MASTER

Ultra cold electron and ion beams

Reijnders, M.P.

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Merijn Reijnders
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Eindhoven University of Technology
Physics Department

**Supervisor:**
dr. ir. O.J. Luiten
Abstract

Ultra short electron bunches are used in many areas of science and technology. For instance, they are used as a probe for observing very fast changes in structures in femtosecond chemistry[1] or needed for the operation of X-ray Free Electron Lasers (XFEL)[2]. The over-all quality of a particle beam, which determines its applicability in applications, can be expressed in terms of the brightness \( B \), the current density per unit solid angle and per unit energy spread[3].

At present the state-of-the-art source of high brightness electron bunches is the radio-frequency (rf) photogun. In these devices free electrons are created by photoemission[4]. This process gives rise to initial velocities of the electrons, which are randomly distributed in all three directions. These electrons with random velocities can be looked at as if they were coming from a thermionic source with an effective temperature of about 5000K.

It can be shown that the brightness \( B \) scales inversely proportional to the initial temperature \( T \) and the initial density[5]. In this project is tried to decrease this initial temperature, in contrast to all other projects aiming at improving the brightness, that try to increase the initial charge density instead. This can be done by photo-ionization of a cloud of laser cooled rubidium atoms just above its ionization threshold[6], so an ultra cold plasma is formed. Cold electrons, down to 10 K, can be extracted by an electric field pulse, 2 – 3 order of magnitude lower than photoemission.

In this master project a Magneto-Optical Trap (MOT) is build to laser cool and trap rubidium atoms. A stable MOT was accomplished that could trap over \( 10^9 \) atoms directly from the background vapor. The measured temperature of 200 ± 100 μK is consistent with the expected Doppler temperature.

In the first ultra cold beam experiment performed, the atoms in the MOT were excited to a highly excited state, a so-called Rydberg state. With use of an electric field pulse, these atoms can be field ionized, so the distribution of states can be measured in time. The spontaneous evolution of the Rydberg cloud to an ultra cold plasma (UCP) due to collisions is observed. More than 1 pC of charge was extracted from the cloud.

With use of an image intensifier, a multi channel plate (MCP), for the first time images are made of the ultra cold electron and ion beams. From the widths of the distribution on these images, an upper estimate of the temperature can be deducted. This gives an upper temperature of 40 K, that corresponds to an initial emittance of about 0.05 mm mrad. The precise dynamics in time of the widths of the measured distribution on the phosphor screen are not completely understood. But the strong correlation with the total measured charge, suggests that space-charge plays an important role.
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Chapter 1

Introduction

Ultra short electron bunches are used in many areas of science and technology. For instance, they are used as a probe for observing very fast changes in structures in femtosecond chemistry[1] or needed for the operation of a free electron laser[2]. With even brighter bunches, exciting new applications may become possible, as for example single-shot time-resolved electron microscopy at sub-picosecond timescales[7].

The behavior of particles in a beam or bunch can be quite complex. The overall quality of a particle beam, which determines its applicability in applications, can be expressed in terms of the brightness $B$, the current density per unit solid angle and per unit energy spread[3]. At present the state-of-the-art source of high brightness electron bunches is the radio-frequency (rf) photogun. In these devices free electrons are created by photoemission by shooting a high power femtosecond laser pulse at a metal cathode surface[4]. The photoemission process itself gives rise to initial velocities of the electrons, which are randomly distributed in all three directions. These electrons with random velocities can be looked at as if they where coming from a thermionic source with an effective temperature of about $5000K$.

It can be shown that the brightness $B$ scales inversely proportional to the initial temperature $B \propto T^{-3/2}$ and the initial density[5]. All current used strategies for improving the brightness of pulsed electron sources have in common that they try to increase the current density at the source. However the current density cannot be increased indefinitely, because this results in non-linear space charge fields which spoil the brightness.

In all those cases, the angular spread, which can be expressed by the effective electron temperature of the source, is kept constant. In this project a totally different approach is chosen: the electron temperature of the source is reduced, instead of the source area. This is accomplished not by photo-emission from a metal cathode, but by photo-ionization of a cloud of laser cooled rubidium atoms just above its ionization threshold[6]. This results in an ultra cold plasma with electron temperatures down to $10K$, 2 – 3 orders of magnitude lower than conventional photoemission based sources. That means that an enormous gain in brightness may be possible with this new concept. To use an ultra cold plasma as a source for pulsed electrons, the electrons have to be stripped from the plasma and accelerated with use of a fast high voltage electric field pulse.

In this master thesis the results are presented of the first experiments done on the creation of charged particle beams out of an ultra cold plasma. First the basic theory of laser cooling and ultra cold plasma is explained briefly. A short introductional investigation of the use of a electric field pulse as diagnostic tool for the longitudinal phase space is given. Furthermore the laser cooling setup and the measured properties of the realized magnetic optical rubidium trap are presented.

Subsequently, the results of the first experiments are discussed in detail. The spontaneous
evolution of a cloud of highly excited atoms, so-called Rydberg atoms, to an ultra cold plasma is observed, by detecting ions and electrons stripped by an electric field. For the first time the spatial profile is measured for ions and electrons originating from the ultra cold plasma.
Chapter 2

Ultra cold plasma

In conventional plasma physics cold plasmas are characterized by an electron temperature of about $10^4$ K ($kT \approx 1$ eV). Because of the fast developments in the field of laser cooling, a totally different regime of plasmas is made possible, the so called ultra cold plasmas (UCP). These plasmas can be created by photo- or near photo-ionization of a laser cooled cloud of atoms [6]. The temperatures of these plasmas are orders of magnitude lower than in conventional cold plasmas, with a ion temperature in the order of 1K and electrons temperature around $10^4$K [8].

In our experiment we will use the cold neutral atoms as a source of electron and ion beams. In this chapter first the basics of the techniques of laser cooling and trapping will be treated. Once a trapped cold gas has been created it can be ionized in two ways: by photo-ionization just above its ionization limit or by field ionization of highly excited states, the so called Rydberg states. For this reason also the basic properties of Rydberg atoms are treated.

2.1 Laser cooling and trapping

In this section the basic principles of Doppler laser cooling and the magnetic optical trap (MOT) are briefly described. After the first demonstration of laser cooling of neutral atoms in 1986 by S. Chu [9], the field of laser cooling matured rapidly. It can now be used as a fairly simple tool to generate a sample of ultra cold neutral atoms, with temperatures below 1 mK. This is often used for high resolution spectroscopy and for a large range of atomic collisions experiments.

2.1.1 Laser cooling

The elementary laser cooling method is so-called Doppler cooling. It relies on the radiation pressure acting on a neutral atom induced by nearly resonant laser light. To simplify the discussion, it will be limited to a one dimensional model of a two-level atom at low laser intensity. Despite its simplicity, it contains most of the features of the more complicated systems.

If a ground state atom absorbs a resonant photon from the laser, the momentum of the photon, given by $\hbar k$, is transferred to the atom. By spontaneous emission the atom falls back to the ground state and emits a photon in an arbitrary direction. Due to this random angular distribution, the momentum transfer of many spontaneous emission events averages to zero, see Fig. 2.1. This results in an average force that is exerted on a stationary atom in the direction of the laser beam. This force is called the scattering force $F_{\text{scatt}}$ and is given by [10]:

\begin{align*}
F_{\text{scatt}} &= (\text{photon momentum}) \times (\text{scattering rate}) \\
&= \hbar k \frac{I}{\hbar \omega} \sigma,
\end{align*}

(2.1) (2.2)
2.1 Laser cooling and trapping

Figure 2.1: An atom experiences a net force in the direction of the laser beam\(^{[10]}\).

with \( I \) the laser intensity, \( \omega \) the laser frequency and \( \sigma \) the effective photon scattering cross section given by:

\[
\sigma = \sigma_0 \frac{1}{1 + (\frac{\omega - \omega_0}{\Gamma/2})^2},
\]

with \( \omega_0 \) the atomic resonant frequency and \( \Gamma \) the natural line width of the transition. The resonant photon scattering cross section \( \sigma_0 \) depends, in the simplest case of a two level system, only on the wavelength:

\[
\sigma_0 = \frac{3\lambda^2}{2\pi}.
\]

By defining the saturation parameter \( s = I/I_0 \) the scattering force can be rewritten to the form:

\[
F_{\text{scatt}} = \hbar k \frac{\Gamma}{2} s \frac{1}{1 + (\frac{\delta}{\Gamma/2})^2},
\]

with \( \delta = \omega - \omega_0 \) the frequency detuning of the laser frequency \( \omega \) in comparison to the atomic resonance frequency \( \omega_0 \), and \( I_0 \) the saturation intensity given by \( I_0 = \frac{\hbar k^2 \sigma}{12\pi} \) and \( k \) the wave number given by \( k = 2\pi/\lambda \).

For a moving atom with velocity \( \vec{v} \), the laser light is Doppler shifted with a frequency shift of \( \delta_d = -k \cdot \vec{v} \). An atom traveling towards the laser beam will experience a blue shift. If the laser is red-detuned from the atomic resonance, it becomes closer to resonance, so the scattering rate will increase. For an atom moving in opposite direction, the same thing happens but in opposite direction, so the scattering rate will drop.

In the case of two monochromatic counter-propagating laser beams, the force on the atom is just the sum of the forces of both laser beams. This is true because we assumed low intensity, so the population of the excited state can be neglected. There is no coupling between both laser beams because of stimulated emission or power broadening. For a stationary atom, the net force will be zero, because both forces cancel each other out, see Fig. 2.2a. But for a moving atom, see Fig. 2.2b, the balance of forces is shifted, so a net force is produced opposite to the velocity, the so-called damping force. Because of this friction an atom experiences, this technique is called an optical molasses. In Fig. 2.3 the force caused by each laser beam is displayed and its sum, given by:

\[
F_{\text{mol}} = \frac{\hbar k \Gamma}{2} \frac{1}{s} \left( \frac{1}{1 + (\frac{\delta + k\vec{v}}{\Gamma/2})^2} - \frac{1}{1 + (\frac{\delta - k\vec{v}}{\Gamma/2})^2} \right).
\]

As can be seen in the figure, for small velocities the net force can be approximated by a constant damping force. For such low velocities, \( kv \ll \Gamma \), the expression of the force of the
Figure 2.2: The principle of an optical molasses is shown in this figure. Two counter propagating red-detuned laser beams generate the force needed for laser cooling. 

a) For a stationary atom, the forces of both lasers cancel each other out, so the net force is zero. 

b) A moving atom, shown in the rest frame, experiences a net force opposite its movement, because the right laser beam comes closer to resonance due to its doppler shift, so the scattering rate increases\cite{10}.

Figure 2.3: For a low intensity optical molasses with red-detuned laser light, the force experienced by an atom, is simply the sum of the scattering forces. For velocities smaller than the Doppler velocity $v_D$, the net force can be approximated by a constant damping force\cite{10}.
2.1 Laser cooling and trapping

molasses $F_{\text{mol}}$ can be reduced to $F_{\text{mol}} = \alpha v$, with $\alpha$ approximated by:

\[
\alpha = 2k \frac{\partial F}{\partial \omega} = 4\hbar k^2 s \frac{-2\delta \Gamma}{[1 + (2\delta / \Gamma)^2]^2}
\]

(2.7)

(2.8)

From this simple model it looks like that the atoms would be cooled to zero temperature. But unfortunately, at very low velocities this model breaks down. The discrete nature of the spontaneous emission process causes a random walk in velocity phase space, as displayed in Fig. 2.4. The net effect averages to zero, but the fluctuations cause a non-zero mean square velocity that causes heating of the atoms. This heating process competes with the cooling term and results in a lowest reachable temperature, the Doppler temperature, which is given by:

\[
T_D = \frac{\hbar \Gamma}{2k_B}
\]

(2.9)

and a corresponding velocity, the Doppler velocity $v_D$ given by:

\[
v_D = \sqrt{\frac{\hbar \Gamma}{2m}}.
\]

(2.10)

By using three pairs of orthogonal laser beams, a three dimensional optical molasses can be created. In the intersection of these six beams, atoms will be cooled. But because of the non-zero final velocity, the atoms will diffuse out of the intersection area and escape.

2.1.2 The magneto-optical trap (MOT)

To create a trap, not only a velocity dependent damping force is needed, but also a spatial restoring force. This can be done by adding a magnetic field gradient to an optical molasses. If the polarizations are chosen carefully, the magnetic field can cause an imbalance in the optical scattering force in such way it creates a restoring force.
Figure 2.5: The basic working principle of a Magneto-Optical trap ($J = 0$ to $J = 1$ transition) is shown. Because of the field gradient and the resulting linear Zeeman splitting, the sub level splitting is position dependent. When the atom is illuminated by a pair of red-detuned counter-propagating laser beams, an imbalance in the scattering force occurs, because one laser can only excite to one sub-level due to the selection rules for the transitions\cite{10}.

In Fig. 2.5 the working principle of a MOT is explained, for again a two state system with ground level $J = 0$ and excited state $J = 1$. The magnetic field induces a linear Zeeman shift near the trap center, so the excited level is splitted in three magnetic sub-levels ($M_J = 0, \pm 1$). For simplicity the $z$-axes is used as quantization axes, instead of the magnetic field vector itself. The two counter-propagating laser beams have to be circular polarized, one beam $\sigma^+$, the other $\sigma^-$. The transition selection rules state that $\sigma^+$ light only can drive $\Delta M_J = +1$ transitions, and $\sigma^-$ light only $\Delta M_J = -1$. On the left side of the center of the trap $M_J = +1$ is closer to resonance and on the right side $M_J = -1$. This causes the extra imbalance that will function as a restoring force.

This force can be derived in the same way as the force in the optical molasses, by also adding the Zeeman shift in Eq. (2.6), so:

$$F_{MOT} = \hbar k \frac{\Gamma}{2} \beta \left( \frac{1}{1 + \left( \frac{\delta + kv + \beta_\omega}{k} \right)^2} - \frac{1}{1 + \left( \frac{\delta - kv - \beta_\omega}{k} \right)^2} \right),$$

with $\beta$ the restoring force coefficient defined as:

$$\beta = \frac{g \mu_B}{\hbar} \frac{dB}{dz},$$

with $g$ the Landé-factor, $\mu_B$ the Bohr magneton and $\frac{dB}{dz}$ the magnetic field gradient in $z$-direction.

The motion of an atom near the center of the trap, has a Zeeman shift that is small compared to the detuning, so it can be looked at as a damped-harmonic oscillator with a force approximated by:

$$F_{MOT} = -\alpha v - \frac{\alpha \beta}{k} z$$
with $\alpha$ the dampening coefficient introduced for the optical molasses.

![Diagram](image)

**Figure 2.6:** (a) The gradient needed for the MOT can be generated by a pair of coils in anti-Helmholtz configuration[10]. (b) On axes this result in a linear field gradient.

This one-dimensional MOT can be expanded to three-dimensions by using a quadropole magnetic field, generated by two coils in anti-Helmholtz configuration, see Fig. 2.6 and by choosing the right polarizations for all the three pair of perpendicular laser beams.

From literature, the most common reported MOT has a trap depth of about 1 K, contains up to $10^9$ particles at a density of about $10^{10}$ cm$^{-3}$ and a temperature of 300 $\mu$K[11]. The largest MOT produced so far to our knowledge, in both number of atoms as well in density, is reported by Ketterle et al. They succeeded in creating a MOT containing $10^{10}$ atoms at a density of $10^{12}$ cm$^{-3}$ with a temperature of approximately 1.2 mK[12].

## 2.2 Rydberg states

A Rydberg atom is an atom with a single valence electron in a high principle quantum number $n$ state[13][14]. In this section a short introduction to the properties of Rydberg atoms will be given.

Due to the high excited state, the classical radius of the valence electron orbit is large compared to the size of the ionic core. The core electrons shield the valence electron from the nucleus, so it only experiences a nucleus with an effective charge of a single proton, like in a highly excited hydrogen atom. Therefore a Rydberg atom also behaves very much like a hydrogen atom. The energy levels of high Rydberg states are thus also very close to the levels of hydrogen, given by the Rydberg formula:

$$E_n = -\frac{R}{n^2}, \quad (2.14)$$
with $R$ the Rydberg constant of $R \approx 13.6$ eV. In Fig. 2.7 some high excited states of sodium are compared with the hydrogen levels and indeed agree well.

With use of perturbation theory a correction on the hydrogen energy levels can be calculated for non-hydrogen atoms in Rydberg states. For high angular moment $l$-states, the valence electron is always far from the ionic core, so only a small correction is needed. For lower $l$-states, the correction is larger because it can penetrate and polarize the ionic core. This correction term is called the quantum defect $\delta_l$. With this correction term, the original hydrogen Rydberg energy formula is changed to

$$E = \frac{-R}{(n - \delta_l)^2} \quad (2.15)$$

with $\delta_l$ the quantum defect for a particular $l$-state.

Except this correction term, the whole concept of high excited Rydberg atoms can be understood from the simple classical Bohr model. The radius of the charge distribution of the valence electron scales as $n^2$. So high Rydberg states can have very large radii, for example $n = 50$ has a radius of 2500 times that of ground state Hydrogen atom! Also the lifetime of Rydberg states have a totally different timescale, they scale, same as hydrogen, as $n^3$ for high $n$.

In Rydberg states the excited electron is in a large, loosely bound orbit and the dipole moment scales with the radius. Therefore these Rydberg states are, in contrast with their ground states, very sensitive to external electric fields. Classically, an external electric field lowers the coulomb potential of the atom on one side, see Fig. 2.8. Classical field ionization occurs whenever the energy of a bound state exceeds the energy of the saddle point of the potential[15], so it can escape. The classical ionization limit $E_c$ is given for azimuthal angular momentum quantum number $m = 0$ by:

$$E_c = \frac{1}{16} \frac{1}{n^4}. \quad (2.16)$$

For states with higher $m$, a slightly higher field is necessary, because in those states the electron is kept away from the saddle point by the centrifugal barrier.

Highly excited Rydberg atoms can be ionized with very low fields, for example only $65.9$ V/cm for $n = 47$, compared with impractical large field of $3.2 \times 10^8$ V/cm for the ground state of hydrogen[16].

These field ionization properties can be used to detect the different Rydberg states. In most cases this is done by applying a time dependent electric field, at the timescale of $\mu$s[13]. A particular Rydberg state has its precisely defined ionization field, and thus a precisely defined ionization point in time. By measuring and analyzing the time dependent signal from the electrons or ions extracted from the Rydberg gas, the initial state distribution can be measured.

### 2.3 Ionization and creation of a UCP

An UCP can be created by ionization of a cloud of cold atoms. There are two main ways to accomplish that. The most direct way, is to photo-ionize the cold atoms just above the ionization threshold, see Fig. 2.9b. The excess energy, the difference in energy between the photon and the ionization threshold, is transferred to kinetic energy. Due to the large difference in mass, most energy is transferred to the electrons and not to the ions. A lower excess energy than given by the laser bandwidth is not possible, so this limits the initial temperature of the electrons. If needed a higher initial temperature is possible by adjusting the laser wavelength.

When the atoms are ionized, the electrons have much larger velocities than the ions. So some electrons will escape from the ionized cloud, because they are not trapped as long the total charge is zero. But as more and more electrons escape, a potential well is created that
Figure 2.7: Energy level diagram for H and Na. The Na Rydberg levels of high l-states are very close to the Hydrogen levels[10].
Ultra cold plasma

Figure 2.8: Schematic overview of classical field ionization. The original coulomb potential of an atom (dashed line) is lowered by an applied external electric field (solid line).

Figure 2.9: There are two closely related ways to create a UCP. a) Creation of a UCP by exciting a Rydberg level that evolves, under the right circumstances, spontaneously to a UCP. b) Creation of a UCP by photo-ionizing over the threshold. The excess energy is mostly transferred to the kinetic energy of the electrons, because of the much larger mass of the ions[17].
2.3 Ionization and creation of a UCP

![Potential Energy](image)

Figure 2.10: Initially neutral atoms \(A\) are ionized. Because there exists no correlation in the position of the atoms, the generated potential landscape after ionization is very erratic. This potential energy is converted to kinetic energy of the fast moving electrons\[17\].

starts to trap the remaining electrons. If this happens a UCP is born. The ions are of course not fixed in space, so the cloud will expand due to initial velocities and the repelling Coulomb force. But because of the large difference in mass, this timescales is much slower than of the timescale of electron motion that cause the creation of the plasma.

The trapped atoms in the MOT are distributed randomly. There exists no spatial correlation between the atoms, so the created potential landscape in the plasma is very irregular. The electrons and ions are on average not at a potential energy minimum. This extra potential energy will be rapidly converted to kinetic energy, and thus heating the plasma. This intrinsic heating effect is called correlation heating. This heating is the fastest for the electrons at a timescale corresponding to its inverse plasma frequency \(1/\omega_p = \sqrt{m_0/\epsilon_0 n_e e^2}\), with \(m\) the electron mass, \(\epsilon_0\) the electric permittivity of free space and the electron density \(n_e\). This heating continues until the equilibrium temperature is obtained of \(kT \approx \frac{e^2}{4\pi\epsilon_0 n_e^2}\). At typical parameters with \(n_e \approx 10^9\) cm\(^{-3}\), a plasma of 10 mK is heated to about 10 K on a timescale of ns.

An other option to create a plasma is illustrated in Fig. 2.9a. Instead of tuning the pulsed laser above the ionization threshold, it is tuned just below and used to excite the atoms to a high level Rydberg state. It was shown by M.P. Robinson et al\[18\], that a cloud of Rydberg atoms spontaneously evolve in an UCP under the right circumstances.

This process happens in a timescale of \(\mu s\), in which the position of most of the atoms in the Rydberg gas can still be considered to be frozen\[19\]. Both options create a comparable UCP, but the precise dynamics of this method are quite complex and totally different. In a cloud of Rydberg atoms, some atoms are always ionized by, for example, black body photo-ionization or dipole-dipole interactions\[18\]. These initial electrons cause an avalanche effect that ionizes about 70% of the Rydberg atoms, and transfer 30% of the atoms to lower lying Rydberg states.
Chapter 3

Ultra cold electron and ion beams

3.1 Brightness and emittance

The properties and resulting dynamics of particle beams and bunches is quite complex. To make it possible to compare different particles beams, it is important to have some quantity to express the average beam quality, a measure for the potential usefulness in applications. For this purpose a local quantity, the 6D brightness, can be defined as the current density per unit solid angle and per unit energy spread:

\[ B = \frac{\partial^3 I}{\partial A \partial \Omega \partial U}. \]  

(3.1)

The average of this property over the whole bunch, the so called average 6D brightness \( \bar{B} \) can be used as a single comprehensive quantity, given by:

\[ \bar{B} = \frac{I}{\Delta A \Delta \Omega \Delta U}. \]  

(3.2)

Here \( I \) is the beam current, \( \Delta A \) the surface area of the beam cross section, \( \Delta \Omega \) the solid angle of the beam and \( \Delta U \) the beam energy spread. In Fig. 3.1 an illustration is given of these properties for a focussed electron beam. The quantity \( \bar{B} \) is independent of the beam size or divergence, but it depends on the beam energy; when the averaged energy increases, the beam divergence will automatically decrease, because of the higher velocity in the beam direction, which leads to an increased brightness.

To enable comparison of beams of different energies, a Lorentz-invariant definition of brightness can be used, the so called normalized 6D brightness[4]:

\[ B_n(\vec{r}, \vec{p}) \equiv e m^2 c^2 f(\vec{r}, \vec{p}), \]  

(3.3)

Figure 3.1: An illustration of an electron beam focus with \( \Delta A \) and \( \Delta \Omega \) the solid angle.
with \( f(\vec{r}, \vec{p}) \) the local phase space density, normalized to the number of particles \( N \) in the bunch, \( e \) the elementary charge and \( m \) the mass of the particle.

In most cases the exact distribution function \( f(\vec{r}, \vec{p}) \) is not known, but a model can be used. The simplest, to some extent realistic model for a particle bunch is a 6D Gaussian distribution in phase space with its axes aligned along the \( x, y, \) and \( z \) coordinate axes. This distribution is given by:

\[
f(\vec{r}, \vec{p}) = f(0, 0) \exp \left( -\frac{x^2}{2\sigma_x^2} - \frac{y^2}{2\sigma_y^2} - \frac{z^2}{2\sigma_z^2} - \frac{p_x^2}{2\sigma_{px}^2} - \frac{p_y^2}{2\sigma_{py}^2} - \frac{p_z^2}{2\sigma_{pz}^2} \right).
\] (3.4)

For this uncorrelated distribution the peak normalized 6D brightness \( \hat{B}_n \), the maximum local phase space density at \( B_n(0,0) \), can be expressed as:

\[
\hat{B}_n = \frac{mc^2}{(2\pi)^3} \frac{Ne}{\sigma_x \sigma_{px} \sigma_y \sigma_{py} \sigma_z \sigma_{pz}}.
\] (3.5)

For a thermal source, the momentum distribution follows a Boltzmann distribution, so \( \sigma_{px}^2 = \sigma_{py}^2 = \sigma_{pz}^2 = mkT \). In that specific case, the peak normalized 6D brightness can be calculated as:

\[
\hat{B}_n = \frac{n_0e}{mc} \left( \frac{mc^2}{2\pi kT} \right)^{3/2},
\] (3.6)

with \( n_0 \) the initial density, \( T \) the effective particle temperature and \( k \) the Boltzmann constant. This equation shows the essence of the project. The normalized brightness can be improved not only by increasing the source density but also decreasing the initial temperature.

In practice there are in many cases linear correlations in the phase space distribution, for example position-momentum correlations due to focusing or expansion of the particle beam. In that case Eq. (3.5) would give a significant underestimation of the actual peak brightness, because the product of the correlated root-mean square (rms) components can be much larger than the rms value of the product itself. The three directions, \( x, y \) and \( z \) are in most cases in good approximation decoupled, so this overestimation can be solved by the introduction of the normalized root mean-square emittances \( \epsilon_x, \epsilon_y, \) and \( \epsilon_z \). In the \( x \)-direction it is defined as[4]:

\[
\epsilon_x = \frac{1}{mc} \sqrt{<x^2> <p_x^2> - <xp_x>^2},
\] (3.7)

where the brackets \(< >\) mean averaging over all particles in the bunch. The other emittances \( \epsilon_y \) and \( \epsilon_z \) are defined in the same manner. The emittance can be understood as the area in the phase-space plane of the minimal ellipse enclosing the distribution. This is illustrated in see Fig. 3.2. If \( \vec{p} = \gamma m \vec{v} \) is used then it is a Lorentz-invariant quantity, i.e. energy independent. The emittance is a measure for the focusability of the beam in one particular direction.

With use of these three emittances and the total charge of a bunch, a practical estimate for the peak normalized 6D brightness can be given as:

\[
\hat{B}_n \approx \frac{1}{mc} \frac{Ne}{(2\pi)^3 \epsilon_x \epsilon_y \epsilon_z}.
\] (3.8)

For many applications however, the peak current and the transverse normalized emittances are the most important parameters, while the longitudinal energy spread is not so critical. For this reason the so-called transverse brightness \( B_\perp \) is often used to express bunch quality, which is defined by [20]:

\[
B_\perp \equiv \frac{i}{4\pi^2 \epsilon_x \epsilon_y},
\] (3.9)
Figure 3.2: The projections of two bunches on the x-p_x phase space plane and the enclosing ellipsoids. One bunch has a low emittance and the other a much larger emittance, because of the curvature, while the actual surface areas are comparable.

where \( \hat{J} \) is the peak current. In the absence of position-momentum correlations in the longitudinal-direction, the transverse brightness \( B_\perp \) is simply proportional to the product of the peak normalized 6D brightness \( \hat{B}_n \) and the rms energy spread \( \sigma_u \):

\[
B_\perp = \sigma_U \sqrt{2\pi} \hat{B}_n, \tag{3.10}
\]

where \( \sigma_U = \frac{e\hat{J}}{m} \sigma_{p_z} \) with \( p_{z0} \) the average momentum of the bunch.

It is easy to show that for a thermal electron source with peak current density \( \hat{J} \), Eq. (3.9) can be rewritten to[4]:

\[
B_\perp = \frac{mc^2 \hat{J}}{2\pi kT}, \tag{3.11}
\]

with \( T \) the effective electron temperature. This equation also illustrates the basic concept of the project. The transverse brightness is inverse proportional to the initial source temperature \( T \), so improving the brightness can be done by decreasing the initial temperature.

### 3.2 Current state of the art sources

Electron sources can be divided into two main categories: continuous and pulsed sources. High-Brightness continuous sources are mainly used for electron microscopy. The conventional state of the art sources that are used for that purpose are the Cold Field Emission Guns (CFEG). The current highest experimentally measured brightness, both normalized and transverse brightness, is produced by recently developed carbon nano tubes (CNT) field-emitters[21]. CNTs can produce up to 1 \( \mu \)A of continuous current from a few nm\(^2\) source area, with a transverse brightness of about \( 10^{15} \) A m\(^{-2}\) rad\(^{-2}\). This is about an order of magnitude below its thermal brightness limit, if an effective source temperature of \( kT \approx 0.5 \) eV is assumed.

For applications that need higher peak currents, such as for example hard X-ray free electron lasers (XFELs)[2], pulsed electron sources have been developed. The best performing high current pulsed picosecond electron sources are radio-frequency (rf) photoguns. In these devices electron bunches are created by pulsed laser photoemission from a metal cathode. The electrons are subsequently accelerated in strong rf fields of, typically, 100 MV/m.
The best performing rf photogun at present, the photogun of the Accelerator Test Facility at Brookhaven National Lab, can produce about 0.5 nC electron bunches with \( I \approx 120 \) A and \( \epsilon_z \approx 0.8 \) \( \mu \)m, corresponding to a beam brightness \( B_1 = 5 \times 10^{12} \text{A} / (\text{rad}^2 \text{m}^2) \) [20], about an order of magnitude smaller than the thermal transverse brightness limit given by Eq. (3.11). The brightness achieved is limited by nonlinear space charge forces. Recently, it was shown that the unwanted effects of space charge forces can be almost completely eliminated by proper shaping of the radial intensity profile of the femtosecond photo-excitation laser [22]. This would make it possible to reach the thermal brightness limit corresponding to \( T = 10^3 - 10^4 \) K.

### 3.3 Ultra cold electron bunches

All current strategies for improving the brightness of photo-emission based electron sources have in common that they try to increase the current density at the source. This is done by making the source area very small, but by decreasing the source area the problem of non-linear space charge fields become larger, so the current has to be decreased to retain the low emittances.

In all those cases the angular spread, or the effective electron temperature of the source, is not changed. In this project, the ultra cold bunches project [5], a totally different approach is used. In stead of increasing the current density, the temperature of the source is decreased.

This new source is similar to an rf photogun in the sense that it will produces bunches with a comparable bunch charge from a comparable source area. But instead of a metal cathode, a laser cooled cloud of rubidium atoms, trapped in a MOT is used as source of the electrons. The temperature of the electrons in the plasma created by just above the threshold photoionization of this cold Rubidium cloud, can be as low as a mK. But due to correlation heating, as described in Sec. 2.3, the temperature of the electrons will increase to about 10 K on a ns-timescale. But this is still about 2-3 orders better in temperature than electrons generated from a metal cathode! So a large improvement in brightness should be possible.

#### 3.3.1 Method

The generation of an electron bunch from a cold cloud of atoms can be done in a four-step procedure, schematically illustrated in Fig. 3.3. A cold cloud of rubidium atoms is trapped in a Magneto-Optical Trap in a volume of a about 1 mm\(^3\) with densities up to \( 10^{18} \text{m}^{-3} \) [12] and a temperature below the 1 mK.

In the next step a fraction of the cold atom cloud is excited to an intermediate excited state with use of an \( \mu \)s laser pulse. It is possible to simply use the MOT trapping laser beams, but also a special shaped laser beam so the shape of the cloud of excited atoms can be changed.

After that, a laser pulse perpendicular to the excitation laser is given. There are two possibilities by changing the wavelength of this laser pulse. It is possible to do just above threshold ionization of the excited atoms, so a UCP is directly formed in this volume [6]. Or just below the ionization threshold so an cloud of highly excited Rydberg atoms is created in the volume. As described in Sec. 2.2 the Rydberg atoms will interact and also form an ultra cold plasma after some delay of several \( \mu \)s.

In this way up to 100 pC, about \( 5 \times 10^8 \) electrons, can be created in a volume of \(~ 1 \text{ mm}^3\) . For a MOT a loading rate of over \( 1 \times 10^{11} \) atoms/s is possible [12], so a total charge up to 10 nC may be extracted per second.

In the last step, the electrons are extracted out of the plasma to form an electron bunch. The electric field has to be at least an order of magnitude stronger than the minimum required field to pull the electrons and ions in the plasma apart. For a 1-mm-sized 1 pC bunch this typically means that a voltage of 10 kV has to be applied across a 1 cm gap. This field also has to be
Figure 3.3: Schematic of the four-step procedure to realize a pulsed UCP electron source[5].

switched on extremely rapidly, sub nanosecond, to prevent emittance growth by space-charge forces induced by nearby ions during acceleration.

In the case of below threshold excitation, the electric field pulse of the last step can be given before there is time to form an UCP. So an electron bunch may be created directly from field ionizing the Rydberg state. In this case the heating processes in the plasma, such as correlation heating, might be avoided, so even lower temperature bunches might be possible. To investigate if this is really possible, more study has to be done on the exact dynamics and timescales of field ionization of Rydberg atoms.

In both schemes is the initial electron density proportional to the product of the intensities of the excitation and the ionization laser beams. This offers an excellent opportunity to control the initial charge distribution by adjusting the profile of both lasers. The unwanted effects of non-linear space-charge forces may be almost completely eliminated by proper shaping of the initial electron density. This can be done by choosing the initial conditions such that the bunch automatically evolves in a waterbag, a homogenous filled ellipsoidal bunch, that only has linear space-charge forces[22]. If there are only linear forces, there is no emittance growth and the expansion can be reversed by simple optics.

There are two interesting initial charge distributions possible that have this behavior, a pancake with a half-circle radial initial charge density profile and a cigar like bunch with a initial parabolic longitudinal charge density distribution. Using Eq. (3.7) it can be shown that the normalized rms emittance of both bunches with radius $R$ is given by

$$
\epsilon = R\sqrt{kT/5mc^2}.
$$

Both configurations have there own specific advantages as will be seen in the next sections.
3.3.2 Velocity bunching

An extra advantage of this setup is, that the extracted bunches are automatically compressed in drift space after they are accelerated. This happens because electrons in the back of the bunch experience a larger acceleration potential than those in the front. So the electron in the back of the bunch have a larger velocity than electrons in the front, so they will catch up. This "velocity bunching" leads to sub-ps bunch lengths, without the need for a magnetic compressor or special compression rf-cavity.

3.3.3 GPT simulations

To get a better insight in the characteristics and abilities of the ultra cold electron source, numerical simulations have been performed, and the results are published[6]. The simulations are performed with the General Particle Tracer(GPT) code[23]. In this code, the equations of motion are solved for a set of particles in a given external electric magnetic field with taking the internal space-charge forces into account.

A very basic cold atom setup is assumed in the simulations. It consist of a cylindrical symmetric accelerator structure of two thin parallel conductive plates, with both a circular hole so electrons can exit the structure. A voltage pulse of 1 MV and a risetime of $\tau_r = 150$ ps is used. Simulations are done for a pancake as a cigar bunch.

The velocity bunching effect for the cigar bunch is the most pronounced. At a position about 2 cm outside the diode a rms bunch length of only $\sigma_z = 20$ fs is achieved, that corresponds to a peak current $I > 25$ A. The lateral brightness at this point is $B_\perp \geq 5 \times 10^{14}$ A/(rad$^2$ m$^2$), comparable to CNT performance. The pancake bunch compresses almost immediately to a bunch length of $\sigma_z = 150$ fs corresponding to $I = 25$ A and stays almost constant over many cm in its trajectory. This behavior is due to a balance between space charge forces and the smaller velocity bunching, because of the small longitudinal velocity differences. The lateral brightness of this configuration is $B_\perp = 5 \times 10^{13}$ A/(rad$^2$ m$^2$), an order of magnitude higher than state-of-the-art rf photogun performance!

From these simulations, the pancake bunches can be characterized by a high charge, a small energy spread and robust, stable behavior. In the other hand, cigar bunches offer the highest peak brightness and shortest bunch lengths at a specific position, but at costs of energy spread. The achieved brightness values in the simulation are still 1 – 2 orders of magnitude removed from the thermal limit. So in principle even higher brightness values may be possible by for example optimizing the accelerating structure, so the emittance growth by non-linear fields will be reduced.

3.3.4 Ion beams

The ideas described in this section can also be used to produce very high brightness continuous and pulsed ion beams. In principle only the polarity of the voltage pulse needs to be inverted to change the electron source to an ion source. Because of the large difference in mass, the timescales of the ions are totally different.

For example the correlation heating is orders of magnitude slower for ions. This may make the production of even colder ion beams possible and easier compared with electrons. The technical problem of producing sub-nanosecond electric field pulses needed for acceleration is avoided. But the precise possibilities and limitations of such ultra cold pulsed ion sources still have to be investigated. The same thing also has to be done for a continuous setup, where for example an ion beam is extracted from a MOT or directly from a slow atomic beam with use of a constant accelerating field and ionization laser.
Chapter 4

Beam diagnostics

In this chapter a short overview is given of the most basic diagnostics that are used or can be used to measure the properties of the ultra cold electron and ion beams. In general this means mapping the phase space of the bunch. For most cases it is sufficient to measure the projection of the 6D phase space distribution on the 2D $x$-$p_x$ and $y$-$p_y$ transverse phase spaces and the projection on the 2D $z$-$p_z$ longitudinal phase space separately. The focus of this chapter is in particular on the more difficult task of measuring the longitudinal phase space distribution.

4.1 Charge

The most basic diagnostic tool is a Faraday cup, which simply consists of a conductive surface that collects the charged particles. The current will flow through a connected wire to a current amplifier, so the beam current can be measured. Unfortunately it is not possible to measure the time dependent current electronically at a timescale of picoseconds, so for a pulsed source with (sub)-picosecond bunches no data can be collected of the longitudinal profile. But the total charge per bunch can be obtained by integrating the signal after first collecting the charge in some capacitor.

The exact electronic amplification factor is normally known, so this device is often used for charge calibration. For higher energies the basic principle is still the same, but a more complicated design is necessary to prevent that electrons produced by secondary emission do not escape.[24]. At the energies we aim at in the UCP project (30 keV) the secondary electron emission rate is so low, that it causes no problem.

4.2 Beam profile

To observe the spatial distribution in the lateral direction a phosphorescent screen can be used. The screen is placed in vacuum, where electrons entering the screen excite the molecules. In turn the molecules emit fluorescence light, that can be imaged outside the system with a CCD camera. If a higher sensitivity is needed, an image intensifier in the form of a multi channel plate (MCP) can be mounted in front of the phosphor screen.

A commonly used phosphorescent material is ZnCdS:Ag (type P20)[25]. Electrons entering the screen excite the molecules, that decay and emit a photon. The emission spectrum from the phosphor ranges from 470 nm to 670 nm with a maximum at 550 nm (yellow/green)[24]. The fluorescence decays in about 1 ms to 1%. The efficiency is about 0.06 emitted photons per eV, but depends on the grain size, thickness, electron energy and the used conductive layer[25]. This thin conductive layer of aluminium or indium tin oxide (ITO) is added to prevent the build-up
of charge. Aluminium has the advantage that it increases the efficiency by reflecting back the emitted photons in the backward direction to the CCD camera, but only electrons with energies above 3 keV can penetrate through the layer. In addition it is possible to use this conductive layer as an Faraday cup.

To increase the detection sensitivity, a multi channel plate (MCP) can be added. This device consists of an array of millions of very thin conductively coated glass capillaries with a diameter of several μm, see Fig. 4.1a for a schematic illustration. A voltage difference is applied between the front and the back side of the plate, which makes each channel function as an independent secondary-electron multiplier tube: an incoming electron or ion generates some secondary electrons when it hits the wall of the channel that are accelerated, and in turn generate more electrons, so an avalanche is started. At the output a large number of electrons are released, see Fig. 4.1b. These electrons can be detected with a Faraday cup or, to have spatial resolution, with a phosphor screen.

![Micro channel plate and Single MCP channel](image)

**Figure 4.1: An schematic overview of an MCP detector[26].**

For our first experiments, a Burle MCP detector mounted on a P20 phosphor screen has been used. This detector has an effective detection area with a diameter of 40 mm. The MCP channels have a diameter of 10 μm and a length of 400 μm. Three independent voltages can be applied, on the front of the MCP, the back of the MCP, and on the conductive layer of the phosphor screen. An extra grounded metal grid, with a filling factor of about 50%, is placed in front of the detector to shield the rest of the setup from the electric field applied on the front of the MCP.

At normal operations 5 kV is applied to the phosphor screen. Depending on whether electrons or ions have to be detected, a positive or negative voltage is applied to the front of the MCP. A voltage of about 900 V is applied over the MCP channels which determines the gain factor. A current amplifier is placed between the high voltage supply and the phosphor screen, so the conductive layer on the phosphor screen can be used as a Faraday cup measuring time dependent signals up to several ns.

The gain of the MCP detector should be calibrated with a true Faraday cup, which has not been done yet. To make an estimate of the charge we assume the maximum specified gain of $10^4$. The measured charges are therefore a lower bound for the actual charge.

### 4.3 Transverse phase space

More information of the bunch can be obtained by mapping the transverse $x$-$p_x$ and $y$-$p_y$ phase space. The simplest way to do this, is to measure the local angular spread at every position by...
selecting a small portion of the bunch by a moveable slit, see Fig. 4.2. The local angular spread
is simply the transverse momentum spread divided by the $z$-momentum $p_z$.
By moving the position of the slit, one can measure the whole transverse phase space, so an
distribution as in Fig. 3.2 can be reconstructed. This information can be used to calculate the
emittance of the bunch.

4.4 Longitudinal phase space

In contrast with most conventional high-brightness electron sources like rf-photoguns, there is no
synchronized femtosecond laser or synchronized rf-pulse available in this project. It is technically
very hard to synchronize an rf-pulse or femtosecond laser pulse with our acceleration voltage
pulse with a time jitter lower than, say, 1 ps. Commonly used methods such as electric-optical
detection[28] or kicker cavities[29][30] are therefore not possible to measure the longitudinal
charge distribution.
Fortunately a perfectly synchronized high voltage pulse is available, the accelerating pulse
itself. In this section a simple analytical model will be constructed to make a first estimate of the
properties and capabilities of a streaking setup with use of this pulse. The proposed streaking
setup consists of two parts, a set of streaking electrodes and a bending magnet, see Fig. 4.3.
With those two parts, it can map the longitudinal phase space $z-p_z$ of the bunch on the $x - y$
plane of the detector.

When the bunch passes through the streaking electrode structure, a fast electric field pulse
is given. Particles that enter with some delay in respect to the reference particle, experience a
different average field so they get an different deflection angle. This deflection angle translates,
through drift space, in a $x$-position on the detector, so in principle the initial $z_0$ phase space is
mapped on the $x$-axes of the detector.
After the electrodes, the particles enter a bending magnet. For simplicity the z-direction is chosen to be in the forward direction of the reference particle. Because of dispersion in the bending magnet, it images particles with different z-momenta at different y-positions on the detector. Therefore it translates the initial $p_z$ phase space to a distribution on the y-axis of the detector. So the two parts in the system together, generate an image of the initial phase space on the detector.

### 4.4.1 Streaking setup

**Electrode setup**

For an analytical model a simple electrode geometry is used, consisting of two rectangular metal plates parallel to the y-z-plane and separated by a distance $h$ in the x-direction, see Fig. 4.4. The particle will travel in the z-direction between the plates, which have a length $L$. The width is much larger than the bunch radius, so the field can be assumed homogenous.

The voltage over the electrodes is modeled as linear ramp, with a slope parameterized by the rise time $\tau$ and the maximum voltage $V_{\text{max}}$. In the ideal case there are no fringe fields, so the electric fields are given by a step function in space. In that case the homogenous field between the plates is given by:

$$E(t) = E_0 \frac{t}{\tau} + E_{\text{off}}. \quad (4.1)$$

The extra offset term $E_{\text{off}}$ is chosen such that the average field experienced by a reference particle through the electrodes is zero and the field strength $E_0$ is defined as $E_0 = \frac{V_{\text{max}}}{h}$. 

---

**Figure 4.3:** Schematic overview of the proposed streaking setup with additional bending magnet.
Induced magnetic field

Because of the linear time dependence of the electric field, also a constant magnetic field will be present. From the Maxwell equations it can be directly obtained that:

$$\frac{cB}{E} \sim \frac{L}{c\tau},$$  \hspace{1cm} (4.2)

with \( L \) the characteristic length and \( \tau \) the characteristic timescale of the system and \( c \) the velocity of light. From this equation the importance of the magnetic field can be expressed by the ratio of the maximum electric and the magnetic force:

$$\frac{F_B}{F_E} \approx \frac{1}{2} \frac{v_0 L}{c^2 \tau}.$$  \hspace{1cm} (4.3)

For typical parameter values this is much smaller than 5%. Furthermore the magnetic motion cancels out in first order, if the streaking by the electric field is ignored, because of the reversal in sign in the first and second part between the electrodes. Therefore the magnetic field can be neglected in the analytical model.

For relativistic velocities, the Lorentz force will become comparable in strength to the Coulomb forces of the electric field and thus cannot be neglected anymore. Therefore it is chosen to use simple Newtonian mechanics for the model, because without description of the magnetic field effects the model cannot be used for relativistic particles anyhow.

Electron trajectories

For simplicity the effects of space-charge are neglected in the model and only the behavior of single particles are calculated. An electron exiting an accelerator diode traveling in the \( z \)-direction will pass between the two electrodes plates with an electric field given by Eq. (4.1). The change in momentum in the \( x \)-direction due to the Coulomb force is given by:

$$\frac{dp_x}{dt} = qE(t).$$  \hspace{1cm} (4.4)

At \( t = t_0 \) the particle enters the electrode structure at \( z = 0 \). For the reference particle it is chosen that \( t_0 = 0 \). The relative time of the particle between the electrodes is thus defined as \( t' = t - t_0 \). So the \( x \)-momentum can be calculated as:

$$p_x(t_0, t') = \int_{t_0}^{t_0 + t'} qE(t)dt$$

$$= q \left( \frac{1}{2} E_0 \tau t'^2 + E_0 \tau t' + E_0 \tau t_0 t' \right).$$  \hspace{1cm} (4.5)
The $x$-position is then given by:

$$x(t_0, t') = \int_0^{t'} \frac{1}{m} p_x(t_0, t'') dt'' \tag{4.7}$$

$$= \frac{q}{m} \left( \frac{1}{6} E_0 \tau \rho^2 + \frac{1}{2} (E_{off} + E_0 \tau_0) \tau^2 \right) \tag{4.8}$$

The time it takes for a reference particle with $p_z = p_{z0}$ to pass the electrode structure is $\Delta t_{L0} = \frac{L_{str} m}{p_{z0}}$, so $E_{off}$ is defined as:

$$E_{off} = \frac{1}{2} E_0 \frac{\Delta t_{L0}}{\tau} \tag{4.9}$$

$$= \frac{1}{2} E_0 L_{str} \frac{m}{p_{z0}} \frac{1}{\tau} \tag{4.10}$$

which cancels out the momentum increase for the reference particle after passing the electrode structure. The $x$-displacement is not canceled out, because only the momentum change is zero, not the averaged momentum itself. In Fig. 4.5 the trajectories of three particles between the electrodes with three different starting delays are plotted.

**Deflection**

The $x$-momentum of an electron after leaving the streaking electrodes, and to a lesser degree its position, determine the characteristics of the image that will be formed. The length $L_{str}$ is chosen such that the reference particle spends a time equal to the switching time $\tau$ between the electrodes, so a fair comparison is possible between different situations with different energies or switching times:

$$L_{str} = \frac{p_{z0}}{m} \tau. \tag{4.11}$$

For particles with $p_z = p_{z0}(1 + \Delta)$ the $x$-momentum after the electrodes is calculated as:

$$p_{zL} = -q E_0 \left( \frac{t_0}{(1 + \Delta)} - \frac{\tau \Delta}{2(1 + \Delta)^2} \right), \tag{4.12}$$

where $\Delta$ is the relative momentum deviation from the reference particle. The change in $x$-position is given in the same way by:

$$x_L = q E_0 \frac{1}{m} \frac{1}{2} \left( \frac{t_0}{(1 + \Delta)^2} - \frac{\tau t_0}{(1 + \Delta)^3} \right) \tag{4.13}$$

$$\approx q E_0 \frac{1}{m} \frac{1}{2} \left( \frac{1}{\tau^2} - \frac{\tau t_0}{(1 + \Delta)^3} \right) \tag{4.14}$$

For imaging purposes, only the deflection slope is important. The slope $\theta_x$ is given by:

$$\theta_x = \frac{p_{zL}}{p_z} = -q E_0 \frac{1}{p_{z0}} \left( \frac{t_0}{(1 + \Delta)^2} - \frac{\tau \Delta}{2(1 + \Delta)^3} \right). \tag{4.15}$$

For the case there is no spread in $z$-momentum ($\Delta = 0$), this simplifies to:

$$\theta_x = -q E_0 \frac{1}{p_{z0}} t_0. \tag{4.16}$$

This equation shows the essential behavior. The deflection slope is linear with the time delay and the magnitude of this effect does not depend on the switching time and only on the maximum achievable field. This is a direct consequence of setting the electrode length by Eq. (4.11).
Figure 4.5: Plot of trajectories of an electron between the electrodes with a separation of 10 mm, a voltage pulse of 30 kV with a rise time of 200 ps. The middle line is the reference electron. The other two are for electrons with a relative delay of 10 ps and -10 ps. a) The x-momentum as function of z. b) The x-position as function of z.
Figure 4.6: The x-deflection is plotted as function of the delay time for the reference electron (middle line) and two particles with ±1% momentum deviation with same parameters as in Fig. 4.5
To do some first order investigations on the feasibility of the more complex setup, the model of the streaking electrodes is converted to a first order model in matrix form. Like in geometric optics, calculations in particle optics can be simplified by using transfer matrix notation. An optical system containing multiple optical elements can be modeled by multiplying the different transfer matrixes consecutively[31]:

\[ M_{\text{system}} = M_{\text{element}_N} \cdot M_{\text{element}_{N-1}} \cdots M_{\text{element}_2} \cdot M_{\text{element}_1} \]  

(4.17)

A particle can be described in phase space for our purpose by the vector with six components:

\[
\vec{P}_n = \begin{pmatrix} x \\ p_x \\ y \\ p_y \\ \Delta \\ z_0 \end{pmatrix}
\]

(4.18)

where \( z_0 \) is the initial position in the bunch relative to the reference particle and is related to the start time as \( t_0 = z_0 \frac{m}{p_z} \), and \( \Delta \) is the initial relative \( z \)-momentum deviation.

The first order streaking electrode model can thus be written in a \( 6 \times 6 \) transfer matrix:

\[
T_{\text{str}} = \begin{pmatrix}
1 & (x|p_x) & 0 & 0 & 0 & (x|z_0) \\
0 & 1 & 0 & 0 & (p_x|\Delta) & (p_x|z_0) \\
0 & 0 & 1 & (y|p_y) & 0 & 0 \\
0 & 0 & 0 & 1 & 0 & 0 \\
0 & 0 & 0 & 0 & 1 & 0 \\
0 & 0 & 0 & 0 & 0 & 1
\end{pmatrix}
\]

(4.19)

The \( x \) and \( y \)-direction are assumed decoupled, so the matrix is relatively simple. All \( x \) and \( p_x \) matrix elements can be found by multi-dimensional Taylor expansion of Eq. (4.12) and Eq. (4.14). The first order terms of the expansion that are not zero are given here:

\[
(x|p_x) = -q \frac{V_{\text{max}}}{h} \frac{1}{2} \frac{1}{p_{z0}}
\]

(4.20)

\[
(x|z_0) = -q \frac{V_{\text{max}}}{h} \frac{1}{2} \frac{1}{p_{z0}}
\]

(4.21)

\[
(p_x|\Delta) = -q \frac{V_{\text{max}}}{h} \frac{1}{2} \frac{1}{p_{z0}}
\]

(4.22)

\[
(p_x|z_0) = -q \frac{V_{\text{max}}}{h} \frac{m}{p_{z0}}
\]

(4.23)

In \( y \)-direction there are no forces acting on the particles, so it is simply drift space[31] of the length \( L_{\text{str}} \). Thus \( (y|p_y) \) is given by:

\[
(y|p_y) = \frac{L_{\text{str}}}{p_{z0}}
\]

(4.24)
4.4 Longitudinal phase space

Figure 4.7: Schematic overview of a reference particle through a homogenous sector magnet with parallel pole faces [31].

4.4.2 Bending magnet setup

A charged particle in a magnetic field experiences a Lorentz force perpendicular to its velocity, and will bend the trajectory of the particle. The most simple setup is a sector bending magnet with a constant homogenous magnetic field inside, see Fig. 4.7. The sector magnet has a focussing effect, with $l_1$ and $l_2$ the object and image distance. This total magnet system from object to image plane can be written in a transfer matrix, as is shown in Appendix A.

The energy resolution of the bending magnet is independent of the sector angle $\phi_0$ [31], so in principle a small magnet can be used. But at small $\phi_0$, about $\phi_0 < 45^\circ$, the total drift length through the magnet configuration, increases rapidly, so the total setup only increases in length by reducing the angle $\phi_0$.

4.4.3 Total system

The total system is described by the matrix $T_{sys} = T_{mag} \cdot T_{str}$. For completeness, the whole matrix is given in appendix B. But for the image, only the resulting $x$ and $y$ positions on the detector are important, not the expressions for the momenta.

The $x$-position on the detector is, as calculated with the matrix from the initial parameters $x_0, p_{x0}, \Delta_0$, and $z_0$ described by:

$$x_{screen} = x_0 + (L_{str} + L_{mag}) \frac{p_{x0}}{p_{z0}} + \frac{1}{2} qE_0 \tau \frac{L_{mag}}{p_{z0}} \Delta_0 - qE_0 \left( \frac{\tau}{2p_{z0}} + \frac{L_{mag} m}{p_{z0}^2} \right) z_0. \quad (4.25)$$

The first term is simply the initial $x$-position. The second term describes the drift of an initial velocity in the $x$-direction through the length of the whole setup. The third term describes the sensitivity to some variations in initial $z$-momentum and the fourth and final term describes the actual streaking. The first part of the fourth term is caused by the displacement of the particle when exiting the electrodes, and the second part, the desired linear deflection difference. Note that the letter part of this term is inversely proportional to the particle energy.

The $y$-position is simpler and given by:

$$y_{screen} = M_y y_0 + M_y L_{str} \frac{p_{y0}}{p_{z0}} + \rho_0 (1 - M_y) \Delta_0, \quad (4.26)$$
where the first term present the chosen magnification of the sector magnet with \( M_y \) the magnification factor, the second term the imaging by the bending magnet of drift of the particle between the electrodes. And the last term describes the dispersion for different momenta.

### 4.4.4 Projection of a Gaussian bunch

These two equations link an initial position in phase space to a position on the detector. To make it possible to say something about the resolution, it is needed to know how a distribution in phase space is imaged. To start with, the longitudinal phase space is set to a point in the origin, so for all particles \( z_0 \) and \( \Delta \) are zero. The \( x - p_x \) space is modeled by a simple uncorrelated 2D Gaussian:

\[
P(x_0, p_{x0}) \propto \left( \exp \left( -\frac{x_0^2}{2\sigma_{x0}^2} - \frac{p_{x0}^2}{2\sigma_{p_{x0}}^2} \right) \right).
\]

(4.27)

This corresponds to a transverse emittance \( \epsilon_x = \sigma_{x0}\sigma_{p_{x0}} \). This gives rise to a gaussian distribution on the \( x \)-direction of the screen with rms width \( \sigma_x \) given by:

\[
\sigma_x^2 = \sigma_{x0}^2 + (L_{str} + L_{mag})^2 \left( \frac{\sigma_{p_{x0}}}{p_{x0}} \right)^2.
\]

(4.28)

And in the same manner in the \( y \)-direction, where an uncorrelated gaussian in \( y - p_y \) space defined by \( \sigma_{y0} \) and \( \sigma_{p_{y0}} \) result in a gaussian distribution in the \( y \)-direction on the screen with:

\[
\sigma_y^2 = M_y^2 \left( \sigma_{y0}^2 + L_{str}^2 \frac{\sigma_{y0}^2}{p_{y0}^2} \right).
\]

(4.29)

A much smaller spot size can be obtained by adding particle lenses to the setup before the streaking electrodes. In the \( x \)-direction the spot has to be focussed over a distance of the total optical length \( L_{str} + L_{mag} \), and in the \( y \)-direction in a length of \( L_{str} \). On the other hand this might create unwanted distortion in the longitudinal phase space you want to image. Furthermore in this simple calculation, the effects of space charge are neglected, but when focussing they will eventually start to dominate.

In the longitudinal direction there exists generally a clear momentum-position correlation, so it is modeled as an rotated Gaussian distribution, see Fig. 4.8a. It is given by the parameters \( \sigma_z \) and \( \sigma_b = \sigma_\Delta \frac{\sigma_x}{\sqrt{\sigma_x^2 + \Delta_{max}^2}} \) that define the rms length of both axes and \( \Delta_{max} \) that result together with \( \sigma_x \) in the given angle.

If in this case the emittances of in the \( x \) and \( y \) directions are set to zero, the result of the longitudinal phase space distribution on the detector can be simply calculated with use of the linear transformation matrix. This results in a 2D Gaussian distribution defined by the following parameters, see Fig. 4.8b:

\[
\sigma_{x,scr} = \sigma_{x0}(x|z_0) + \Delta_{max}(x|\Delta)
\]

(4.30)

\[
y_{max,scr} = \Delta_{max}(y|\Delta)
\]

(4.31)

\[
\sigma_{b,scr} = \frac{\sigma_b}{\sqrt{\sigma_{x0}^2 + \Delta_{max}^2}} \sqrt{[(x|z_0)\Delta_{max} - (x|\Delta)\sigma_{x0}]^2 + (y|\Delta)^2\sigma_{x0}^2}
\]

(4.32)

The image on the screen of a real bunch with some non-zero emittance, is the convolution of this longitudinal distribution on the detector with the gaussian distribution caused by the
emittances. An estimation of the resolution can be made by comparing to different initial distributions, illustrated in Fig. 4.9a. A very short bunch with some given energy spread is imaged on the screen as a rotated Gaussian, because the sensitivity of the streaking process to momentum spread. An other initial bunch distribution, with the same energy spread, but now with some length $\sigma z_0$, has an extra 'rotation'. The distance between both the images $\Delta x_2$ in Fig. 4.9b is only caused by the length of the bunch so $\Delta x_2 = (x|z_0)\sigma z_0$.

The energy resolution of the system is given by:

$$\Delta_{res} = \sigma y_1(y|\Delta).$$  \hspace{1cm} (4.33)

The temporal resolution is estimated by:

$$\Delta t_{res} = \frac{m}{p_{z_0}} \Delta z_0 = \frac{m}{p_{z_0}} \sqrt{\frac{\sigma z_1 + \Delta_{res}^2 (x|\Delta)^2}{(x|z_0)}}.$$  \hspace{1cm} (4.34)

Both dependent on $\sigma z_1$ and $\sigma y_1$, which can be improved straightforwardly by adding focusing to the system as suggested earlier.

**Example calculations**

An example of an streaking image of a bunch is calculated for an electron bunch with an energy of 15 keV, 0.2 mm radius and a temperature of $\approx 10$ K. For the streaking a 30 kV pulse with a risetime of 200 ps is used with two electrodes with a length according to Eq. 4.11, ie. $L \approx 14$ mm and a separation of 10 mm. This is larger than the minimal separation distance of $\approx 6$ mm for these parameters needed to avoid collisions with the electrodes. The results are presented in Fig. 4.10. A bending magnet with $\phi_0 = 45^\circ$, $M_y = -1$ and $\rho_0 = 10$ cm is used.

In this calculation an initial longitudinal phase space distribution is defined by the parameters shown in Fig. 4.8. A negative correlated energy spread of a rms value of 2% is chosen, so $\Delta_{max} = -0.02$ and further the thickness of the distribution given by $\sigma_0 = 2 \times 10^{-5}$. In the first image, Fig. 4.10a, the streaking image due to solely energy spread is presented with $\sigma z_0 = 0$. In Fig. 4.10b the image resulting form a correlated longitudinal distribution with $\sigma z_0$ corresponding with a rms length of 2 ps.

There is no extra focusing, so the spot size on the screen will increase due to the finite emittance. In Fig. 4.11 this spot is shown, if a rms lateral velocity of 100 m/s is assumed for
the electrons. The image measured on a detector is the convolution of this spot together with the streak of Fig. 4.10b, that results again in a Gaussian distribution.

For these parameters the momentum resolution is estimated by Eq.4.33 to be about $\Delta_{res} \approx 0.005$ and the longitudinal resolution is approximated with use of Eq. (4.34) to be $\Delta t_{0, res} \approx 0.5$ ps.

Conclusion and discussion

In the section it has been shown that the electric field pulse can be in principle be used in a streaking setup to measure the longitudinal phase space in a single shot. A simple analytical approximation is given, so the influence of the different parameters can be investigated. The achievable resolution depends on all bunch and electric field pulse parameters.

In this simple model, the space charge effects and all non ideal fields are neglected. To make a more realistic investigation in the possible achievable resolutions, numerical simulations on a setup with for example GPT[23] are needed. Also more complicated setups are possible to increase the resolution. First of all, as suggested before, focussing can be added to decrease the spot size due to initial emittance and with use of a stigmatic focusing sector magnet with inclined field boundaries[31], a higher momentum separation in the y-direction can be achieved.
Figure 4.10: Calculated longitudinal streaking image with use of the matrix method. a) The image for a very short bunch without correlation in phase space b) Same bunch but with an rms length of 2 ps. The axes are in m.
Figure 4.11: Calculated image of only the effect of emittance on the spot size if there is no extra beam focussing. The axes are in m.
Chapter 5

First UCP electron beam and ion beam experiment

In this chapter the performed experiments are explained and the results are presented. In Sec. 5.1 the results are presented of a MOT in a test setup. A simple test setup was used to get experience building a MOT and get a better understanding of the underlying physics and characteristics. In Sec. 5.2 the experimental setup used for the first UCP electron beam and ion beam experiments is explained in detail. A larger vacuum system is used, compared to the test setup, and in addition to two diode lasers as in the test setup, also a tunable dye laser is used. The experimental procedures used for the first performed UCP beam experiments are described in Sec. 5.3. In Sec. 5.4 the results are presented of time resolved experiments done on the spontaneous evolution of a cold Rydberg cloud into a UCP. Finally, in Sec. 5.5 the results are presented of the first ultra cold beam experiments done. For the first time phosphor images are made of particles originating from a UCP.

5.1 First trapped rubidium

A major part of the cold bunches setup is the creation of a rubidium MOT. In this section the practical realization of such a trap is described and the first characterization measurements are presented of a test setup. The laser system and MOT diagnostics described in this section, are also used in the UCP setup, described in Sec. 5.2.

5.1.1 Rubidium

Alkali atoms are often used for trapping and cooling. Of the alkalis rubidium ($^{85}\text{Rb}$) in particular is the workhorse, because among other things a transition from the ground state can be used for cooling, and low cost and stable diode lasers are available that can drive the wanted optical cooling transition.

In Fig. 5.1 the transitions used for cooling are shown. The laser cooling is operated from the ground state $5S(J = 1/2)$ to the fine structure state $5P(J = 3/2)$. Because of the nuclear spin ($I = 5/2$) of $^{85}\text{Rb}$ and the hyperfine coupling with $F = I + J$, both states have a hyperfine structure.

For laser cooling and trapping a closed transition is always needed. The trapping laser uses the hyperfine transition $5S_{1/2}, F = 3 \rightarrow 5P_{3/2}, F = 4$ at a wavelength of 780 nm. Unfortunately this light also excites some atoms to the $5P_{3/2}, F = 3$ state 121 MHz off resonance. This state can decay back to the $5S_{1/2}, F = 2$ ground state level, that can no longer be excited by the trapping laser.
The solution to this problem is to add an extra laser, the repump laser which excites these atoms to the $5P_{3/2}, F=3$ state, that can decay again to the needed $5S_{1/2}, F=3$ state. The most important properties relevant for laser cooling for $^{85}\text{Rb}$ are given in table 5.1[32]. The saturation intensity $I_0$ can be calculated with the equations from Sec. 2.1.

Table 5.1: Properties of $^{85}\text{Rb}$ relevant for laser cooling on the $5S_{1/2}(F=3)\leftrightarrow 5P_{3/2}(F=4)$ transition [17].

<table>
<thead>
<tr>
<th>Quantity</th>
<th>Symbol</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Atomic mass</td>
<td>$m$</td>
<td>85 a.m.u. = $1.41\times 10^{-25}$ kg</td>
</tr>
<tr>
<td>Wavelength</td>
<td>$\lambda$</td>
<td>780.24 nm (in vacuum)</td>
</tr>
<tr>
<td>Natural linewidth</td>
<td>$\Gamma$</td>
<td>5.98 MHz</td>
</tr>
<tr>
<td>Lifetime of $6P_{3/2}(F=4)$</td>
<td>$\tau$</td>
<td>$1/(2\pi\Gamma)$ = 26.63 ns</td>
</tr>
<tr>
<td>Saturation intensity</td>
<td>$I_0$</td>
<td>1.64 mW cm$^{-2}$</td>
</tr>
<tr>
<td>Recoil velocity</td>
<td>$v_{\text{rec}} = \hbar k/m$</td>
<td>0.602 cm s$^{-1}$</td>
</tr>
<tr>
<td>Doppler limit, velocity</td>
<td>$v_D = (\hbar \Gamma/2m)^{1/2}$</td>
<td>11.85 cm s$^{-1}$</td>
</tr>
<tr>
<td>Doppler limit, temperature</td>
<td>$T_D = \hbar \Gamma/2k_B$</td>
<td>142.41 $\mu$K</td>
</tr>
</tbody>
</table>

5.1.2 Diode laser

As described in the last section, two lasers are needed to build a rubidium MOT, the trapping laser and the repump laser. Both lasers should be tuned and locked to a different atomic transition. The specification for the trapping laser is a bit stricter than for the repump laser. It should have enough laser power at the right atomic resonance frequency and a frequency stability better than the natural atomic line width of $\approx 6$ MHz.

For the trapping, a commercial diode laser (Toptica DLX) is used that can generate up to about 950 mW of 780 nm laser light. This is almost a factor 10 higher than the power needed for a typical MOT with a saturation parameter of about 10 for a beam diameter of 10 mm. So the energy loss in the optics can be easily compensated and even higher saturation values can be used if needed. For the repump laser, a commercial diode laser (Toptica DL100) is used that
can generate up to 150 mW laser power at 780 nm.

The lasers are mounted on one end of an optical table. To reduce effects of shocks and vibrations they are mounted on a shock absorbing rubber. The two lasers, together with all the optics need to lock the laser at their frequencies, are mounted in a closed box. In this way the density fluctuations caused by airflow and the movement of optics by temperature are suppressed.

In Fig. 5.2 an overview is given of the setup of both lasers. The setup of the repump laser is the less complicated one. A small portion of the beam, about 5 mW, is split off by a glass plate. This laser beam is used to generate a frequency error signal by Doppler free absorption spectroscopy [33] in a cuvette with rubidium vapor. This signal is used by a feedback PI controller, to adjust the position of a piezo driven grating in the cavity built in the diode laser. By adjusting this grating the frequency can be adjusted and thus locked and kept stable at the right atomic transition.

Figure 5.2: Schematic overview of the setup of both lasers, the trapping laser and the repump laser, with all the locking optics in place[17].
The trapping laser setup is a bit more complex, because it uses polarization spectroscopy instead of absorption spectroscopy, to get a better frequency stability[33]. For future purposes, such as a 2D MOT where atoms are cooled to a line to produce an atomic beam[17], the laser beam is divided in two independent beams.

Furthermore to make it possible to detune the laser in a controlled manner, three acoustic-optical modulators (AOM) are added to the setup. In an AOM the photons of the laser interact with phonons in the crystal, generated by a computer controlled rf signal, and produce a deflected and frequency shifted laser beam. By using them in a double pass configuration, as in Fig. 5.2, the alignment of the laser is not changed with changing the AOM frequency. Also twice the frequency shift is achieved in this way. Unfortunately AOMs always operate in some frequency range, in our case 60 – 100 MHz with center 80 MHz, so small detunings are not possible. To make these possible, not only an AOM is placed in the feedback system, but also in the high power output beams, so the frequency can be shifted back on resonance.

The stability of the laser was measured by looking at the error signal. The peak to peak fluctuations correspond to a stability smaller than 1 MHz, much smaller than the 6 MHz natural linewidth. The two laser beams that exit from the laser have both about 200 mW power and the detuning can be set independently.

5.1.3 Vapor-cell MOT

There exist different ways to load the MOT with rubidium atoms. A typical MOT setup uses a slow atomic beam to load the MOT. If the velocities in this beam are smaller than the capture velocity $v_c$, $v_c \approx 40$ m/s for rubidium [34], the atoms are trapped. It is also possible, to load the atoms directly from a room temperature rubidium background gas, a so-called vapor-cell MOT. The capture velocity is relatively high, so even at room temperature there are enough particles in the tail of the Boltzman distribution with a sufficiently low velocity. The principle of a vapor-cell MOT is shown in Fig. 5.3. Only sufficiently slow atoms will be cooled and trapped, the rest will escape.

There are two processes that can cause trapped atoms to be lost. So-called the first order loss results from collisions of the cold atoms with the hot background gas. The cold atoms gets an velocity higher than the capture velocity. And the second order process, where collisions between the cold atoms change it state, so they are no longer trapped. A typical MOT can have densities up to $10^{-12}$ cm$^{-3}$ and lifetimes of several seconds.

The advantage of loading with an atomic beam, is that the background pressure can be lower than with a vapor-cell MOT. In that case the losses by collisions with the hot background are lowered, so the lifetimes of the atoms in the MOT are increased and also the loading rate can be enhanced with about 2 orders of magnitude because of the larger influx of atoms. This also makes a larger MOT possible, about $10^{10}$ atoms instead of $10^9$ atoms with a vapor-cell MOT. Vapor-cell MOTs with $10^{10}$ atoms have been reported[34], but only with a relative high background pressure, which could cause problems by background collisions with the low emittance electron beams.

For our project, we are not interested in long lifetimes per se, in contrast to, for example experiments aiming at to realize a Bose Einstein condensation from a MOT, but only in the number of atoms trapped in the MOT. Because $10^9$ atoms is large enough for all our initial experiments, we started with a vapor-cell MOT. This makes the experiment a bit simpler, because no atomic beam is needed. The only thing that is required is a rubidium ampul connected to the vacuum chamber.

If the two-body collisions inside a MOT are neglected, the loading behavior of a vapor cell
Figure 5.3: Schematic overview of a vapor cell MOT. In the volume where the laser beams intersect, the atoms can be cooled and trapped. Only particles from the background, with velocities below the trapping velocity, are trapped (trajectories A and D). Atoms entering the volume with higher velocities will not be trapped (trajectories B and C)[17].

MOT can be described by:

\[
\frac{dN}{dt} = R_L - N \Gamma_{BG},
\]

(5.1)

where \(N\) is the number of trapped particles, \(\Gamma_{BG}\) the linear loss term due to the background gas. The capture rate \(R_L\) is given by[35]:

\[
R_L = \frac{n_{BG} V^{3/2} v_c^4}{2v_{mp}^3},
\]

(5.2)

where \(n_{BG}\) is the density of rubidium background vapor, \(V\) the capture volume and \(v_{mp}\) the most probable velocity of the background atoms. The solution of this differential equation is given by:

\[
N = \frac{R_L}{\Gamma_{BG}} \{1 - \exp(-\Gamma_{BG} t)\}.
\]

(5.3)

This means that the MOT fills on a timescale determined by the losses due to background atoms. The steady state number of atoms \(N(\infty)\) is given by:

\[
N(\infty) = \frac{R_L}{\Gamma_{BG}} = \frac{n_{BG} V^{3/2} v_c^4}{\Gamma_{BG} 2v_{mp}^3}.
\]

(5.4)

So to capture a large number of atoms in the trap, large MOT beams are needed. They result in a large capture volume \(V\) and also in a high capture velocity \(v_c\). Both the loading rate \(R_L\) and the loss rate \(\Gamma_{BG}\) depend approximately linear on the rubidium background pressure, as long as the mean free path of the atoms are larger than the dimension of the trap. In that regime is the steady state atom number independent of the rubidium pressure, but the loading time scales inversely proportional to the rubidium background pressure.
5.1.4 MOT diagnostics

The most important properties of a MOT can be obtained from looking at the fluorescence light emitted by the atoms in the trap. The power of this emitted light can be used to calculate the total number of atoms in the trap. If a calibrated CCD camera is used, also information of the spatial distribution can be obtained, so the size of the sample and the density can be measured.

The fluorescence power $P$ emitted is given by:

$$P = \hbar \omega \Gamma f N,$$  \hspace{1cm} (5.5)

with $N$ the number of trapped atoms, $\hbar \omega$ the energy of an emitted photon, $\Gamma$ the decay rate from the excited state and $f$ the fraction of atoms in the excited state:

$$f = \frac{1}{2} \frac{CS}{1 + 4 \left(\frac{C S}{1}\right)^{2} + C S},\hspace{1cm} (5.6)$$

where $S$ is the total saturation parameter of all the MOT beams together, which can be calculated from the intensity of the laser beams. The factor $C$ is a phenomenological quantity that depends on the specific atomic transition. We used $C = 0.7$ based upon the work of Townsend et al\[36\].

The current $I$ generated by a photodiode is then given by:

$$I = \eta P \frac{\Omega}{4 \pi},\hspace{1cm} (5.7)$$

where $\eta$ is the power to current calibration factor, which has a value of $0.5 \text{ A/W}$ at 780 nm for the used photodiode and $\Omega$ the solid angle subtended by the detector. From this current the number of trapped atoms can be calculated.

5.1.5 Vacuum system and magnets of the test setup

To build a MOT not only a stable laser as described in a previous section is needed, but also a vacuum system in which the MOT can be created. To make a large number of trapped atoms possible, the pressure has to be sufficient low, in the order of $10^{-8}$ mbar thus reducing the losses in the trap due to elastic and inelastic collisions with hot background particles.

To test the stability of the lasers and get experience building and aligning the MOT optics, a relative simple test setup was built, which will only be described very briefly. The setup used for the first ultra cold beam experiments will be described in more detail in Sec. 5.2.

The test setup consisted of a small vacuum chamber, with 8 vacuum ports, that was pumped down to about $10^{-8}$ mbar using a turbo and an ion-getter pump. A rubidium ampul was placed in a small vacuum tube and connected to the setup with a small mechanical valve, to let the rubidium vapor into the vacuum chamber. The required magnetic field gradient was generated by a pair of coils in anti-Helmholtz configuration, see Sec. 2.1.2. Because of the small dimensions of the vacuum chamber, the coils were mounted outside the vacuum just above and below the chamber. They can generate a magnetic gradient of about 20 G/cm at maximum current. The trapping laser beam was blown up in size with a telescope to about 2 cm diameter, and mixed with the repump laser. The resulting beam was split in three different beams by polarizing beamsplitters and all sent through a quarter lambda plate to adjust the polarization. Instead of using six independent MOT beams, these three beams were retro-reflected to create the other three counter propagating laser beams. This makes the alignment and the optics easier, but due to absorption in the MOT a shadow is cast in the retro reflected beam, which may cause an imbalance in cooling power. For the first experiments however this is not a critical problem.
In this setup a calibrated photodiod was added which can collect a solid angle of about $\Omega = 1.5 \times 10^{-3}$ sr. Also a CCD camera was mounted to measure the dimensions of the cold atom cloud. The CCD camera was calibrated with respect to the calibrated photodiod.

5.1.6 Results

A snapshot of the first cold atom cloud produced in the setup is shown in Fig. 5.4. After the first observation, the alignment and settings were improved, so a more gaussian cloud profile was achieved.

Number of atoms and density

The number of atoms was measured by fluorescence using the calibrated photodiode described previously. In Fig. 5.5 the steady state number of atoms is plotted as function of the power of the trapping laser and the power of the repump laser. The other parameters of the MOT, the magnetic field gradient and the detuning, were set on their optimum values determined in a measurement with constant laser power. The detuning of the trapping laser was set at $-18$ MHz and the magnetic field gradient at 8 G/cm. The rubidium background gas pressure was measured by absorption, and was estimated to be about $5 \times 10^{-8}$ mbar.

In the top part of Fig. 5.5 the number of atoms measured is shown against the power of the repump laser. In this measurement the power of the trapping laser was fixed at 60 mW. The number of atoms start to saturate around a power of 30 mW. This corresponds with a saturation parameter of $s \approx 10$. This value is already much higher than reported in most of the literature. This can be explained by the relative high power of the trapping laser. For low saturation values there is only a $\approx 10^{-4}$ probability to excite the the wrong state 121 MHz off resonance, but for higher saturation values, this value can be much higher because of power broadening.

In the lower part of the figure, the number of atoms is plotted against the power of the trapping laser. In this case the power of the repump laser is kept constant at 60 mW. In this region, there is a clear linear dependence visible between the the laser power and the number of trapped atoms. There is no sign of saturation yet. The most important conclusion here is that
Figure 5.5: The number of atoms measured in the MOT by measuring the fluorescence signal with a photodiode as function of the power of the trapping laser and the repump laser. (a) The repump laser shows saturation behavior, while the trapping laser (b) shows almost a linear relation.
it is possible to trap more than $10^9$ atoms in the MOT at a power of 120 mW. This corresponds to a saturation parameter $s \approx 70$. This number of atoms is well beyond the number needed for the first beam and plasma experiments.

The density of the cold atom cloud is measured by fitting a 2D Gaussian to the fluorescence profile measured with a CCD camera. Because only a single camera was used, the projection in the $x$-$z$ plane was observed. The distribution is assumed to be symmetrical in the $x$-$y$ plane ($\sigma_x = \sigma_y$), because both field gradients in those directions are the same. The maximum density $n_0$ can be calculated as:

$$n_0 = \frac{N}{(2\pi)^{3/2}\sigma_x \sigma_y \sigma_z}.$$  \hspace{1cm} (5.8)

with $N$ the number of atoms in the trap, measured with the calibrated photodiode. For the MOT under normal operation circumstances with $1.2 \times 10^9$ atoms, a width in $x$-direction of $\sigma_x = 0.18$ mm and in $z$-direction of $\sigma_z = 0.20$ mm are found. This gives a density of $10^{10}$ atoms/cm$^3$.

**Loading rate**

By measuring the time dependent signal of the photodiode, when the trapping lasers are switched on, the loading process of the MOT can be monitored. In Fig. 5.6(a) six loading curves are shown, measured at different trapping laser detunings. Also the fits of Eq. (5.3) to the data with two parameters, the loading rate $R_L$ and the linear loss rate $\Gamma_{BG}$, are shown. The data is in very good agreement with the fits. In the case of a 2.5 MHz detuning different behavior was observed. After about 1.5 s the MOT position shifted to a new optimal position, which can be seen by the small deviation from the fit. This effect is due to the shadow forming in the retro reflected beam at higher densities.

The fluorescence signals in steady state have rms fluctuations of less than 2%, measured over more than 20 s, so the MOT shows very stable behavior.

In Fig. 5.6(b) the fitted lifetimes $1/\Gamma_{BG}$ are displayed as a function of detuning. The lifetimes increase and saturate to a maximum value of about 0.75 s at a detuning of $-18$ MHz. All the lifetimes are much longer than the typical extraction timescale of the bunches from the plasma, so losses by background collisions should not be a problem. The maximum loading rate of about $R_L \approx 8 \times 10^8$ atoms/s is also found at this detuning. This value is high enough to extract about 10 pC charge with a repetition frequency of 10 Hz. But if needed it can be increased up to 2 orders of magnitude by loading from a high flux cold atomic beam, that can for example be generated by a 2D MOT.[17].

**Temperature**

To make an estimate of the temperature of the cold atom cloud, a simple expansion and recapture experiment is done. The results are displayed in Fig. 5.7. When the MOT is fully loaded, the trapping laser is turned off with use of the AOM in less than 1 $\mu$s. After a delay time $\Delta t$ of several ms, the laser is switched on again, so the atoms are recaptured quickly. Because of ballistic expansion, a fraction of the atoms has expanded outside the capture volume of the MOT during the delay time and is therefore lost. The rate of recapture is much faster than the loading rate from the background gas, so the lost fraction can be measured directly.

The fraction $f_r$ of atoms recaptured as a function of delay time can be described by a simple expansion model derived by Molenaar[37]:

$$f_r = \text{erf} \left( \frac{\nu_{\text{max}}}{\nu} \right) - \frac{2}{\sqrt{\pi}} \left( \frac{\nu_{\text{max}}}{\nu} \right) \exp \left( - \left( \frac{\nu_{\text{max}}}{\nu} \right)^2 \right).$$  \hspace{1cm} (5.9)
5.1 First trapped rubidium

Figure 5.6: (a) Loading curves of the MOT plotted for 6 different detunings values of the trapping laser, with their corresponding fits (dashed lines) of Eq. (5.3). (b) Lifetimes as a function of detuning, resulting from the fits in (a).

Figure 5.7: The measured fraction of atoms left in the MOT plotted against the interruption times $\Delta t$ of the trapping lasers. The dashed line is a fit of Eq. (5.9) to the data.
First UCP electron beam and ion beam experiment

Here \( \bar{v} = \sqrt{2k_BT/m} \) is the thermal velocity of the atoms in the MOT at a temperature \( T \) and \( v_{\text{max}} = R_c/\Delta t \) with \( R_c \) the capture radius of the MOT beams.

In Fig. 5.7 a fit is shown of Eq. (5.9) to the measured \( f_r \) as a function of \( \Delta t \). From this fit a temperature of \( 200 \pm 100 \mu\text{K} \) is found. This temperature is close to the expected Doppler temperature of \( 142 \mu\text{K} \). The large uncertainty in the temperature is mainly due to the large uncertainty in the exact capture radius \( R_c \). If needed a more precise temperature can be obtained by measuring the ballistic expansion using absorption imaging.

5.2 UCP setup

5.2.1 Dye laser

A tunable pulsed dye laser of the type Quanta-Ray PDL3, is used to ionize rubidium just above threshold or excite the atoms to a Rydberg state. To generate pulses with a wavelength around 480 nm, Coumarin 480 dye is used. It is optically pumped with the third harmonic of a Nd:YAG laser that operates at 10 Hz, producing 10 ns laser pulses with a bandwidth of about 15 GHz, a surface area of about 1 mm\(^2\) and an energy of about 1 mJ.

The wavelength is continuously measured with a calibrated home-built monochromator, with a resolution of about 0.03 nm. The wavelength of the laser can be adjusted with a computer controlled grating. This system can be used to automatically go to the right wavelength and if needed compensate for drift.

5.2.2 Vacuum setup

In Fig. 5.8 a schematic overview of the UHV vacuum system is shown. It consists of a large trap chamber with 18 ports. This was chosen so a future accelerator structure can be mounted inside. There are twelve 40 mm diameter (CF40) ports, which are mainly used for the MOT beams, ionization laser beam and MOT diagnostics. The four windows used for the MOT beams have a 780 nm AR coating, and two windows for the ionization laser have 480 nm AR coating, to reduce unwanted reflections.

Furthermore the chamber has four 200 mm diameter (CF200) ports. One of them on the side is used to mount the accelerator rod structure, described in Sec. 5.2.4, an other one on top is used for the multi channel plate detector. A single 100 mm diameter (CF100) port is available that can be used in the future to mount an atomic beam source, like a 2D-MOT[17]. And finally the last port, a 160 mm diameter (CF160) port, is connected to the vacuum pumps.

The setup is pumped down with a 250 l/s turbo molecular pump, backed by an oil free turbodrag pump and diaphragm pump. Using these pumps, the chamber is pumped down to a pressure of about \( 10^{-6} \) mbar. At this pressure a 150 l/s ion-getter pump is turned on, that can pump down the chamber to a final pressure of about \( 10^{-9} \) mbar. Heating elements are placed around the chamber for bakeout. A manual valve is added between the turbo pump and the rest of the setup, so the turbo pump can be turned off to minimize the vibrations of the optical table, that can for example disturb the stability of the laser.

A rubidium ampul is connected to one of the CF40 ports. It can be openend and closed with a manual valve, so the inlet of rubidium in the chamber can be controlled. A heater is mounted on the ampul, so a higher vapor pressure can be generated.

5.2.3 Magnetic fields

A MOT a magnetic field gradient of about 10 G/cm is created at the trap center by means of a pair of coils, mounted inside the vacuum chamber. The coils both consist of 4 turns with
The trap chamber is pumped down with a turbo molecular pump and two ion getter pumps. Above the chamber an optical platform is placed to mount all the MOT optics and a multi channel plate detector to observe electrons and ions.
First UCP electron beam and ion beam experiment

Figure 5.9: Measurement of the timings of the electric and magnetic fields in a typical experiment. The magnetic field, black line, is measured with a Hall probe. The field can be turned on and off within 150 μs. The electric field, gray line, is measured with a high voltage probe.

For the operation of the MOT, it is important to have zero magnetic field precisely at the center of the trap. There are always non-zero magnetic fields present like the earth magnetic field and the field from the ion-getter pumps, so extra compensations coils are placed around the vacuum chamber. They can be used to null the magnetic field at the center, and can also be used to change the position of the cloud of cold atoms over several millimeters, by moving the magnetic field minimum.

5.2.4 Pulsed extraction field

For the first experiment a simple accelerator structure has been used, consisting of four parallel rods separated by 20 mm. The MOT is created at the center of the rod structure. The upper two rods are connected to ground, while the lower two are connected through a high voltage vacuum feedthrough to a high voltage source, which can deliver both positive and negative pulses. In Fig. 5.10 an illustration is shown of the rod structure mounted together with the MOT coils.

The electrostatic field generated by the rods can be approximated by the 2D field of infinitely long rods. This 2D field is calculated with the code Superfish [38].

The high voltage pulser can generate pulses up to 3 kV with a risetime of about 2 μs. Due to this relatively slow risetime, the field can be approximated by the static field multiplied with a time dependent factor given by the high voltage pulse shape.
Figure 5.10: Illustration of the accelerator rod structure together with the MOT coils. The two upper rods are grounded, while the lower two are connected to a pulsed high voltage source. In the middle between the rods a small cloud of atoms is shown.

5.2.5 Optics

MOT beams

A schematic overview of the optics used to create the MOT beams is given in Fig. 5.11. A symmetric setup is chosen so all diagonal MOT beams have the same length. This can be used in the future to build a dark spot MOT without the need to change the whole optical setup.

An extra platform is placed on top of the vacuum chamber to mount the optics. For the first experiment, a retro-reflected three beam MOT is used, as in the test setup. In the future this can be easily expanded to six independent beams to improve the behavior of the MOT at higher densities, but the behavior as seen in the test setup was already good enough for the first experiments.

The arrangement of the beamsplitters is such that two counter propagating beams are derived from the same polarizing beamsplitter. This makes the adjustment of the intensity balance easy, because only one half lambda plate has to be rotated.

Both the trapping laser beam and the repump laser beam are blown up in size by a telescope pair to a diameter of approximately 2 cm. An adjustable aperture is added to clip the laser, so unwanted reflections can be minimized, which make the detection of the MOT more difficult. In typical operation we used 120 mW for the trapping laser, which results in a saturation parameter $s \approx 70$. The repump laser power is typically $\approx 50 \text{ mW}$.

MOT diagnostics

A calibrated photodiode together with 2 lenses, as shown in Fig. 5.12, are mounted inside a black cylinder. It is used to measure the 780 nm fluorescence light of the MOT. The extra focus with the pinhole is used to minimize the influence of unwanted background light and laser reflections. In this configuration the photodiode collects light from a solid angle of $\Omega = 6.6 \times 10^{-4}$ sr.

In the future, two additional CCD cameras will be installed, so the spatial distribution and
Figure 5.11: An schematic illustration of the optics used to create the MOT beams. A symmetric setup is chosen so in the future a dark spot MOT[34] can be created without changing all the optics. The black lines are beams at the top level, and the gray lines are beams at the low level. In the first experiments a three beam retro reflection configuration is used, so the top level optics are left out and the beam is reflected back with a single mirror.

Figure 5.12: Schematic overview of the setup used to detect the fluorescence light from the MOT. An extra focus lens with pinhole is added to minimize the influence of stray-light.
size of the MOT can be observed as well. The number of atoms measured with the photodiode can be used for calibration of the CCD cameras. From the calibrated cameras all the information needed can be acquired: the total number of particles, the spatial dimensions, and the precise position.

5.3 Experimental procedure

The first experiments in the ultra cold bunches setup have been done to test and show it is possible to produce and extract electrons and ions from a UCP. The UCP is created by excitation of the rubidium atoms to a Rydberg state which subsequently ionize due to collisions and black body radiation. In Fig. 5.13 a schematic illustration is given of the experimental procedure of these first experiments:

A) First, an ultra cold sample of rubidium gas is trapped in the MOT as described in the previous section. B) The next step is to turn off the trapping lasers, which can be done in less than a µs using AOMS. The atoms are not cooled anymore, but because of the low average velocities, they can be considered frozen at µs timescales. In step (C) the magnetic field of the MOT is turned off in about 150 µs. This is done because the magnetic field will affect the trajectories of the charged particles in the beam.

After the magnetic field is turned off, the trapping laser is again switched on, so the \(5P_{3/2}, F = 4\) state is repopulated. After about 20 µs, the atoms are excited to the \(44d\) Rydberg state (D), using the 380 nm pulsed dye laser. After a delay \(\Delta t\), which can be varied between 0 to about 80 µs, the electric field pulse is turned on (E).

The field has a rise time of about 2 µs. It rips the plasma apart and field ionizes all Rydberg atoms in states with principle quantum number \(n \geq 37\). By applying either a positive or a negative pulse, electrons or ions are accelerated in the direction of the MCP detector. In Fig. 5.14 an overview is given of the detection of particles from the UCP. A phosphor screen image is registered by the CCD camera and simultaneously the current pulse is measured with an oscilloscope.

Immediately after this procedure, the MOT coils are turned on again, so the MOT can reload. This procedure is repeated at 10 Hz, the repetition frequency of the dye laser.
Figure 5.14: Schematic overview of the setup used to create and detect ultra cold atoms and ions from the UCP. A) CCD camera and image, B) MCP detector with phosphor screen, C) Time trace on oscilloscope, D) MOT beams, E) Ionization laser, F) The MOT coils, G) High voltage pulse
5.4 Time resolved beam experiments

5.4.1 Results and discussion

Electron signal

In Fig. 5.15 three typical electron traces are shown, measured with the MCP detector as described in the previous section. These traces show the normalized MCP signal, a measure for current, as function of time for different delay times $\Delta t$ between the Rydberg excitation pulse and the beginning of the electric field pulse. To enable a comparison between the multiple curves, the trigger of the electric field pulse is set to $t = -2 \mu s$, so the peak corresponding to the electrons due to the initial 44d state arrives at $t = 0$ at the detector. The electrons exit the accelerator structure with an energy of about 10 eV. The time of flight of the electrons from the MOT to the MCP detector is in the order of 200 ns depending on the precise energy. Also the curves are normalized to the total charge, the area surface of the signal, to make the comparison easier.

Since time of flight can be neglected, compared to the electric field timescale, signals to the left of the central 44d peak correspond to electrons stripped at lower fields. In a simple interpretation these are electrons from Rydberg states with $n > 44$ or electrons stripped from the plasma. At the right of the central peak, electrons stripped at higher electric fields are detected corresponding to Rydberg states with principle quantum number $n < 44$. The lowest state that the electric field pulse still can field ionize is about $n = 37$. In reality however the ionization fields also depend on the angular momentum quantum number $l$, nearby $n$ states can overlap and make a precise one on one translation of the peaks to the different Rydberg states more complex.

At a delay time of about 0 $\mu s$, the sharp central peak corresponds to the initial 44d state that is excited by the laser. The electric field pulse, which is measured simultaneously with a high voltage probe, is calibrated with use of the peak of the 44d state. A second, less resolvable
peak to the right corresponds to electrons from the 46p state. By dipole-dipole interactions between the Rydberg atoms in the 44d state, some are excited to this p-state[39]. Collisions of electrons with the Rydberg atoms result in a distribution over different l-states with the same quantum number $n$. This l-mixing[40] causes the broader wing on the right, because of the higher ionization field needed for those states.

At a time delay of 8 $\mu$s electrons are detected that strip at very low fields of a few V/cm, and thus are very weakly bound. The ratio of atoms in a lower state and the number in the initial state, indicates the creation of a very weakly bound plasma as expected. Free electrons are created by black body radiation and collisions ionize more Rydberg atoms. The fastest electrons escape, so the slower moving electrons are trapped in the created positively charged potential well of the ions, as described in Sec. 2.3.

At a delay time of 22 $\mu$s a single very clear plasma peak is detected at even lower stripping fields. This is expected because the plasma is expanded so the trapping potential is lower, and the electrons are more weakly bound. These results are qualitatively consistent with other measurements on the spontaneous evolution of a cloud of Rydberg atoms into a UCP[40].

**Ion signal**

In Fig. 5.16 the traces are displayed of two different measurements on ions. Because of the large mass difference between electrons and ions, the time of flight timescales are totally different. For electrons the time of flight (about 200 ns) can be neglected compared to the ionization electric field pulse. For ions that gain an energy of about 400 eV in the accelerator structure however, the time of flight is about 10 $\mu$s, which is much larger than the electric field pulse. This makes a precise quantitative analysis more difficult.

The density of the initial cold atom cloud in the MOT was varied by changing the power of the repump laser. At a lower density, there are of course less atoms excited to the initial Rydberg state, but also the atoms are not converted to a plasma after approximately 7 $\mu$s delay time. At higher initial density, more atoms are excited, but they also convert rapidly into a plasma in that same delay time, because of the increased interactions between the cold Rydberg atoms.

**Plasma fraction**

The measured time traces can be divided into two parts. The low electric field part, which contains the charge that is collected from states that strip from lower electric fields than the 97 V/cm needed for the initial 44d Rydberg state. It contains particles that are stripped from the plasma and from atoms in a higher Rydberg states than the initial excited state.

The second part, the high electric field part, contains the collected charge of particles from the initial excited Rydberg state and the lower lying Rydberg states, down to $n = 37$. For electrons the edge between the regions is set at $t = -0.5$ $\mu$s, see for example Fig. 5.15. For ions a value of $t = 16.5$ $\mu$s is used.

We define the plasma fraction as the measured charge in the low electric field part divided by the total measured charge. In Fig. 5.17 the plasma fraction is shown for both electrons as ions. The actual fractions might be lower, because no Rydberg states with $n < 37$ are detected.

The fraction of both the ions as well the electrons start to increases in about 8 $\mu$s to about 55%. The electron signal stays almost constant at this value, but the plasma fraction of the ion signal increases to a maximum of about 90% at a time delay of 25 $\mu$s before it starts to decay.

The difference in behavior can be explained by the difference in detection efficiency between the electrons and ions. After the plasma is created in about 10 $\mu$s it will expand, so the particles become less deeply bound, so they get stripped of at lower fields. This results in electrons and
Figure 5.16: Two ion traces measured both at a delay time of 7 μs, but for two different initial Rydberg atom densities. The gray line is the case for a low initial density, only the original 44d state is present. At higher density, the black line, a large plasma peak appears. The time of flight for ions is about 10 μs, so the electrons of the original state arrive at the detector about 17 μs after excitation.
ions with energies of only several eV. The influence of small magnetic and electric stray fields becomes important at those energies for electrons. Due to the larger mass, the fields only have a relative small effect on the slow ions. Therefore we suspect that because of some stray fields, those electrons never reach the detector. So for the electrons a lower plasma fraction is found. In practice we also see this difference in sensitivity to stray fields. A small change in the magnetic compensation coils can have a large effect on the electron signals, but only a small influence on the ion signals is observed.

Also the measured decay constant of both the ions and electrons is different. While the plasma is expanding and gets less deeply bound, fast moving electrons can escape more easily than heavy ions. This can explain the difference in decay time that is observed in the measurement.

Numerical simulations on the spontaneous evolution of a cold Rydberg cloud to a UCP are done by Pohl et al.[41]. They started with a cloud of Rydberg atoms excited to \( n = 70 \). The calculated plasma fractions show a qualitatively comparable behavior: they first rise in a few \( \mu s \) to 75% after which they decay again because of collisions. The remaining Rydberg atoms redistribute to a lower excited state around \( n = 25 \) in about 4 \( \mu s \).

**Total charge**

For a more detailed discussion, the charges measured in both parts of the signals, as discussed in the last section, are displayed in Fig. 5.18 for the electrons and in Fig. 5.19 for the ions. These measured charges in fact constitute a lower bound from the real charges, because for the MCP the maximum specified gain of \( 10^4 \) was assumed. Because of different sensitivity for electrons and ions and the different gain settings used, the absolute values of the electron and ion signals...
The signal of the Rydberg atoms decays because a large part of the population is converted to a plasma. Furthermore, the Rydberg atoms can be redistributed by collisions and spontaneous emissions to low excited states below the detection threshold of \( n = 37 \).

The total charge signals of the electrons and ions are plotted in Fig. 5.20 as a function of the delay time. Electrons and ions show a different behavior. The charge in the electron signal decreases continuously over time. This is for the most part because slow electrons from the plasma have a low detection efficiency, so they escape without detection and second because collisions and spontaneous decay to non-detectable Rydberg states.

The ion signal shows different behavior because the detection of low energy ions is much higher. After about 10 \( \mu s \), about the time needed for the generation of a plasma, the total charge starts to increase again to about 90% of the initial signal after which it starts to decay like the electron signal. This behavior can be explained if the redistributed low excited Rydberg levels that are not detectable, as in the simulation described in previous section, are ionized or excited to higher Rydberg states by collisions with electrons in the plasma.

For lower densities, where the interaction between the Rydberg atoms and plasma becomes less important, the decay of the Rydberg atoms can be estimated with spontaneous emission. The exponential decay time constant for the spontaneous decay to non-detectable states is extrapolated from the work of Oliveira et al.\[42\] for the 44d state. It is estimated to be about 90 \( \mu s \) and plotted in Fig. 5.20 as a dashed line. This timescale is of the same order of magnitude as the measured decay times.
Figure 5.19: The total charge of the ion signal above the initial Rydberg state (filled circles) and below the initial Rydberg state (open squares) as function of the delay time.

Figure 5.20: The total charge of the electrons (filled circles) and ions (open circles) as function of the delay time. The calculated exponential decay of the Rydberg atoms is shown as a dashed line.
5.4 Time resolved beam experiments

Time of flight

A measurement series done at different electric field pulse amplitudes can be used to make an estimation of the position of the cloud inside the rod structure. In Fig. 5.21 the time of flight of ions from the cloud of 44d Rydberg atoms to the detector is plotted against the input voltage of the μs high voltage pulser. At 60 V input it produces a peak voltage of about 1.5 kV.

The results of two GPT simulations are also shown in the figure. The field in the simulation is approximated by multiplying the calculated static field for an infinitely long rod structure, with a time dependent factor, obtained from a fit of the measured voltage pulse. The measured results are in reasonable agreement with the calculated time of flights for ions, corresponding to initial starting positions between −3.0 and −6.0 mm with respect to the center of the rod structure.

5.4.2 Conclusions

The time dependent measurements show clearly the generation of an ultra cold plasma, by exciting cold atoms in the MOT to the 44d Rydberg state. The spontaneous evolution from cold Rydberg atoms to plasma is observed under the right conditions. The measured plasma fractions as a function of delay time show similar qualitative behavior as earlier measurements and simulations in literature.

The difference between the electron and ion signal can be explained by the difference in detection efficiency for low energy particles. For better detection of these plasma particles, a
faster or custom shaped electric field pulse might be needed.

On the basis of time of flight measurements of ions the starting position is estimated to be between the $-3.0$ and $-6.0$ mm with respect to the center of the rod structure.

5.5 Beam imaging experiments

5.5.1 Phosphor screen images

A set of typical CCD images of the phosphor screen is shown in Fig. 5.22, measured for both electrons and ions. In the left part images are shown of the the electrons and ions created at a low density. The initial density of the atoms in the MOT is controlled by the repump laser intensity. At this low density, the interaction between the Rydberg atoms is too low to create a plasma. So the image created on the phosphor screen is only due to particles that were in the initial 44d Rydberg state or nearby lying state by redistributions.

Spatial distribution

In right part the images are plotted at higher initial density, so a plasma is formed. The images consists both of particles stripped from the plasma as well particles field ionized from the initial Rydberg state. The dimensions of the images on the screen are clearly larger than in the case of low initial density and comparable to the size of the phosphor screen (40 mm).
Figure 5.23: The results of two dimensional Gaussian fits to determine the widths of a series of ion MCP images as function of delay time. The total charge measured is shown as open circles, and the charge in the plasma fraction as filled circles.

Figure 5.24: The results of two dimensional Gaussian fits to determine the widths of a series of electron MCP images as function of delay time. The total charge of the signal is indicated by open circles.
A series of images made at different delay times are fitted with a two dimensional Gaussian distribution. The fitted widths, defined as the $1/\sqrt{\sigma}$ of the maximum value, are plotted in Fig. 5.23 and Fig. 5.24 as a function of delay time for the ions and electrons, respectively. Both axes of the fitted Gaussian distribution are shown, parallel to the rod structure, and one perpendicular. Also the total charge and the charge in the plasma fraction which were already discussed in previous section (see Fig. 5.18 and Fig. 5.19) are shown.

For both the electrons and the ions, the widths seem to follow the total charge of the signal. The ions that have an extra peak in the total charge also show an extra peak in both measured widths, while the electron follow the decreasing profile of the total charge nicely. In the ion series there is a clear offset of the curve in the perpendicular direction, while for electrons there is a less clear separation between them.

The initial width of the cold atom cloud has not been measured directly, but is estimated to be maximal 1 mm at most. There are three major effects that contribute to the increased width on the detector. First of all the random velocity distribution, characterized by its temperature, that causes ballistic expansion of the cloud.

Furthermore the space charge plays an important role. The plasma itself is charged, but still can be seen as quasi-neutral, because only a small fraction of the electron escapes. But when the stripping electric field is turned on and the plasma is pulled apart, the particles that travel to the detector experience a large repulsive force dependent on the charge density of the moving cloud. That increases the width on the detector.

The last effect is caused by the non-linear electric fields due to the rod geometry. In the parallel direction there is no electric field component, but in the perpendicular direction, it acts like a defocussing lens, that can also increase the size of the cloud on the detector.

The clear correlation between the total charge and the measured widths, show that the space charge term is probably the dominant term. Furthermore the offset of the width curves at the ions perpendicular to the rods might be explained by the extra defocussing effect that the ions experience because of the non-linear fields.

**Space charge expansion**

In Fig. 5.25 and Fig. 5.26 the measured widths of the ion and electron images are plotted as a function of the total charge. For the ions the curve shows a square root dependence which saturates at higher charges. The curve for the electrons does not show distinct saturation behavior. But the electrons are harder to interpret because of the effects of stray fields.

If the expansion is assumed to be solely due to space charge, it can be compared to a simple model of an expanding homogenous charged cloud, an approximation is given by:

$$r = r_0 \sqrt{1 + \left(\frac{t}{t_0}\right)^2},$$

with the time constant $t_0$ given by:

$$t_0 = \frac{1}{2} \sqrt{\frac{2\pi}{e}} \sqrt{\frac{m}{e} \sqrt{r_0^3}} \sqrt{\frac{r_0^3}{Q}},$$

where $r$ is the radius after time $t$, $r_0$ the initial radius on $t = 0$ and $Q$ the total charge. This approximation is only valid for $t >>> t_0$.

From this simple model also the qualitative square root behavior in the charge is found. If an initial distribution of a maximum radius of 1.0 mm is assumed with the minimal measured charge of 2 pC, this model gives approximately a five times higher width than measured.
Figure 5.25: The measured widths of the ion images parallel (filled circles) and perpendicular (open squares) to the rods as function of the total charge. The dashed line is a square root fit of the data for illustration.

Figure 5.26: The measured widths of the electron images parallel (filled circles) and perpendicular (open squares) to the rods as function of the total charge. The dashed line is a linear fit to the data.
The space charge forces might be much lower, because they are not created at one point in time, but are spread over some time delay, as in the measured time traces. To investigate the influence of the initial conditions on the resulting width on the screen, a set of simulations has been performed with GPT. But because the exact values of many parameters are unknown, such as the initial size and position of the cold atom cloud, many degrees of freedom exists.

By adding a timespread to the initial situation in the simulation, a reasonable agreement can be obtained for the parallel direction, but the perpendicular direction with a defocussing field, is off by about an order of magnitude. The effect of the non-linearity of the field gives a much larger divergence in the simulation then observed, if a position between the −6 and −3 mm is assumed that was estimated with the time of flight. So this simple idea, that the major contribution to the width increase is due to space charge, cannot explain the behavior in the perpendicular direction.

**Emittance**

The measured widths of the screen images can be used to make an upper estimate for the emittance and the source temperature. By using the widths plots vs charge, the width for very low charges can be found by extrapolation of the data. In this way the effect of space charge can be eliminated. If we assume that the expansion at 'zero' charge is only due to ballistic expansion, because of the initial temperature of the particle cloud, an upper limit can be calculated. For both the electrons as the ions, with a time of flight of 200 ns and 10 μs respectively, a temperature of about 40 K is found.

The electron temperature is in agreement with literature[43], but the ion temperature is an order of magnitude higher than the temperature reported[44]. This indicates that probably other effects besides the ballistic expansion play a role.

The temperatures can be used with Eq. (3.7) to get the initial emittance of the bunch:

\[
\epsilon = R \sqrt{kT/5mc^2}.
\] (5.12)

This results with an estimated radius of 1.0 mm in an initial emittance of \( \epsilon = 0.05 \) mm mrad. Performing this procedure on the measurement with maximum charge, an emittance of \( \epsilon = 0.5 \) mm mrad is found with \( R = 8 \) mm and corresponding temperature of \( T = 105 \) K. These values are upper limits so the real values are probably lower.

**5.5.2 Conclusions**

For the first time spatial images have been obtained from a phosphor screen of electron and ion beams originating from the a UCP. The exact dynamic process that causes the widths on the screen is unknown, but the measurements show that space-charge effects play a major role.

From these measurements an upper bound for the temperature can be estimated of 40 K for both the electrons as the ions. This implies an initial emittance of maximally \( \epsilon = 0.05 \) mm mrad, more than an order of magnitude better than all existing photo-emission electron sources.
Chapter 6

Conclusion and outlook

A stable rubidium MOT has been built that can contain more than $10^9$ atoms. The time dependent measurements show clearly the generation of an ultra cold plasma, by exciting cold atoms in the MOT to the 44d Rydberg state. The spontaneous evolution from cold Rydberg atoms to plasma is observed under the right conditions. The measured plasma fractions as a function of delay time show similar qualitative behavior as earlier measurements and simulations in literature.

For the first time spatial images have been obtained from a phosphor screen of electron and ion beams originating from the a UCP. The precise dynamics is not yet completely understood. Many initial conditions, such as the initial MOT distribution and position are not exactly known for the presented measurement series. Therefore it is hard to do a quantitative analysis of the results, but the measurements show that space-charge effects play a major role.

In the near future a series of new experiments can be done, in which the MOT distribution is measured simultaneously by two calibrated CCD cameras, so the initial distribution and position are known. Furthermore the MCP detector has to be calibrated, so an absolute measure for the collected charge can be obtained. With this extra information, a more realistic study of the results is possible by using for example GPT[23] simulations. If needed, measurements of Rydberg states can be done with a DC electric field, so the electric field characteristics can be estimated and compared to the field calculated with the Superfish[38] code.

For future experiments a new accelerator structure has been designed that fits in the current vacuum setup, see Fig. 6.1. It is designed for pulses up to 30 kV with a rise time of 10 ns. The accelerating structure has been built and is being tested. With this structure the truly bunched electron and ion beams can be created and prove the potential of the ultra cold beam source. Simulations performed with GPT suggest an emittance of 0.15 mm mrad can be obtained with a beam current of about 50 mA.
Figure 6.1: Schematic overview of the 30 kV accelerator structure. On the left the 3D view of the structure. On the right side the cross section of the accelerator.
Appendix A

Sector bending magnet

A sector bending magnet with a constant homogenous magnetic field inside and with parallel pole faces[31] is shown in Fig. 4.7. It has a focusing effect, so an image object relation exists between the plane \( l_1 \) from the entrance pole, to the image plane \( l_2 \) from the exit pole. These lengths can be expressed in terms of the chosen magnet properties, \( \phi_0, \rho_0 \) and \( M_y \):

\[
l_1 = \frac{\cos \phi_0 - (1/M_y)}{\sin \phi_0}, \quad (A.1)
\]

\[
l_2 = \frac{\cos \phi_0 - M_y}{\sin \phi_0}, \quad (A.2)
\]

with \( \phi_0 \) the sector angle, \( \rho_0 \) the radius of the magnet and \( M_y \) the magnification. The length of the trajectory of a reference particle from one to the second plane is thus given by:

\[
L_{magnet} = \rho_0 \left( \phi_0 + \frac{2 \cos \phi_0 - M_y - (1/M_y)}{\sin \phi_0} \right). \quad (A.3)
\]

Due to its simple form, the transfer matrix is used that images from the object plane directly on the image plane. The coordinates themselves are assumed to be rotated, so the reference particle is still moving in the \( z \)-direction after the sector magnet. In the needed \( 6 \times 6 \) matrix form it is defined as:

\[
T_{magnet} = \begin{pmatrix}
1 & (x|p_x) & 0 & 0 & 0 & 0 \\
0 & 1 & 0 & 0 & 0 & 0 \\
0 & 0 & (y|y) & 0 & (y|\delta) & 0 \\
0 & 0 & (p_y|y) & (p_y|p_y) & (p_y|\delta) & 0 \\
0 & 0 & 0 & 0 & 1 & 0 \\
0 & 0 & 0 & 0 & 0 & 1
\end{pmatrix}, \quad (A.4)
\]

with in the \( x \)-direction only drift space, because no forces are exerted in that direction, so:

\[
(x|p_x) = \frac{1}{p_{z0}} L_{magnet}. \quad (A.5)
\]

And in \( y \)-direction the bending itself, with:

\[
(y|y) = M_y, \quad (A.6)
\]

\[
(y|\delta) = \rho_0 (1 - M_y), \quad (A.7)
\]
\( (p_y|y) = \frac{1}{\rho_0} \sin \phi_0 p_{x_0}, \)  

(A.8)

\( (p_y|p_y) = \frac{p_{x_0}}{M_y}, \)  

(A.9)

\( (p_y|\Delta) = p_{x_0} \sin \phi_0. \)  

(A.10)
Appendix B

Transfer matrix streaking setup

The total transfer matrix of the streaking setup is simple the multiplication of both the matrices of the streaking and the magnet part, \( T_{sys} = T_{mag} \cdot T_{str} \). To be complete, the result is given here:

\[
T_{sys} = \begin{pmatrix}
1 & (x|p_x) & 0 & 0 & (x|\delta) & (x|z_0) \\
0 & 1 & 0 & 0 & (p_x|\delta) & (p_x|z_0) \\
0 & 0 & (y|y) & (y|p_y) & (y|\delta) & 0 \\
0 & 0 & (p_y|y) & (p_y|p_y) & (p_y|\delta) & 0 \\
0 & 0 & 0 & 0 & 1 & 0 \\
0 & 0 & 0 & 0 & 0 & 1
\end{pmatrix}, \tag{B.1}
\]

with the particular matrix elements of the \( x \)-direction given by:

\[
(x|p_x) = \frac{L_{str} + L_{mag}}{p_{z_0}}, \tag{B.2}
\]

\[
(x|\delta) = -\frac{1}{2}q \frac{V_{max} L_{mag} \tau}{h p_{z_0}}, \tag{B.3}
\]

\[
(x|z_0) = -\frac{1}{2}q \frac{V_{max} \tau}{h} \left( \frac{\tau}{p_{z_0}} + \frac{2m L_{mag}}{p_{z_0}^2} \right), \tag{B.4}
\]

\[
(p_x|\delta) = \frac{1}{2}q \frac{V_{max}}{h} \tau, \tag{B.5}
\]

\[
(p_x|z_0) = -q \frac{V_{max}}{h} \frac{m}{p_{z_0}}, \tag{B.6}
\]

And the elements for the \( y \)-direction given by:

\[
(y|y) = M_y, \tag{B.7}
\]

\[
(y|p_y) = M_y \frac{L_{str}}{p_{z_0}}, \tag{B.8}
\]

\[
(y|\delta) = \rho_0 (1 - M_y), \tag{B.9}
\]

\[
(p_y|y) = -\frac{1}{\rho_0} p_{z_0} \sin \phi_0, \tag{B.10}
\]
\[(p_y|y) = \frac{L_{str} \sin \phi_0}{\rho_0} + \frac{p_{z0}}{M_y} \tag{B.11}\]

\[(p_y|\delta) = p_{z0} \sin \phi_0. \tag{B.12}\]

where \(L_{str}\) is given by Eq. (4.11) and \(L_{mag}\) by Eq. (A.3).
Bibliography


