters were varied throughout the course of flux measurements and not all data might have been taken under the same conditions. All intensity measurements were performed with the Faraday cup (described in Sec. 3.2.1), whose output was read out using the picoammeter.

To express the current of the picoammeter in terms of the number of atoms hitting the steel plate, one needs to convert 1 nA to atoms/s. Knowing the distance between the skimmer and the steel plate $l = 10.8$ cm, the size of the aperture of the mask $d = 0.55$ cm, the electron ejection efficiency of He^* $\eta = 69 \pm 9\%$ and $1 \text{ C} = 6.24 \times 10^{18}e$ where e is the unitary charge, one gets

$$(4.1) \quad 1 \text{ nA} = 10^{-9} \frac{\text{C}}{\text{s}} = 6.42 \times 10^9 \frac{\text{atoms}}{\eta \text{ s}} = (9.05 \pm 1.21) \times 10^9 \frac{\text{atoms}}{\text{s}}.$$

Then the flux of atoms per steradian follows as

$$(4.2) \quad \Phi = \frac{1 \text{ nA}}{\text{sr}} = (4.44 \pm 0.21) \times 10^{12} \frac{\text{atoms}}{\text{s sr}}.$$

Apart from atoms in the metastable triplet state, one still has a negligible contribution of the extreme ultraviolet (XUV) photons ($< 1\%$ [39]) and a significant amount of helium in the metastable singlet state 2^1S_0 ($\sim 10\%$ [57]), whose radiative lifetime is 19.7 ± 1.0 ms

TABLE 4.1. Intensity of the helium beam for three nozzles with different orifice diameters.

Orifice diameter [mm]	Intensity [10^{14} atoms s $^{-1}$ sr $^{-1}$]
0.30	0.84 ± 0.18
0.50	2.44 ± 0.51
0.75	0.90 ± 0.19

[58] and whose electron ejection probability from the steel plate is $\eta' = 53\%$ [51], to the current.

4.1.1. Nozzles comparison. Primarily in the experiment the nozzle with 0.5 mm orifice was used. Two other nozzles (orifice diameter: 0.3 mm and 0.7 mm) were installed. Yet their performance was disappointing, since the discharge was “running” in a different regime¹ than with 0.5 mm. Thus the flux of atoms was few times lower than with the 0.5 mm nozzle, what presents Table 4.1.

The reason for the discrepancy in the flux magnitude is different for each of the “non-standard” nozzles. For the 0.3 mm nozzle, the majority of the particles does not get through the orifice leading to the emergence of another kind of discharge. In the case of the 0.75 mm nozzle, the other kind of discharge prevails over the “right” one, because most of the helium leaks out through the opening of the nozzle directly into the source vacuum chamber and too little helium flows backwards through the glass tube. For very specific values of inlet pressure, voltage and needle position, we were able to achieve discharge in the proper regime, although it was very unstable and it could not be sustained. Hence, the nozzle with 0.5 mm orifice was chosen for further measurements.

4.1.2. Skimmer comparison. Further optimisation of the flux was done by comparing output through three different skimmers listed in Table 3.1. According to the theory (see Section 1.3.2) the flat skimmer should deliver the lowest flux and the other

¹The colour of the discharge glow is turquoise-white rather than blue-white.

TABLE 4.2. Intensity of the helium beam for three different skimmers (for respective dimensions see Table 3.1) at room temperature $T_0 = 298.3$ K.

Skimmer type	Intensity [10^{14} atoms s^{-1} sr^{-1}]
short, curved	8.44 ± 1.77^a
long, curved	2.44 ± 0.51
flat	0.9 ± 0.19

^a This value was measured with the new skimmer holder whose design is superior to the older one. With the older holder we achieved a value $(1.8 \pm 0.38) \times 10^{14}$ atoms s^{-1} sr^{-1} .

two should perform equally well. Data presented in Table 4.2 were taken at room temperature $T_0 = 298.3$ K. They confirm that the flat skimmer performs the worst, although unexpected is the striking difference in the output through the two curved skimmers. This may be owing to the fact that we used different skimmer holders, since the distance between the nozzle and the skimmer wall was too small to fit the long curved skimmer in the same holder as the short one. The choice of the optimal skimmer was quite straightforward, the short curved skimmer was used during the rest of the experiment.

4.1.3. Final flux of atoms. This section summarises the flux measurements. Figure 4.1 presents number of atoms per second per steradian in the beam as a function of applied voltage for three different inlet pressures. For two inlet pressures ($p = 6$ mbar and $p = 20$ mbar) the number of metastable helium atoms clearly saturates for voltages above 3 kV. This may be a result of the fact that the metastable helium density is so high that further production of He^* is suppressed by Penning ionisation, cf. Appendix A. Due to unknown reasons for $p = 13$ mbar this saturation does not take place. However, the source may be regarded as predictable, since its output is on the order of $\sim 10^{14}$ atoms s^{-1} sr^{-1} and above. On the one hand one wants to get as high number of atoms in the MOT as possible, therefore one would like to operate the source at its maximal output values, on

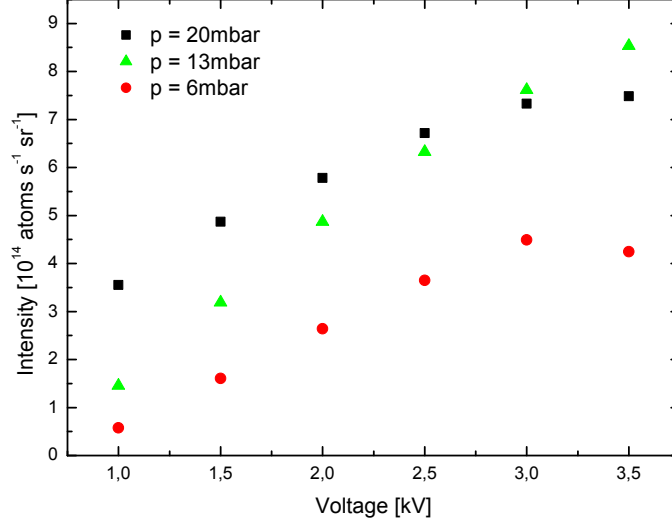


FIGURE 4.1. Helium beam intensity as a function of voltage which is applied to the tungsten needle for three different inlet pressures. The measurements were done at cryogenic temperature. On the horizontal axis are given absolute values of applied voltage. Size of the points may be regarded as an error bar.

the other hand driving the source at higher voltages leads to quicker wear of the needle tip and the discharge is less efficient. Moreover higher operating voltages result in faster atoms coming out of the source. This is the subject of the following section.

4.2. Time-of-flight

The time-of-flight (TOF) measurements were performed with a setup detailed in Sec. 3.2.2 at cryogenic temperature. The velocity distribution of the supersonic beam is given by Eq. (1.20), the corresponding TOF signal on the detector $I(t)$, is provided by the formula [59]

$$(4.3) \quad I(t) \propto \left(\frac{L}{t\sqrt{2k_B T/m}} \right)^5 e^{-\frac{m(L-t\bar{v})^2}{2t^2 k_B T}},$$

where L is the distance between the detector and the chopper, T and $\bar{v}$ are given by Eqs. (1.17) and (1.14), respectively. A typical TOF spectrum is shown in Fig. 4.2 with

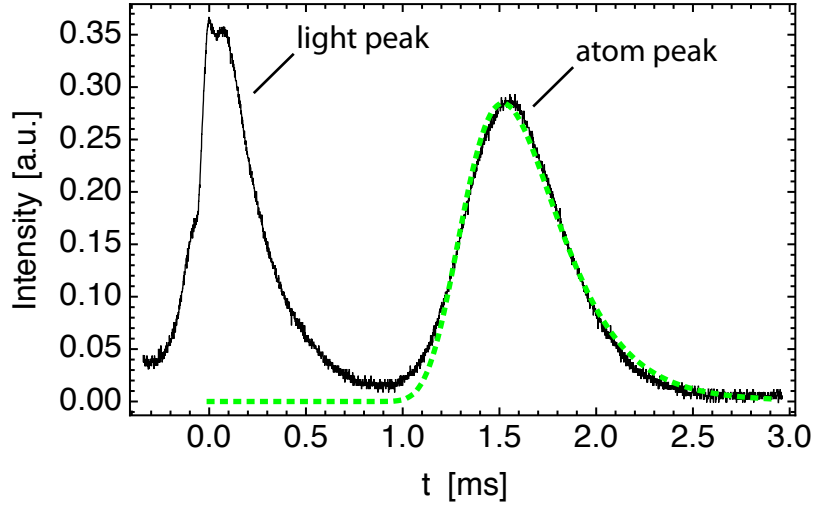


FIGURE 4.2. Typical time-of-flight signal (solid line) with fitted curve (dashed). The sharp peak is the light peak and it defines the zero point on the time axis. It was taken for $V = -2$ kV with fitted parameters $M = 5.8 \pm 0.5$, $T_0 = 58 \pm 2$ K and $\bar{v} \approx 810 \pm 29$ m/s.

a fitted curve from Eq. (4.3). A zero time reference is given by the XUV photon peak from the source. The Mach number of our beam is on the order of 5 which is a lower value than theoretically calculated $M_{\text{theory}} = 12.8 \pm 1.3$ from Eq. (1.19) for $x = 4 \pm 1$ mm. The difference may emerge from the fact that the theoretical value is given for the beam at the point where it enters the skimmer, although TOF flight signals were measured at the distance $L = 1.26$ m behind the skimmer. So one would have to assume that the Mach number does not change during the flight. From the taken spectra one can also calculate the peak velocity $\bar{v}$ and plot it versus the driving voltage (Fig. 4.3) and the needle position (Fig. 4.4).

Figure 4.3 presents the peak velocity of the atomic beam as a function of voltage applied to the needle for three different pressures. Except for the first three points at $p = 20$ mbar, whose behaviour is not well understood², the data points for given voltage mutually differ by 5 – 15%. This result is expected, since higher inlet pressure increases the temperature in the discharge region.

²This abnormality may result from insufficient cooling power of the cryostat, yet it still does not explain why the points lie on the same line when increasing the voltage.

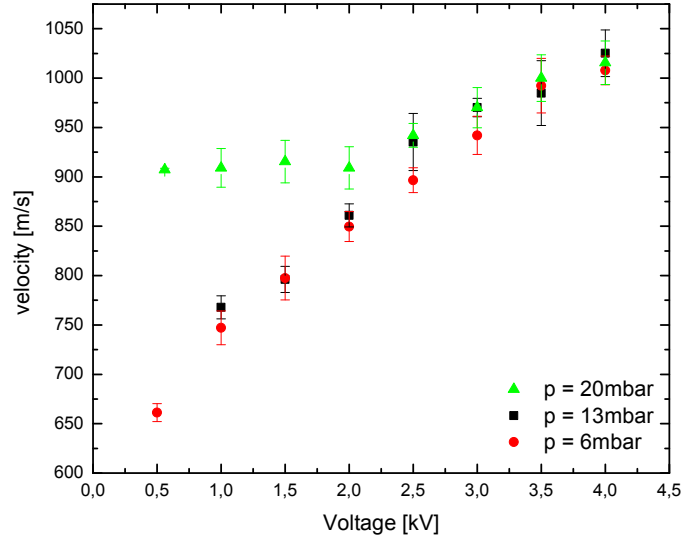


FIGURE 4.3. Results of the time-of-flight measurement. Peak velocity of the helium atoms vs. the voltage applied to the anode for three inlet pressures at cryogenic temperature. Error bars result from the fitting error.

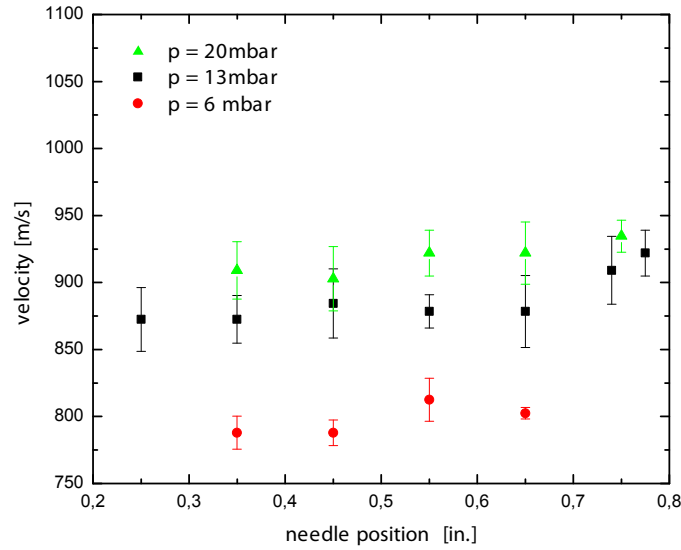


FIGURE 4.4. Peak velocity of the helium atoms vs. position of the tungsten needle for three inlet pressures at cryogenic temperature. The voltage applied to the needle was set to $V = -2$ kV.

Figure 4.4 implies that changing the needle position has no significant influence on the velocity of the atoms³. It influences only the discharge itself if it continues to “run” or

³The same was observed in the flux measurements, though no systematic data were taken.

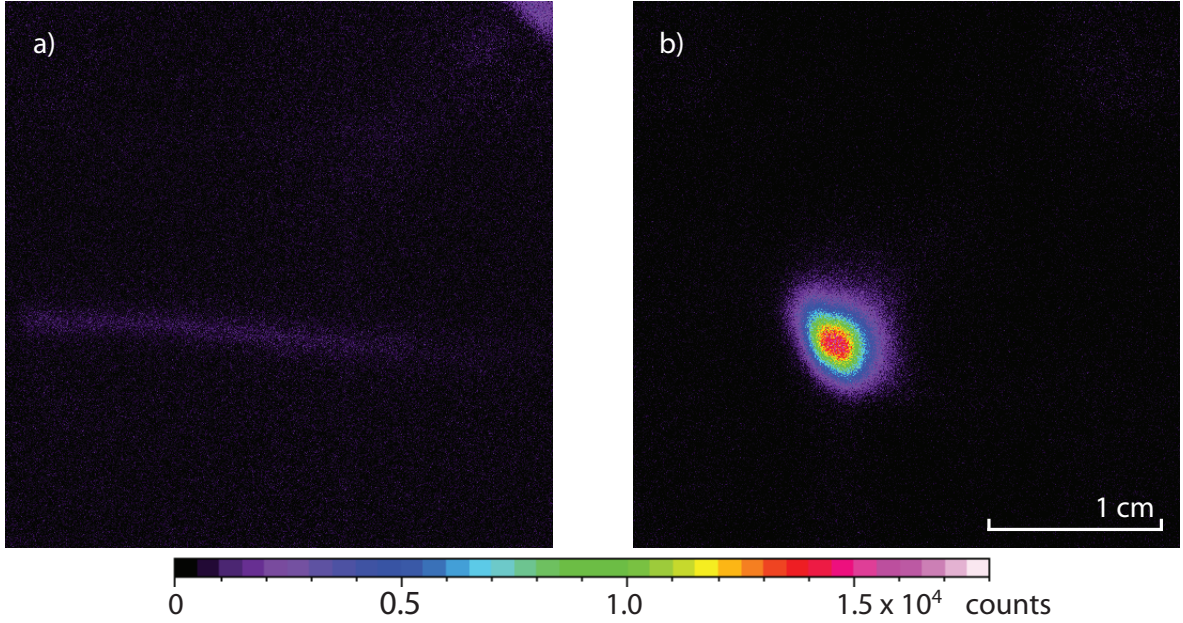


FIGURE 4.5. Fluorescence signal from the CCD camera taken with 300 ms exposure time. **a)** Uncollimated metastable helium beam. Only its small part fluoresces, because the laser is just resonant with atoms with zero transverse velocity. **b)** Optimised fluorescence spot of the collimated atomic beam. Area of the spot (FWHM) is about $6.2 \pm 1.1 \text{ mm}^2$ and the enhancement factor is about 26.

not. Data points for $p = 6 \text{ mbar}$ illustrate that only for a certain range of needle position the discharge “runs”.

Since the output of the source is on the same order of magnitude (cf. Section 4.1.3), we decided to set the upper limit of the available peak velocities to $\sim 850 \text{ m/s}$ and optimise the future of the experiment for this value. This peak velocity is about 20% lower than in other metastable helium sources [41, 49], at the same time producing a comparable number of He^* atoms per second per steradian. Slower atoms require a shorter Zeeman slower, which leads to smaller atom losses due to divergence of the beam (the following Section describes collimation of the helium beam, yet it is still slightly divergent) in the Zeeman slower.

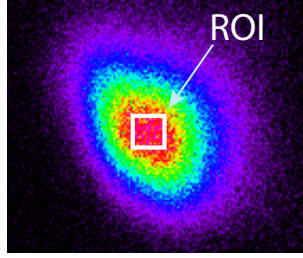


FIGURE 4.6. Zoomed in collimated beam with region of interest (ROI) marked as a white square.

4.3. Collimation

Figure 4.5a presents the fluorescence of the uncollimated beam⁴. Since only atoms whose transverse velocity is zero are resonant with the light sheet, we see just a narrow excerpt of the quickly expanding helium beam. Switching on the collimation beams leads to emergence of a bright spot in the fluorescence (Fig. 4.5b), what is a signature of a well-working collimation. The full width at half maximum (FWHM) area of the spot is about $6.2 \pm 1.1 \text{ mm}^2$. Estimation of the enhancement factor over the uncollimated case was done by selecting a $0.94 \times 0.94 \text{ mm}^2$ square region of interest⁵ (ROI, see Fig. 4.6). Within the ROI the mean count number, the standard deviation of the mean and the maximum count number were calculated for both uncollimated and collimated beam. The results are presented in Table 4.3. The enhancement factor was taken as the ratio between the mean values for collimated and uncollimated beams and it is approximately equal to 26.

⁴It can be seen that the uncollimated beam is hard to spot when reading the hard copy of the Thesis.

⁵Height of the area of ROI was selected such that it is equal to the width of the fluorescence stripe of the uncollimated beam.

TABLE 4.3. Number of counts of selected ROI
(region of interest) for uncollimated and colli-
mated beam.

Counts	Uncollimated beam	Collimated beam
Mean	541	13814
Standard deviation	341	1229
Maximum	1545	15970

CHAPTER 5

Summary and Outlook

Starting this Master's Thesis project in October 2008, apart from optical tables and an air conditioning, the lab was empty. Lots of effort was done to develop the lab to the state in which it is right now. This thesis describes part of the work which was done during the last year. In the scope of the Master's Thesis the source of metastable helium atoms was built. The increase in the intensity of the atomic beam was achieved by additional collimation with laser light right after the skimmer.

The intensity of the metastable helium beam was measured before collimation. It is approximately $\sim 1 - 8 \times 10^{14} \text{ atoms s}^{-1} \text{ sr}^{-1}$. At the same time the peak velocity is in the range of $650 - 1050 \text{ m/s}$. In comparison to other helium BEC groups [41, 49], these values are promising as our source delivers the same order of magnitude in intensity, though at lower peak velocity. This enables building a shorter Zeeman slower, which results in smaller atom losses due to divergence of the helium beam, therefore creating a metastable helium MOT may be easier.

In the second part of the thesis the helium beam was collimated in two dimensions right after exiting the skimmer using a converging wavefront technique. Laser beams were 6 mm in diameter and each carried 25 mW of power. The collimation length was 13 cm. With laser collimation we gained an approximately twenty-six-fold increase in the number of atoms per second per steradian.

Future work includes mounting and characterisation of the Zeeman slower, which had already been built. It is 1.36 m long and is comprised of two solenoids, which overall contain 1.6 km of copper wire with a rectangular cross section ($3 \times 1 \text{ mm}^2$). The wire is wound on a double-wall vacuum pipe with an inner diameter of 3.5 cm. Afterwards, the next step includes building up the MOT chamber and the magnetic trap in order to trap the atoms. By the time of finishing this thesis, the Zeeman slower is being characterised.

The first part of the MOT chamber is on the optical table with vacuum pumps, being tested. Finally the magnetic trap was designed and it is now being built.

APPENDIX A

Helium atom

Helium is the second lightest element and the simplest multielectron atom. Furthermore it is a noble gas, that is all its electron shells are closed, and has two stable isotopes: ^4He and ^3He with the natural abundance of 99.99986% and 0.00014%, respectively [60]. Of particular interest is the ^4He isotope, which is a boson and which is used in the experiment. Due to the fact that its nucleus consists of two protons and two neutrons its nuclear spin vanishes, $I = 0$, there are no hyperfine states in the energy level structure (see Fig. A.1), thus the level diagram is simpler than in ^3He .

A.1. Spectroscopic properties of ^4He

Helium in its ground singlet state 1^1S_0 has no optically available transitions, hence for radiative cooling one has to use the 2^3S_1 state, which is a metastable triplet state (He^*) with a lifetime $\tau = 7900\text{ s}$ [61]. The reason for such a long lifetime of the metastable triplet state is that the deexcitation to the ground is forbidden for [58]:

- electric dipole transition since the parity does not change;
- electric quadrupole, because of $J = 1 \rightarrow J' = 0$ transition;
- magnetic dipole due to the $\Delta S = 0$ violation.

Furthermore, only contribution of the spin-dependent relativistic correction operators to the magnetic dipole operator [62] give rise to the finite lifetime of the 2^3S_1 state.

A.2. De-excitation processes

If the metastable helium atom collides inelastically with another particle, whose ionisation energy is less than the internal energy of He^* (19.8 eV), then the He^* can decay to the ground state and the other atom is ionised. In our experiment, He^* may collide with either another metastable particle or a background gas particle (typically H_2), whose

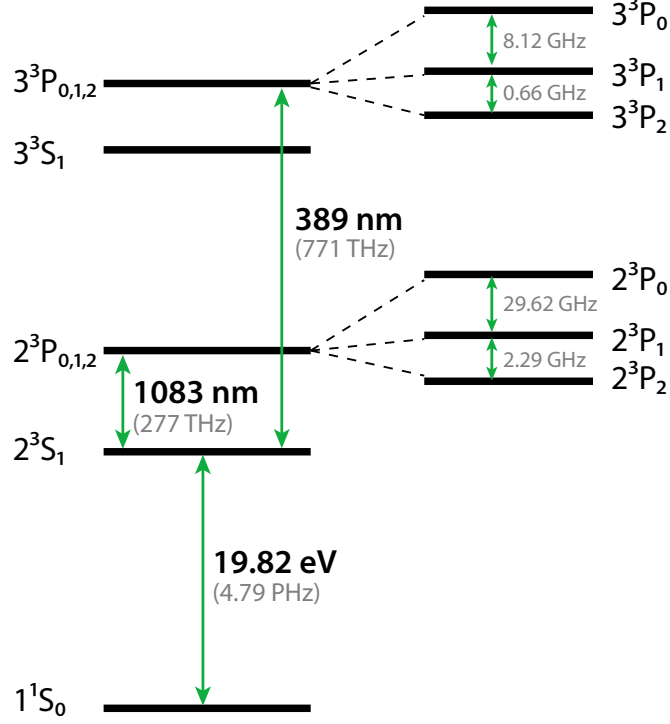
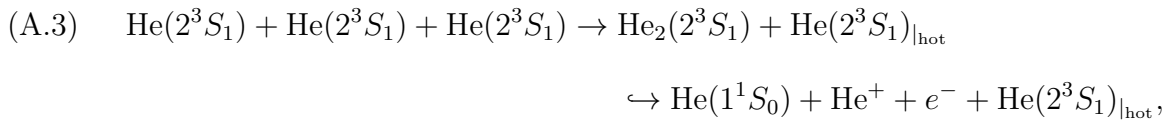
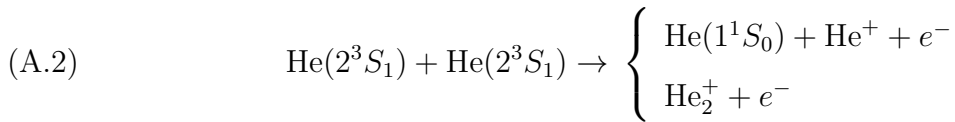
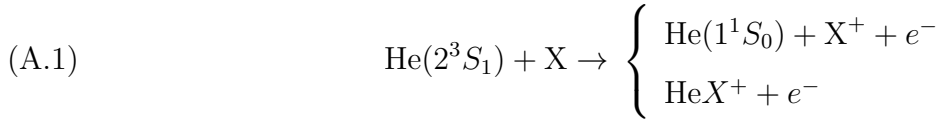


FIGURE A.1. Selected energy levels of the ${}^4\text{He}$ atom. The spectroscopy data is taken from [63].

ionisation energies are, respectively, 4.77 eV and 15.5 eV [64]. Each ionisation event leads to unwanted atom losses.

This ionisation process was predicted in 1927 by Penning [65], thus it is often being referred to as Penning ionisation (or Penning collision). Depending on the collision partner, one can distinguish several reactions giving rise to the Penning ionisation



where X denotes a background molecule. Each of the processes described by Eqs. (A.1)–(A.3) dominates at different densities of metastable helium.

TABLE A.1. Spectroscopic data for transitions $2^3S_1 \rightarrow 2^3P_2$ and $2^3S_1 \rightarrow 3^3P_2$ in ^4He . All the data are taken from [29].

		$2^3S_1 \rightarrow 2^3P_2$	$2^3S_1 \rightarrow 3^3P_2$
Wavelength	λ [nm]	1083.33	388.98
Lifetime	τ [ns]	98.04	106.83
Linewidth	$\Gamma/2\pi$ [MHz]	1.62	1.49
Cross section for absorption^a	σ_{ge} [10^{-15} m^2]	560.4	72.2
Saturation intensity^a	I_s [mW/cm ²]	0.17	3.31
Recoil frequency	$\omega_r/2\pi$ [kHz]	42.46	329.35
Capture limit:	v_c [m/s]	1.76	0.58
	T_c [mK]	1.49	0.16
Doppler limit:	v_D [cm/s]	28.44	27.25
	T_D [μK]	38.95	35.75
Recoil limit:	v_r [cm/s]	9.2	25.6
	T_r [μK]	4.075	31.61

^a The values for the cross section and the saturation intensity apply for the strongest transition between magnetic sublevels.

He*+X: These are the collisions with background gas particles which dominate at low densities (up to $\sim 10^8$ atoms/cm³ [66]), when collisions between two helium atoms are rare. Equation (A.1) implies that either metastable helium collides with a background molecule, transferring its internal energy to the other molecule which in the end is ionised or a molecular ion, through associative ionisation, is formed. Each of the two reactions yield an ion, and in both cases He* is lost from the trap.

He*+He*: If one increases the density of the sample, then the reaction from Eq. (A.2) is predominant, and both atoms undergo the same physical processes as particles in Eq. (A.1). The molecular ion He₂⁺ is created with a probability of a few percent [67].

He*+He*+He*: Since Eq. (A.3) is a three-body process and mainly occurs at the densities $> 10^{12}$ atoms/cm³, it is usually found at the end of evaporative cooling, on the road to the Bose-Einstein condensate (BEC). Therefore it is not so important in the scope of this work. In the collision a fast decaying excited He₂ molecule is created and a “hot” He* (~ 400 μ K), whose energy corresponds to the binding energy of the He₂.

To prevent the second category of Penning ionisation, one has to spin polarise the atomic ensemble, what in the case of metastable helium corresponds to transferring the atoms to the 2^3S_1 , $m_F = 1$ magnetic level. Then the value of the total spin, on the left-hand side of Eq. (A.2), equals 2, whereas on the right-hand side 0 or 1, hence due to spin conservation this collision is forbidden. It has been experimentally shown that usage of this technique suppresses the unwanted collisions between helium atoms by a factor of 20 [68].

Penning ionisation does not play a crucial role in obtaining a collimated beam of metastable helium atoms, although in the discharge region of the source (see Subsec. 3.1.3), due to the high density of helium, Penning collisions might be very important. However, if one wants to create a BEC with He*, what is the ultimate goal of the experiment, then one should remember to spin-polarise the atoms in the magneto-optical trap (MOT), before loading them into the magnetic trap.

APPENDIX B

Curriculum Vitae

Personal information

Surname(s) / First name(s)

Address(es)

Telephone(s)

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Nationality(-ies)

Date of birth

Kotyrb Mateusz

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Polish

December 26, 1985

Education

September, 2008 – January,
2010

Wien Universität, Wien, Austria

Physics
prof. A. Zeilinger Group

October, 2004 – June, 2008

Jagiellonian University, Kraków, Poland

Interdisciplinary Studies of Mathematical and Natural Sciences
Specialization: Atomic Physics
prof. W. Gawlik Group

October, 2007 – February, 2008

Wien Universität, Wien, Austria (ERASMUS Scholarship)

Physics
prof. M. Arndt Group

September, 2000 – June, 2004

II LO im. A.F. Modrzewskiego, Rybnik, Poland (High School)

Mathematical and Physical Faculty

September, 2002 – June, 2004

Youth Palace, Katowice, Poland

Member of Creative Group Quark
Physics Faculty

Spoken languages

Mother tongue(s)

Polish

Self-assessment
European level^()*

English

German

Understanding		Speaking		Writing
Listening	Reading	Spoken interaction	Spoken production	
C1 Proficient user	C1 Proficient user	C1 Proficient user	C1 Proficient user	C1 Proficient user
B2 Independent user	B2 Independent user	B2 Independent user	B2 Independent user	B2 Independent user

^(*) Common European Framework of Reference (CEF) level

Skills

Operating systems

Linux, Windows, Mac OS X

Programming languages

C, Pascal, LabView

Other Computer skills

L^AT_EX, Mathematica, MatLab, Origin, Illustrator, Photoshop, Acrobat, 3ds Max, SolidWorks

Driving licence

B category – cars up to 3,5 tons

Publications

S. Pustelny, M. Gring, A. Wojciechowski, K. Sycz, M. Kotyrba and W. Gawlik

Novel method of optical magnetometry (in Polish)
Elektronika **5/2008**, 26 (2008)

S. Pustelny, A. Wojciechowski, M. Gring, M. Kotyrba, J. Zachorowski and W. Gawlik

Magnetometry Based on Nonlinear Magneto-Optical Rotation with Amplitude-Modulated Light
J. Appl. Phys. **103**, 063108 (2008)

S. Pustelny, A. Wojciechowski, M. Kotyrba, K. Sycz, J. Zachorowski, W. Gawlik, A. Cingoz, N. Leefer, J. M. Higbie, E. Corsini, A. O. Sushkov, M. P. Ledbetter, S. M. Rochester, D. F. Jackson Kimball, and D. Budker

All-optical atomic magnetometers based on nonlinear magneto-optical rotation with amplitude modulated light
Proc. of SPIE, **6604**, 660404 (2007)

M. Kotyrba

Experimental investigation of nonlinear Faraday phenomenon in rubidium vapour (in Polish)

	Internal department paper M. Kotyrba <i>Investigation of resonances of coupled LC circuits</i> (in Polish) Internal department paper
Posters	W. Gawlik, M. Gring, M. Kotyrba, S. Pustelny, A. Wojciechowski, J. Zachorowski, D. Budker, A. Cingoz, N. Leefer <i>Sensitive optical magnetometry based on nonlinear magneto-optical rotation</i> World of Photonics – CLEO, Munich, Germany 2007 S. Pustelny, A. Wojciechowski, M. Kotyrba, K. Sycz, J. Zachorowski, W. Gawlik, A. Cingoz, N. Leefer, J. M. Higbie, E. Corsini, M. P. Ledbetter, S. M. Rochester, D. F. Jackson Kimball, V. V. Yashchuk, and D. Budker <i>Magnetometric aspects of nonlinear magneto-optical rotation with modulated light</i> Sunny Beach, Bulgaria, 2006
Hobbies	cycling, stock exchange, computers, football, volleyball, tennis, logical puzzles, games

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