THE THERMIONIC DIODE AS A SENSITIVE SPECTROMETER: 
ITS CONSTRUCTION, THEORY AND APPLICATIONS

A dissertation submitted to the
SWISS FEDERAL INSTITUTE OF TECHNOLOGY ZURICH

for the degree of
Doctor of Natural Sciences

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ADAG Administration & Druck AG
Zurich 1989
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CURRICULUM
The thermionic detector or diode is the key element in the experimental set-up for the measurement of parity violation in atomic Cesium currently under way in our laboratory. Although this device has been known for a long time its design remained purely empirical and its electrical characteristics understood only qualitatively.

The present thesis describes the construction of diodes especially suited for ultra-rare laser-induced transitions (such as the ones in the parity experiment). These diodes are exceptionally long (up to 30 cm !), have indirectly heated cathodes and were designed with a high degree of symmetry in order to achieve field-free interaction regions.

This thesis furthermore presents the first quantitative theory of the current vs. voltage characteristics, and compares it with measurements; very good agreement results. A comparison of the theoretically estimated gain with carefully measured values gives order-of-magnitude agreement. The simple model of looking at the thermionic diode as an ion capacitor is shown to describe its gain vs. chopping frequency behaviour very accurately.

An extension of the theory to the presence of axial magnetic fields is also given. It was successfully tested in an analytically tractable situation. The experimental study of the behaviour in axial magnetic fields led us to discover a method that enables one to detect fields down to 5 mGauss within the diode.

As an application of the diode to an actual physical problem we investigated (parity-conserving) circular polarization asymmetries in the same resonant 2-photon transition as will be used to study parity violation, i.e. the chain 6S \( \left( \lambda_1 \right) \rightarrow \).
Abstract

7S (\lambda_2) \rightarrow 15P in atomic Cs. Two counterpropagating laser beams \lambda_1, \lambda_2 with circular polarizations q_{1,2} = \pm 1 enter the diode through windows and interact with Cs vapour in it. A special "Stokes Monitor" was developed to monitor the intensity and the purity of polarization of light behind the windows, i.e. **within** the diode. The measured asymmetries vs. light intensity are compared to the predictions of a rate equation model proposed by J. Hoffnagle; i.e. a model which includes saturation and collisional population redistribution among the sublevels of the final Rydberg state (with simplifying assumptions) is tested.
KURZFASSUNG

Der thermionische Detektor, kurz Diode genannt, ist das zentrale Element im experimentellen Aufbau zur Messung der Paritätsverletzung in atomarem Cäsium. Obwohl diese Art von Detektor seit langer Zeit bekannt ist, blieb sein innerer Aufbau empirisch bestimmt und sein elektrisches Verhalten nur qualitativ verstanden.

Die vorliegende Arbeit beschreibt Dioden, die speziell für sehr schwache, laser-induzierte Übergänge (wie z.B. im Paritätsverletzungsexperiment) konstruiert wurden. Diese Dioden sind außergewöhnlich lang (bis zu 30 cm!) und mit indirekt geheizten Kathoden versehen. Bei der Konstruktion wurde auf hohe Symmetrie geachtet, um die Wechselwirkungs-Region möglichst feldfrei zu halten.


Eine Erweiterung der Theorie, welche zusätzlich axiale magnetische Felder berücksichtigt, wird vorgestellt und in einem analytisch behandelbaren Fall mit dem Experiment verglichen. Im Laufe der Untersuchungen des Verhaltens der Diode im Magnetfeld haben wir eine Methode gefunden, mit der man im Innern der Diode selbst Magnetfelder bis herunter zu 5 mGauss messen kann.
Als Anwendungsbeispiel wurde die Diode in einem Experiment zur Messung paritätserhaltender Zirkularpolarisations-Asymmetrien eingesetzt. Benutzt wird dabei der gleiche resonante 2-Photonen Übergang wie im eigentlichen Paritäts-Experiment, d.h. die Übergänge $6S (\lambda_1) \to 7S (\lambda_2) \to 1S$ in atomarem Cäsium. Zwei gegeneinander laufende Strahlen $\lambda_1, \lambda_2$ mit Zirkularpolarisationen $q_1(\lambda_1) = \pm 1$ wechselwirken mit Cäsiumdampf in der Diode nach Durchgang durch optische Fenster. Um die Intensitäts- und Polarisations-Abweichungen innerhalb der Diode messen zu können, haben wir einen "Stokes Monitor" entwickelt. Die gemessenen Asymmetrien in Funktion der Lichtintensität werden mit den Voraussagen eines von J.Hoffnagle vorgeschlagenen Ratenmodells, welches Sättigung und stossinduzierte Populationsumverteilungen (mit vereinfachenden Annahmen) im Rydberg-Zustand berücksichtigt, verglichen.
INTRODUCTION

The main aim of our group, the measurement of parity violation on the M1 transition $6S_{1/2} \rightarrow 7S_{1/2}$ in cesium in the absence of external electric and magnetic fields, is conceptually a very simple experiment. It was first proposed by Bouchiat and Bouchiat in 1974 [Bou 74,75], and consists in measuring a circular dichroism of the order of $10^{-4}$, i.e. a small difference of the absorption coefficients for left- and right-circularly polarized light. The signal of this transition for the two polarization states ($\xi = +1, -1$) of the light is given by

$$S(\xi) \sim |M1|^2 + 2\xi \text{Im}(M1 \cdot E1^*) .$$

M1 is the real (parity allowed) magnetic dipole transition amplitude, whereas $E1_{pv}$ represents the purely imaginary (parity violating) electric dipole amplitude due to the weak interaction. Thus parity violation can be observed by an asymmetry measurement:

$$A_{pv} = \frac{S(+1) - S(-1)}{S(+1) + S(-1)} \approx 2 \text{Im}(E1^*)/|M1| .$$

The reasons for which this effect has not yet been observed under field-free conditions are the extremely small oscillator strength ($f \approx 4 \cdot 10^{-18}$, i.e. $\sigma_{abs} \approx 10^{-28} \text{cm}^2$) of the M1 transition, the laser-induced fluorescence of the Cs$_2$ near the M1 detection wave length ($1.47 \mu\text{m}$, see Fig. 0.1) and the huge black-body background from a Cs oven at $\approx 200 ^\circ\text{C}$).
The first successful measurements of the field-free M1 transition were achieved in our laboratory by fluorescence spectroscopy [Hof 81], [Hof 82]. Because of the extremely large Cs$_2$ background and the poor signal to noise ratio of this method a novel detection scheme for 7S atoms had to be developed.

The key element in this new scheme is the thermionic diode as described in the Chapters I and II. The detection of the 7S atoms is achieved by pumping them to a Rydberg state (e.g. 15 P$_{3/2}$) that is readily ionized by collisions with ground state atoms [Wei 84], see Fig. 0.2. This resonant Doppler-free detection scheme has already been tested successfully [Her 86;1].
Further investigation of this detection scheme was necessary since the second transition (E1) is easily saturable and the polarization correlations of the two laser beams had to be studied. This is described in detail in Chapter III.

Fig. 0.2: Term scheme of the Cs atom with transitions used in the Doppler-free detection scheme.
During the time the thermionic diodes were developed in our laboratory, they were used in several spectroscopic applications involving Cesium Rydberg states, viz.

- Measurement of the doublet intensity ratios of the 7S_{1/2} \rightarrow nP_J series in Cesium, the 7S_{1/2} states being excited by electron impact (Franck-Hertz effect) [Hof 85;2].

- Doppler-free two-photon spectroscopy of the 8S state of Cesium was used to measure the h.f. splitting and the term energy of that state [Her 85].

- Measurements of tensor polarizabilities of the 40D and 60D states in Cesium, the Rydberg states being populated by two-photon transitions (\omega_1 \neq \omega_2, no resonant intermediate state) [Her 86;4].

- Stark spectroscopy of the forbidden 6S_{1/2} \rightarrow nP_J two-photon transitions in an electric field as a technique to measure small electric fields [Her 86;5].

- Collisional ionisation probabilities of highly exited S_{1/2}, D_{3/2} and D_{5/2} of Cesium and measurement of the doublet intensity ratios of the 6S_{1/2} \rightarrow nD_J two-photon transitions (\omega_1 \neq \omega_2, no resonant intermediate state) [Her 86;2].
The thermionic diode - a powerful detector in spectroscopy

The thermionic diode, working in the "space-charge-limited" mode, was found to be an efficient ion detector already in 1923. K.H. Kingdon [Kin 23] saw a significant deviation in the voltage-current characteristic of his gas-filled diodes from the vacuum behaviour for voltages above the ionisation potentials of neon and mercury vapour. The theory he developed was based on a suggestion of I. Langmuir that these deviations were due to trapped positive ions. Independently G. Hertz [Her 23] used this space charge reduction effect to measure the ionisation potentials of helium and argon. He produced the ions by injecting an electron beam into an electrostatically shielded box containing a heated tungsten filament.

In 1925 the thermionic diode was used the first time as a detector in spectroscopy by F.L. Mohler et al. [Moh 25,26]. They detected with rather poor resolution some lines of the cesium principal series.

As recently as 1964 Popescu et al. rediscovered this old technique and they improved it in the subsequent years. From the early 1970's, when tunable dye lasers became available, the thermionic diode has become a powerful detector in laser spectroscopy; for a review, see an article of K. Niemax [Nie 85].

Our diodes were designed and constructed as detectors for two-photon excitation experiments and as the key elements in a cesium PNC-experiment.
I.1 Construction and assembly of the thermionic diodes

Having done some measurements with "standard" thermionic detectors such as described in [Nie 85], we realized that we had to improve the design for our purposes. To get a detector working on its whole length under optimal bias conditions (i.e. without a voltage drop along the cathode), we decided to use indirectly heated cathodes. (The advantages of this were already mentioned by D.C.Thompson [Tho 82]).

![Fig. 1.1: Schematic drawing of the indirectly heated cathode.](image)

Fig. 1.1 shows such a cathode schematically. To minimize unwanted magnetic fields in the diode, we adopt a bifilar heating (tungsten wire, diameter 150 \(\mu\)m). As the cathode surface proper we use an other tungsten wire (diameter 50 \(\mu\)m) tightly wound around a \(\text{Al}_2\text{O}_3\) ceramic bi-bore tube (diameter 0.9 mm) in a self-fixation technique (i.e. no glue or screws are used to hold the winding on the tube, see Fig. 1.1).

To collect as many ions as possible from the interaction volume, we designed double diodes as shown in Fig. 1.2; more details about the construction are presented in Appendix A. The two diodes "see" the electrostatically shielded interaction volume I through grids made of 50 \(\mu\)m tungsten wires wound on a rectangular molybdenum support with 0.3 mm spacing. To characterize this type of diode with respect to a simple coaxial one the mean anode radius \(<r_a>\) has to be calculated (see Appendix B); for Fig. 1.2, \(<r_a> \approx 14\) mm.
Fig. 1.2: Schematic cross section through the diode:

$C_1, C_2$ cathodes consisting of:
- $\text{Al}_2\text{O}_3$ tube (diam. 0.90 mm) with two bores
- $f$ heating wires (tungsten, diam. 150 $\mu$m)
- $W$ emitting surface (tungsten wire, diam. 50 $\mu$m, tightly wound around the ceramic tube in a self fixation technique);

A anode, electrically connected to the grids $G_1$ and $G_2$;

I interaction volume;

M Macor support with
- $\text{Al}_2\text{O}_3$ isolating tubes
- $h_1, h_2$ counterpropagating heating currents (tungsten rods diam. 1.0 mm)

P Pyrex housing (outer diam. 74 mm).
The whole double diode is enclosed in a pyrex cell having a side arm containing cesium metal. The longest diode constructed by us has an active length of 30 cm and consists of three double diodes in series.

Great care was taken in the design of the electrical feed-throughs to assure a long detector live time (over two years were achieved). The problem to be solved was to protect the feed-throughs from the reduction of the tungsten-oxide layers on the tungsten-glass interface by hot cesium vapour. The solution is shown in Fig. 1.3. First, a feed-trough is prepared by the standard technique; next, after cleaning the tungsten surfaces, a thin tube (of the same glass as used in the first step) with an inner diameter well matched to the tungsten rod is carefully molten on the inward end. This step is done in a hydrogen atmosphere to prevent oxidation.

![Fig. 1.3: Schematic drawing of Cs resistant feed-through.](image)

The pyrex cell is provided at its ends with plane parallel windows. These windows are optically contacted to a support tube to minimize the birefringence effects standard sealing techniques produce (Fig. 1.4). A He-Ne laser is used to adjust the windows before sealing the support tubes to the diode envelope.

A further reduction of the residual window birefringence by a factor of two was achieved by taking for the windows and the support tubes glass (Schott-glass Duran 50) from the same batch.
The measured residual birefringence is ≈ 10 mrad.

After assembly, the double-diode is connected to a vacuum pump and baked out for several days at about 300 °C until the pressure falls below \(10^{-8}\) Torr for at least one day. Then the filaments are heated for several hours with a current of 0.6 A to further clean the cathode surfaces. Thereafter about 0.2 g of cesium is destilled into the side arm and the diode sealed off from the vacuum system.

**I.2 Burning-in of the diode:**

Before the diode can serve for spectroscopic measurements the cathode surface has to be activated by a so called "cooking procedure". Four to six weeks were found to be necessary for this. We suppose that some slow crystallisation changes on the cathode surface and perhaps some further reduction by the cesium vapour occur in this period.

Our diodes are operated in a two-chamber oven, with temperatures stabilized to ± 1 °C for the diode chamber and to ± 1/4 °C for the side arm chamber.
The anode current is measured by a current to voltage converter and the variable anode voltage is provided from a stabilized and adjustable source. For the schematics of the cathode heating and the anode-cathode circuits see Appendix C.

The measured temperature vs current characteristic of our indirectly heated cathodes is provided in Appendix D. With this characteristic we can determine the temperature of the cathode surface $T_C$ with an estimated accuracy of $\pm 20 \, \text{°K}$.

The details of our cooking procedure are as follows:

a) First the cesium vapour pressure is set to about 10 mTorr and the vapour temperature to about 200 °C.

b) The filament temperature is set to $\approx 1500 \, \text{°K}$ (i.e. heating current $\approx 1.1\, \text{A}$; the max. temperature allowed for $\text{Al}_2\text{O}_3$ ceramics is $\approx 2100 \, \text{°K}$).

c) The anode current is monitored on a chart-recorder. To check the behaviour of the diode, the ion life time is measured once a day. The ions are produced through two-photon ionisation of cesium-dimers (Cs$_2$) by focussing some

![Chart recorder output of the diode anode current vs. time with $\approx 20 \, \text{mW Ar}^+\text{-laser light switched on and off. (time increases from right to left!)}$](image_url)
10 mW from the Argon-Ion Laser (all lines) into the interaction region. These ions cause an increase of the anode current $\Delta i_a$ by space charge reduction (see Chapter II). A typical chart-recorder output is shown in Fig. 1.5.

For these measurements the filament current is first switched off for about one hour to allow the cesium to rebuild a layer on the cathode surface, and then the filament current is set to 0.7 to 0.8 A and the anode-cathode voltage to about 0.5 V. If $\tau_{\text{off}}$, the mean live-time of the ions (for definitions see Fig. 1.5), is about 10 sec or more, the diode is said to be ready.

I.3 Characteristic and gain measurements:

After the "cooking procedure" current-voltage characteristics of the diode are measured to find the optimum operating conditions. Fig. 1.6 shows a plot of a typical characteristic with the additional measurements of signal current ($\Delta i_a$), mean ion life-time ($\tau_{\text{off}}$) and the collection time ($\tau_{\text{on}}$). These measurements were done with a prototype double diode (geometry slightly different from the diodes used now, see Appendix B).

The last curve shown on Fig. 1.6 is the signal current to mean ion life-time ratio. This ratio shows to be more or less constant for anode-cathode voltages from 0 to about 1.5 V; i.e. the anode current change is proportional to the number of ions in the diode:

$$\frac{\Delta i_a}{q\Gamma^+\tau_{\text{off}}} = \text{const.}$$
Fig. 1.6: Diode test measurements at a Cs pressure of 
\approx 4 \text{ mTorr}.

- current-voltage characteristics
- signal current
- ion collection time $\tau_{on}$
- mean ion life time $\tau_{off}$
- signal current to ion life time ratio
where \( \Gamma^+ \) is the production rate of positive ions. If we define the diode gain as the ratio of anode current change to ion current change we get:

\[
G_s = \frac{\Delta I_a}{\Delta n^+} = \frac{\tau}{\tau_{\text{off}}}
\]

With this definition the ion collection efficiency is included in \( G_s \). By setting

\[
Q^+ = \rho^+ \Delta V = qn^+
\]

for the total positive charge in the diode,

\[
\tau_{\text{on}} = \tau_{\text{off}} = \tau
\]

neglecting the measured differences between \( \tau_{\text{on}} \) and \( \tau_{\text{off}} \) and

\[
i_o = qr^+
\]

for the ion current into the diode we get the differential equation for a modulated ion current:

\[
(1.3.1) \quad \frac{dn^+}{dt^+} = -\frac{n^+}{\tau} + \Gamma^+ \cos(\omega t)
\]

and the solution

\[
(1.3.2) \quad n^+(\omega) = \frac{\Gamma^+ \tau}{\sqrt{1 + \omega^2 \tau^2}} \cos(\omega t - \arctan(\omega \tau))
\]

i.e. a \( 1/\omega \) dependence for \( \omega \tau \gg 1 \).

Equation (1.3.1) suggests the diode to be an ion equivalent RC network with an external current source (\( \tau = RC \)):

\[
\begin{align*}
\frac{dq}{dt} &= -\frac{Q}{RC} + i_o \cos(\omega t)
\end{align*}
\]
Since the ion mobility is much less than the electron mobility, the anode current response is assumed to be instantaneous $\Delta i_a(\omega) \sim qn^+(\omega)$. Thus we get for the diode gain:

\[
G(\omega) = G_s \frac{1}{\sqrt{1+\omega^2\tau^2}}
\]

\[
G_s = \frac{\Delta i_a(\omega \rightarrow 0)}{qn^+(\omega \rightarrow 0)} \cdot \tau
\]

We have used two independent techniques to measure diode gains.

i) Single photon ionization of Cs for measuring the gain at $\omega = 0$ ($G_s$). A high pressure xenon lamp and an interference filter (3'125 Å) are used to produce a UV beam. Switching this beam on and off produces an anode current change which can be recorded on a chart. Using the photoionization cross section from Hofsäess [Hof 79] $G_s$ can be calculated:

\[
G_s = \frac{\Delta i_a}{qn^+} \quad \text{Typical value: } G_s = 8,3(2,5) \cdot 10^8
\]

The specified uncertainty is mainly due to the power (flux) measurement.

ii) Two-photon transitions in Cs (6S $\rightarrow$ 15D) followed by collisional ionization [Wei 84], [Her 86;2] are used to measure gains at $\omega \neq 0$ ($G(\omega)$). Two laser beams counterpropagating through the interaction volume of the diode are used here. The frequency of one of the dye lasers is scanned over the transition. The anode current $\Delta i_a(\omega)$ is measured by phase sensitive detection technique. $G(\omega)$ is determined by two successive measurements, one with filament on, the other with filament off, i.e. with gain switched on and off.
The trapping time $\tau$ is measured as shown in Fig. 1.5. With $\tau = 2.5$ sec, equation (1.3.2) and an accuracy of 10 to 15 % for the "filament-off" measurement, we found a typical value: $G_0 = 1.2(0.2) \cdot 10'$

The two measurements of $G_0$ agree. Furthermore, the measured frequency dependence of the signal current follows equation (1.3.2) exactly; i.e. our simple model describes the dynamic properties of the thermionic diode very well.

### I.4 Background and noise of the thermionic detector

The sensitivity of any detector is limited by noise. This is especially important with the thermionic diode since the signal current is only a small change on a large quiescent current. We distinguish between intrinsic (detector dependent) and external noise sources:

**Intrinsic noise sources:**

i) electronic noise from current-voltage converter and bias voltage  

ii) noise current from the finite cell resistance $R_i$; $R_i$ varies with the Cs pressure and hence this noise contribution is related to the stability of the reservoir temperature ("excess" noise)

iii) shot noise of the quiescent anode current  

\[
(i_a \approx 350 \mu A; \quad i_{\text{shot}} = \sqrt{2q i_a \Delta v} \approx 11 \text{ pA/Hz} }
\]

iv) fluctuations of the emission conditions at the cathode ("flicker" noise)
v) The contribution from Cs ionisation at the hot cathode surface leads to a small ion current (estimated ion rate for our diode is $\Gamma_c \approx 680 \text{ /sec}$). The shot noise of this current can be neglected compared to the other noise sources.

**External noise sources:**

vi) shot noise of the ion current: $i_{\text{shot}}^+ = \sqrt{2q i^+ \Delta v}$

vii) fluctuations in the ion production rate due to power (intensity) instabilities.

The signal considered so far ($\Delta i_a$ due to two-photon ionisation of $\text{Cs}_2$) is unwanted background in our measurements involving two counter-propagating laser beams with different wavelengths ($\lambda_1 \approx 540 \text{ nm}$, $\lambda_2 \approx 831 \text{ nm}$). Since the $\lambda_2$ beam by itself produces no "background ions", we get rid of the two-photon background of the $\lambda_1$ beam by chopping the $\lambda_2$ beam and using phase sensitive detection (lock-in amplifier).

Nevertheless the fluctuations of the ion production rate of the two photon background process due to power fluctuations and the shot noise of this ion current contribute still to the noise on the signal. This requires that the power of the $\lambda_1$-laser must be highly stabilized.

The measured contributions of the individual noise sources i) and ii) to the total intrinsic noise current ($\Delta i_e^*/\sqrt{\text{Hz}}$) of the diode at typical working conditions (cathode temperature $\approx 800 \text{ K}$, Cs pressure $\approx 10 \text{ mTorr}$, $V_a = 0.5 \text{ V}$) are shown in Fig. 1.7 (lines are drawn to guide the eye). In this figure input equivalent noise currents are given by taking into account the feedback resistor (10 kΩ) in the current to voltage converter. This noise data were obtained with our "longest" diode (active length $\approx 30 \text{ cm}$).

The total external noise current is given by
\[ (1.4.1) \quad \Delta i_a^*(\omega; \text{ext}) = G(\omega) \cdot i^+_{\text{shot}} \sqrt{1 + 2 \Gamma_P^+ \left( \frac{\Delta P(\omega)}{P} \right)^2} \]

where \( \Gamma_P^+ \), the ion production rate, is assumed to be proportional to the square of the intensity (\( \Gamma_P^+ \sim P^2 \) for background processes) and \( \frac{\Delta P}{P}(\omega) \) the relative power stability at frequency \( \omega \).

![Intrinsic diode noise at typical working conditions](image)

**Fig.1.7:** Intrinsic diode noise at typical working conditions:
- cathode temperature \( \approx 800 \) K (filament hot)
- anode-cathode voltage = 0.5 V

We define the **sensitivity** of the thermionic diode as the ion current needed to produce an anode current change equal to the total intrinsic anode current noise, thus the **noise equivalent ion rate** is
Since $\Delta i_e^\star(\omega)$ and $G(\omega)$ show the same $1/\omega$ dependence down to $\approx 10$ Hz (see Fig. 1.7), the signal to noise ratio ($S/N$) of our thermionic diodes is expected to be constant for chopping frequencies down to that frequency. This was confirmed experimentally; the chopping frequency generally used in our measurements was $v_{ch} \approx 8$ Hz.

With the noise current from Fig. 1.7 and (1.3.2) we find for the sensitivity of our diodes:

$$\text{NEIR (2\pi \cdot 8 \text{ Hz})} \approx 2 \cdot 10^4 \text{ ions/sec}$$

This figure could be reduced only by a factor of less than five by using cathodes with less "flicker" noise, since this kind of noise contributes most to the total intrinsic noise current. D.Popescu et al. [Pop 69] estimated for their diodes a sensitivity of 1 to 100 ions/sec from the shot noise of the anode current ($\approx 3$ \text{ \mu A}) and did not mention at all the other sources of noise.
1.5 Behaviour of the diode in an axial magnetic field

Our main goal, the determination of parity non-conservation (PNC) in atomic Cesium by the measurement of circular dichroism requires a field-free interaction volume in the diode. Since an axial magnetic field of the order of 0.1 Gauss causes a circular dichroism of the same order of magnitude as the PNC effect proper, the whole diode is shielded with a \( \mu \)-metal tube (residual earth field inside the tube without diode \( \approx 3 \) mGauss, from flux-gate measurement). To compensate possible residual axial fields from cathode heating currents (geometrical imperfections) a solenoid (17.1 Gauss/A) is installed between diode and magnetic shield. This solenoid was used to investigate the behaviour of the diode in axial magnetic fields.

An anode current vs. B-field characteristic \((i_a(B))\) is obtained by varying the solenoid current and recording the anode current on a chart (Fig. 1.8a with \( V_a = -0.5V \); Fig. 1.8b with \( V_a = +2.25V \)). While the diode is normally operated at \( V_a > 0 \), the theory for the diode in an axial magnetic field as presented here (Chapter II) is easily tested at \( V_a < 0 \). The decrease of anode current (Fig. 1.8a) with increasing magnetic field is well reproduced by our theory predicting a purely \( \exp(-aB^2) \) behaviour. Upon inspecting the \( i_a(B) \) characteristics more closely, we discovered "fine structures", i.e. reproducible departures from the exponential behaviour. While these are not yet fully explained, we can exploit them to measure the mean residual axial magnetic field within the active volume of the diode and thus obtain an upper bound for the magnetic field in the interaction volume. The "fine structure" (e.g. the "bump" near \( B = 0 \), Fig. 1.8b) emerges most clearly when the Cs pressure is reduced, presumably because it is washed out by collisions, and is emphasized when one plots \( \frac{d}{dB} i_a \) vs. \( B \) (see Fig. 1.8c). For field measurements, the diode is therefore operated at Cs pressures about
200 times lower than the "standard" pressure of \( \approx 10 \) mTorr and at anode cathode voltages from \(-1.0\) to \(+2.5\) V.

![Diagram](image1)

![Diagram](image2)

Fig. 1.8: Anode current vs. B-field characteristics at 0.05 mTorr Cs pressure:

a) at \( V_a = -0.5 \) V  
   b) at \( V_a = +2.25 \) V

The derivative of \( i_a(B) \) is recorded by superposing on the scanned solenoid current a modulation current with a small amplitude at \( \approx 80 \) Hz and detecting the in-phase component of the anode current with a lock-in amplifier (see Fig. 1.8c, which shows the derivative of Fig. 1.8b). By determining the
difference of the zero-crossings of the solenoid current and of the derivative $d_i_a(B)/dB$ (Fig. 1.8c) we achieved a sensitivity of $\approx 5 \text{ mGauss}$.

Fig.1.8c) $d_i_a(B)/dB$ vs. B-field characteristic of Fig.1.8b)
CHAPTER II

The thermionic detector - theoretical considerations

II.1 Measured voltage-current characteristic of the thermionic diode and the Langmuir-Child law

To compare a measured voltage-current characteristic as in Fig. 1.6 with the 3/2-power law of Langmuir [Lan 13] the finite cell resistance $R_1$ has to be considered ($R_1 = 0(100 \, \text{k\Omega})$). The measured and corrected (for the current through $R_1$) characteristic is shown in Fig. 2.1. The full line shows the characteristic as predicted by I. Langmuir, viz. the anode current per unit length ($\Delta z$):

$$
\frac{i_a}{\Delta z} = \frac{2}{9} \left( \frac{2e}{m} \right)^{1/2} \frac{va^{3/2}}{ra^2} ,
$$

(2.1.1)

where $V_a$, $r_a$ denote anode-cathode voltage and anode radius respectively, and $\beta^2$ is a function (see [Lan 23;2]) of the ratio anode radius to cathode radius. With $r_a = 1.60 \, \text{cm}$ and $\beta^2 (r_a/r_c = 32) = 1.09$ this yields

$$
\left( \frac{i_a}{\Delta z} \right)^{2/3} = \left( \frac{8}{81} \frac{e}{m} \frac{1}{\beta^4} \right)^{1/3} V_a = 0.41 \cdot 10^{-3} \, (\text{A/cm})^{2/3} (V_a/V)
$$

From Fig. 2.1 we see that the slope calculated from the 3/2-power law corresponds very well to the asymptotic slope of the measured characteristic; this asymptote, as opposed to Langmuir's prediction, has however finite (negative) intercept on the $V_a$ axis. This disagreement is due to Langmuir's original assumption that the electrons are emitted at the cathode surface with zero velocity. Subsequently Langmuir
[Lan 23; 1] gave an approximate solution for a maxwellian distribution of the emitted electrons:

\[
\frac{i_a}{\Delta z} = \frac{2}{9} \left( \frac{e}{m} \right)^{1/2} \frac{1}{r_a b^2} \left[ V_a - V_b + \frac{1}{4} V_o \left( \log \left( \frac{V_a}{\lambda V_o} \right) \right) \right]^{3/2}
\]

where \( eV_o = 3kT_c/2 \) is the mean initial kinetic energy, \( \lambda \) a constant between 1 and 2 to be determined experimentally, and \( V_b \) the magnitude of a potential minimum between cathode and anode. This formula cannot represent our data since it allows no negative anode-cathode voltages.

Fig. 2.1: Measured diode characteristic with prediction of the Langmuir 3/2 power law (-----) and the modified 3/2 power law (----).
To get a relation that fits the asymptotic slope of the measured characteristic we modify Langmuir's original theory [Lan 13] by introducing a new parameter $V_m$ on the basis of the assumption that all the electrons reaching the anode had a unique energy $qV_m > 0$ at the cathode surface. Energy conservation yields:

$$v = \sqrt{\frac{2 e}{m} (V - V_m)} ,$$

hence we have to replace $V_a$ in (2.2.1) by $V_a - V_m$

\begin{equation}
\frac{ia}{az} = \frac{2}{9} \left[ \frac{2e}{m} \right]^{1/2} \frac{(V_a - V_m)^{3/2}}{r_a B^2} \tag{2.1.2}
\end{equation}

or

$$\left[ \frac{ia}{az} \right]^{2/3} = 0.41 \cdot 10^{-3} \ (A/cm)^{2/3} (V_a - V_m)/V$$

This means that for $V_a \leq V_m$ there is no anode current at all. In Fig. 2.1 formula (2.1.2) is represented by the broken line; $V_m$ is found to be $\approx -710 \ mV$.

Although Eq. (2.1.2) describes the asymptotic behaviour very well, we need a more accurate calculation of the potential distribution in the diode in order to be able to estimate the change in anode current by trapped positive ions. This calculation is presented in detail in the next two sections.
II.2 A new approach for the calculation of the potential \( \Phi(r) \) and the voltage-current characteristic \( i_a(V_a) \) of the thermionic diode

One main input for this calculation is the existence of a potential minimum somewhere between cathode and anode.

There is already experimental evidence for the existence of such a potential-hill for the emitted electrons: It is well known (see for example [Lan 13]) that the anode current reaches its saturation value only at some positive potential difference between anode and cathode. Furthermore we can easily prove that a potential minimum exists between cathode and anode by realizing that electron emission creates a positive charge density \( \rho_c \) within the cathode. Gauss's law yields for the electric field at the cylindrical cathode surface

\[
\mathbf{E}(r_c) = 2\pi r_c \rho_c \mathbf{\hat{r}} = -\frac{\partial \Phi(r)}{\partial r}\bigg|_{r=r_c}.
\]

Since \( \rho_c > 0 \), one has

\[
(2.2.1) \quad \frac{d\Phi(r)}{dr}\bigg|_{r=r_c} < 0.
\]

If we choose \( V_a > V_c \) and assume a smooth potential dependence, we get for the potential slope at the anode

\[
(2.2.2) \quad \frac{d\Phi(r)}{dr}\bigg|_{r=r_a} > 0.
\]

Thus there exists a potential minimum at some radius \( r_b \) between the cathode and the anode. A potential minimum just at the anode, i.e.
can only occur if the anode potential is lower than the cathode potential.

To get the radial dependence of the potential between cathode and anode (concentric, infinitely long cylinders) one has to solve the Poisson equation

\[
\frac{d\Phi(r)}{dr} \bigg|_{r_b=r_a} = 0,
\]

with the following boundary conditions:

\[
\Phi(r_c) = V_c = 0 \\
\Phi(r_a) = V_a
\]

We assume a vapour pressure sufficiently low that collision processes may be neglected.

The working conditions of the diode are characterized by the following parameters:

- cathode temperature \( T_c \)
- cathode work function \( W_0 \)
- cathode emission-factor \( a \)

The emission factor \( a \) is introduced to allow for possible imperfect emission conditions at the cathode. The charge density \( \rho(r) \) is calculated, assuming these working conditions at the cathode surface, in the next three paragraphs.
II.2.1 Velocity distribution and current density outside a hot metallic surface

To calculate the electron density outside a metal the surface potential barrier (work-function $W_o$) that the electrons have to pass upon crossing the surface, has to be considered; since the ratio $W_o/k_B$ is of the order of $10^4$ °K, the Fermi distribution can be well approximated by the Maxwellian velocity distribution for electrons outside the cathode; i.e. the Richardson law for the saturation current-density at the cathode surface holds. Fig. 2.2 shows qualitatively what we already know about the potential $\phi(r)$ from within the cathode up to the anode (Fermi potential $\phi_F=\Phi_F/e$).

Fig. 2.2: Sketch of the expected potential distribution between cathode and anode.
Defining $B = 1/(k_B T_C)$, where $T_C$ is the cathode temperature, the charge density per velocity-intervall $d\nu_0$ can be written as:

$$\frac{d\rho_0}{d\nu_0} = \rho_0 \cdot \frac{du_0}{d\nu_0}$$  \hspace{1cm} (2.2.5)

where $\rho_0 = q \left[ \frac{2\pi m}{\beta h^2} \right]^{3/2} \exp\{-B\nu_0^2\}$

and the normalized velocity distribution:

$$\frac{du_0}{d\nu_0} = 2 \left[ \frac{Bm}{2\pi} \right]^{3/2} \exp\{-Bm\nu_0^2/2\}$$

From this we get for the mean velocity and the mean square radial velocity:

$$<v_0> = \sqrt{2/(\pi Bm)}$$  \hspace{1cm} (2.2.6)

$$<v_0^2> = 1/(Bm)$$

Calculating the outward current density by integration, one gets the Richardson law:

$$j_0 = \int d\nu_0 \left\{ \frac{d\rho_0}{d\nu_0} \nu T_C \right\}$$

$$= \rho_0 <v_0> = A_0 T_C^2 \exp\{-BW_0\}$$

with $A_0 = 4\pi q m k_B^2/h^3 = -120$ A/cm$^2$K$^2$ \hspace{1cm} (q = -|e| !).

As we mentioned before, a cathode emission factor $a$ has to be introduced:

$$j_0 = a A_0 T_C^2 \exp\{-BW_0\}$$  \hspace{1cm} (2.2.7)
With all these definitions equation (2.2.5) can be rewritten as

\[
\frac{d\rho}{dv_0} = j_0 \cdot 2 \left( \frac{1}{2\pi \langle v_0^2 \rangle} \right)^{3/2} \exp \left\{ -\frac{v_0^2}{2\langle v_0^2 \rangle} \right\}
\]

This equation specifies the dynamic boundary conditions at the cathode surface.

### II.2.2 Space charge and potential minimum around a cylindrical cathode

If the static boundary conditions allow a potential minimum \( \Phi(r_b) \) at some radius \( r_b \) then the "slow" electrons are reflected back to the cathode somewhere between \( r_c \) and \( r_b \). Thus we can divide the diode into an "inner" and an "outer" region, as shown schematically in Fig. 2.3:

![Schematic sketch of the inner and outer region defined by the space charge.](image)

**Fig. 2.3**: Schematic sketch of the inner and outer region defined by the space charge.
i) inner region with \( r_c \leq r \leq r_b \)

ii) outer region with \( r_b < r \leq r_a \)

By assuming no collisions for the electrons, i.e. their mean free path to be greater than the anode radius, (o.k. for Cs pressures below 1 mTorr) one has the equation of motion:

\[
\ddot{r} = \frac{qE}{m} = -q \nabla \phi(r)
\]

with the constants of motion:

- z momentum \( m_z = m_z \)
- angular momentum \( L_z = m r_c^2 \phi_c = m r^2 \phi \)
- and total energy

\[
\frac{m}{2} \left( \dot{r}^2 + r_c^2 \phi_c^2 + z_c^2 \right) = \frac{m}{2} \left( \dot{r}^2 + r^2 \phi^2 + z^2 \right) + q\phi(r)
\]

The stationary charge distribution \( \rho(r) \) must obey current conservation:

\[
\rho(r) = \text{div} \ j(r) = 0
\]

First, we consider the anode current per unit length:

\[
\frac{i}{\Delta z} = 2\pi r_c \ j(r_c) = 2\pi r \ j(r) = 2\pi r_b \ j(r_b)
\]

If there is a potential minimum \( \phi(r_b) \) outside the cathode, only electrons with a kinetic energy \( E(r_b) \geq 0 \) contribute to the anode current. Energy conservation yields for the inner region:

\[
(2.2.9) \quad E(r_b) \geq 0 \iff E_r + E_{\phi_c} (1-r_c^2/r_b^2) + q\phi(r_c) - q\phi(r_b) \geq 0
\]

and for the outer region:

\[
(2.2.10) \quad E(r_a) \geq 0 \iff E_r + E_{\phi_b} (1-r_b^2/r_a^2) + q\phi(r_b) - q\phi(r_a) \geq 0
\]
The condition (2.2.10) is always fulfilled and simply means that there are no reflected electrons from the outer region.

From condition (2.2.9) we get the so called "critical (radial) velocity" for electrons at the cathode to reach the anode:

\[ v^*(r_c) \geq \sqrt{\frac{2mq^*(r_b)}{v_{\phi c}^2} \left[ 1 - \frac{r_c^2}{r_b^2} \right]} \geq 0 \]  

(2.2.11)

From (2.2.8) and (2.2.11) the anode current density \( j(r_c) \) at the cathode is calculated by integration:

\[ j(r_c) = \frac{j_0}{v_{\phi c}} \int_{v^*(r_c)}^{\infty} \left\{ \frac{du}{dv_\phi} \right\} \left\{ \frac{dv_\phi}{v_C} \right\} \]

(2.2.12)

\[ j(r_c) = j_0 \text{erfc} \left( \frac{mq^*(r_b)}{1 - 1/x_b^2} \right) + \]

\[ + x_b \exp(-mq^*(r_b)) \text{erf} \left( \frac{m}{x_b} \sqrt{1 - 1/x_b^2} \right) \]

where \( \text{erfc}(...) \), the complementary error function, and \( x_b = r_b/r_c \) was used.

Next, the forward current \( i^+(r) \) between the cathode and the potential minimum has to be calculated. Starting with the dynamic boundary condition (2.2.8) at the cathode surface

\[ i^+(r_c) = 2\pi r_c Az j_0 \]

\[ = 2\pi r_c Az \left\{ dv_\phi \left\{ \rho^+(r_c) \frac{du}{dv_\phi} v_C \right\} \right\} \]
and reexpressing this integral for some radius \( r \), where \( r_c \leq r \leq r_b \), we get with the Jacobi determinant and from energy conservation:

\[
(2.2.13) \quad i^+(r) = 2n_r \Delta \varphi \int_0^{v_r} \left[ \rho^+(r) \frac{du}{dv}(r) v_r \right] dv
\]

where

\[
\rho^+(r) = \frac{j_0}{\langle v_o \rangle} \exp(-Bq\Phi(r))
\]

and

\[
\frac{du}{dv}(r) = 2\left(\frac{1}{2\pi\langle v_o^2 \rangle}\right)^{3/2} \exp\left\{ -\frac{v^2}{2\langle v_o^2 \rangle} \right\}
\]

By integrating (2.2.13) from the critical velocity \( v^*(r) \) at some radius \( r \) to \( \varphi \), we find the anode current-density

\[
(2.2.14) \quad j(r) = j_0 \left[ \exp(-Bq\Phi(r)) \operatorname{erfc} \left( \frac{Bq(\Phi(r_b) - \Phi(r))}{1 - x^2/x_b^2} \right) \right] + \frac{x_b}{x} \exp(-Bq\Phi(r_b)) \operatorname{erf} \left( \frac{x}{x_b} \sqrt{\frac{Bq(\Phi(r_b) - \Phi(r))}{1 - x^2/x_b^2}} \right)
\]

For \( r \to r_c \) we retrieve formula (2.2.12). To calculate the anode current we use (2.2.14) in the limit \( r \to r_b \) and find

\[
(2.2.14') \quad i_a = 2n_r \Delta \varphi j_0 \exp(-Bq\Phi(r_b)).
\]

In the case of \( x_b = x_a \), i.e. potential minimum at the anode, we get an exponential behaviour for this part of the current-voltage characteristic:

\[
(2.2.14'') \quad i_a = 2n_r \Delta \varphi j_0 \exp(-Bq\Phi(r_a))
\]

To calculate the stationary charge distribution \( \rho_s(r) \) in the inner region the two currents contributing to this charge distribution have to be considered, viz.
the forward current \( i^+(r) \)

the backward current \( i^+(r) - i_a(r) \)

\[
(2.2.15) \quad \rho_1(r) = 2 \int_0^\infty dv \int dv^* \left\{ \rho^+(r) \frac{d\nu(r)}{dv} - \rho^+(r) \frac{d\nu(r)}{dv^*} \right\}
\]

The first integral in (2.2.15) is easily calculated and can be interpreted as describing the case where all emitted electrons are reflected back to the cathode, that is when \( \Phi(r_b = r_a) \rightarrow \infty \) (i.e. \( v^*(r) \rightarrow \infty \)).

\[
(2.2.15') \quad \rho_1(r) = 2 \rho^+(r) - \int_0^\infty dv \left\{ \rho^+(r) \frac{d\nu(r)}{dv} \right\}
\]

For the second integral due to a non-vanishing anode current we get:

\[
\rho^+(r) \left[ 1 - \frac{2}{\sqrt{\pi}} \int_0^{y^*} dy \left\{ e^{-y^2} \text{erf}(S) \right\} \right]
\]

where \( y^* = \sqrt{\frac{2q(\Phi(r_b) - \Phi(r))}{1-x^2/x_b^2}} \),

and \( S = \sqrt{\frac{2q(\Phi(r_b) - \Phi(r)) - y^2(1-x^2/x_b^2)}{1-x^2/x_b^2}} \)

and finally for the stationary charge distribution in region 1:

\[
(2.2.16) \quad \rho_1(r) = \rho^+(r) \left[ 1 + \frac{2}{\sqrt{\pi}} \int_0^{y^*} dy \left\{ e^{-y^2} \text{erf}(S) \right\} \right]
\]

For \( r = r_b \) the above integral is zero, and one gets the charge density at the potential minimum: \( \rho_1(r_b) = \rho^+(r_b) \)
(2.2.17) \( \rho_1(r_b) = \frac{j_s}{<v_s>} \cdot \exp(-\delta q(r_b)) \)

This is reasonable since at \( r_b \) no reflected electrons contribute to the charge density. From equation (2.2.13) we further get the dynamic boundary conditions for the outer region, i.e. at the "virtual cathode" of radius \( r_b \).

\[
(2.2.18) \quad \frac{d\rho}{dv}(r_b) = \rho_1(r_b) \frac{du}{dv}(r_b)
\]

II.2.3 Charge density outside the potential minimum

Starting from the "dynamic" boundary conditions at the virtual cathode (2.2.18), we can write with the current conservation (only forward current):

\[
i_a(r_b) = i^+(r_b) = 2\pi r_b \Delta z \int_0^\infty dv \left[ \rho_1(r_b) \frac{du}{dv}(r_b) v_r \right]
\]

\[
= i_a(r) = 2\pi r \Delta z \int_0^\infty dv \left[ \rho(r) \frac{du}{dv}(r) v_r \right] v_{min}(r)
\]

To get the integrand in \( i_a(r_b) \) in function of \( r \) explicitly, we reexpress it at some radius \( r \) with \( r_b \leq r \leq r_a \):

\[
i_a(r) = 2\pi r \Delta z \rho_1(r_b) \exp(-\delta q(r) + \delta q(r_b)) \int_0^\infty dv \left[ \frac{du}{dv}(r) v_r \right] v_{min}(r)
\]
The lower bound for the integration we get by setting $v_{rb} = 0$ in the equation from the energy conservation:

\begin{equation}
(2.2.19) \quad v_{\text{min}}(r) = \sqrt{\frac{v^2}{r_b^2} - 1} - \frac{2}{m} \left[ q'(r) - q'(r_b) \right]
\end{equation}

For $r \rightarrow r_b$ this integration gives the same current density as (2.2.14). Thus we are allowed to write for the charge density in the outer region:

\begin{equation}
(2.2.20) \quad \rho_2(r) = \rho_1(r_b) \exp\left\{-6q^\#(r) + 6q^\#(r_b)\right\} \times \frac{I}{\rho_1(r_b)}
\end{equation}

where \( I = \int dv \int dv_r \int dv_z \left( \frac{du}{dv} \right) \)

For the integral \( I \) we get

\[
I = 1 - \frac{2}{\sqrt{\pi}} \int dy \left\{ e^{-y^2} \text{erf}\left(\frac{u_{\text{min}}(y^2)}{\sqrt{2}v_B^2}\right) \right\}
\]

where \( u_{\text{min}}(y^2) = v_{\text{min}}(r)/\sqrt{2v_B^2} \)
and \( y = v/\sqrt{2v_B^2} \)

For high anode potentials the exponential factor in (2.2.20) can cause floating point overflow in numerical calculations. This problem is eliminated by working with the transformed integral:

\begin{equation}
(2.2.21) \quad \rho_2(r) = \rho_1(r_b) \frac{2}{\sqrt{\pi}} \int dy \left\{ \frac{e^{-y^2} \text{erf}\left(y/\sqrt{x^2/x_B^2 - 1}\right)}{1 + \frac{1}{y^2} \left( 6q^\#(r_b) - 6q^\#(r) \right)} \right\}
\end{equation}
With the charge densities (2.2.16) and (2.2.21) we finally give the differential equation to be solved:

\[
\frac{d^2 \Psi}{dr^2} + \frac{1}{r} \frac{d \Psi}{dr} = -4\pi \rho(r)
\]  
(2.2.22)

with \( \rho(r) \) defined as:

\[
\rho(r) = \begin{cases} 
  \rho_1(r) & \text{if } r_c \leq r \leq r_b \\
  \rho_2(r) & \text{if } r_b \leq r \leq r_a 
\end{cases}
\]

In region 1 we have to solve a boundary-value problem:

i) \( \Psi(r_c) = 0 \)

ii) \( \Psi'(r=r_b) = 0 \)

with the additional condition for the radius of the potential minimum:

iii) \( r_c \leq r_b \leq r_a \)

The solution will give us a value for the potential minimum \( \Psi(r_b) \leq 0 \) and from this we get the anode current by equation (2.2.14')

\[
i_a = 2\pi r_b \Delta z j_0 \exp(-\beta q \Psi(r_b))
\]

In region 2 an initial value problem has to be solved:

i) \( \Psi(r_b) \) from region 1

ii) \( \Psi'(r=r_b) = 0 \)

From this solution we get the anode-cathode potential:

\[
\Psi(r_a) = V_a
\]

Solving the differential equation (2.2.22) for different radii \( r_b \) of the potential minimum produces the current-voltage characteristic of the thermionic diode (naturally, for a specified \( T_c \)).
II.3 Numerical solution of the Poisson equation

For the numerical solution we transform (2.2.22) into a dimensionless equation by choosing

\[ x = r/r_c \quad \text{and} \quad W(x) = -\beta q \phi(x) \]

\[ W''(x) + \frac{1}{x} W'(x) = 4\pi\beta q r_c \rho_i(x) \quad \text{with } i = 1,2 \quad \text{for the inner and outer regions respectively,} \]

and the boundary values

\[ W_i(1) = 0 \]
\[ W_i'(x=x_b) = 0 \]

We further transform Eq. (2.3.1) to be able to use subsequently the method of integration by finite differences [Har 57]. We define

\[ Y(x) = \int x W(x) \]

and obtain

\[ Y''(x) = -\frac{Y_i(x)}{4x^2} + 4\pi\beta q r_c \int x \rho_i(x,Y_i) \quad \text{with the boundary values} \]

\[ Y_i(1) = 0 \]
\[ Y_i'(x=x_b) = Y_i(x_b)/(2x_b) \]

For the calculation, the following five parameters have to be specified:
cathode radius \( r_C \)
anode radius \( r_a \)
cathode temperature \( T_C \)
cathode work function \( W_0 \)
cathode emission factor \( \alpha \)

The emission factor \( \alpha \) was already included into the saturation current-density \( j_0 \).

Now the dimensionless differential equation to be solved reads:

\[
\begin{equation}
\label{2.3.3}
Y''_i(x) = -\frac{Y_i(x)}{4x^2} + ab \frac{r^2}{T_C} \exp(-2W_0) \frac{1}{x} \tilde{\rho}_i(x,Y_i)
\end{equation}
\]

where \( b = \alpha \frac{mc^2}{(\hbar c)^2} \sqrt{\frac{2}{\pi}} mc^2 k_B \)

and \( \tilde{\rho}_i = \frac{<v_0>}{j_0} \rho_i(x,Y_i) \)

The code to solve \((2.3.3)\) is written in Turbo Pascal and runs on a Personal Computer (Olivetti M24). The calculations presented so far are summarized in the flow chart below.
Electron velocity distribution at cathode surface (Richardson law)

cylindrical geometry ϑ(rb) divides diode into two distinct regions

equations of motion current conservation (steady state)

forward and backward current in "inner" region
\[ r_c \leq r \leq r_b \]
\[ \rho_1(r) \]

"virtual" cathode at radius rb
\[ \varphi_1(r_b) = \varphi_2(r_b) \]
\[ \varphi_1'(r_b) = \varphi_2'(r_b) = 0 \]
\[ \rho_1(r_b) = \rho_2(r_b) \]

only forward current in "outer" region
\[ r_b \leq r \leq r_a \]
\[ \rho_2(r) \]

Poisson equation
\[ Y_1''(x) = f(x, Y_1) \]

solve differential equation numerically

\[ x = r/r_c \]
\[ i = 1, 2 \ (\text{inner, outer reg.}) \]
\[ Y_1(x) = -\sqrt{x} \ \text{sgn} \varphi_1(x) \]
II.4 Comparison of calculated and measured current-voltage characteristics

The input parameters for the calculation are determined as follows: The cathode temperature is determined from the measured temperature vs filament current characteristic (Appendix D). The cathode work function for cesiated tungsten is taken from [Cog 68] as the mean of the values for the 110 and 112 surfaces of tungsten crystals, $W_0 \approx 1.70$ eV (this value agrees below 800 K with the value given by [Tay 33] for polycrystalline tungsten). The imperfect cylindrical geometry of our diodes is allowed for by taking a mean radius for the calculation (Appendix B).

Fig. 2.4 shows a measured and the calculated current vs voltage characteristic of our thermionic diodes, i.e. a very good agreement of theory with experiment. The dashed line represents Langmuir's 3/2 power law (2.1.1). The calculated radial potential distribution at typical working conditions ($V_a \approx 0.4$ V; cathode heating current $\approx 0.75$ A) is shown in Fig. 2.5. Fig. 2.6 shows potential distribution, electric field and electron charge density for the same conditions. Fig. 2.7 shows the zoomed electron density at the potential minimum. The chosen "mesh width" for the numerical integration in the outer region is indicated there as well.

A comparison of the measurements with the theory presented must be confined to the case where $x_b > 1$ (i.e. not vanishing space charge outside the cathode), since the assumed potential step $W_0/e$ at the cathode surface (Fig. 2.2) leads to a discontinuity in the electric field for $x_b > 1$.

The theory could be completed for the latter case, i.e. for a description of the saturation behaviour of the anode current, by including the image charge potential (Schottky effect [Sch 14]) near the cathode surface.
$R_c = 0.05 \text{ cm}$
$R_0 = 1.60 \text{ cm}$
$T_c = 825 \text{ K}$
$W_0 = 1.70 \text{ eV}$
$a = 0.6$

Fig. 2.4: Calculated and measured current-voltage characteristic at a Cs pressure of $\approx 4 \text{ mTorr}$. The dashed line represents Langmuir's $3/2$ power law.

Fig. 2.5: Calculated potential distribution at the same conditions as Fig. 2.4 for an anode-cathode voltage of $\approx 0.4 \text{ V}$.
Fig.2.6a) : Calculated potential distribution  
(same as Fig. 2.5; note log-scaled R axis).

Fig.2.6b) : Calculated electric field (same conditions as  
Fig. 2.5; note log-scaled R axis).
Fig. 2.6c: Calculated electron density (same conditions as Fig. 2.5; note log-scaled R axis).

Fig. 2.7: Blow-up of charge density near minimum (Fig. 2.6c). Dashed lines indicate mesh-size used in numerical integration.
II.5  Estimation of possible gain of the thermionic diode

From the results of the calculations presented in Chapter II.4 we can give an estimation of maximum possible gain. By assuming

i) only few positive ions are in the diode, i.e. no distortion of the potential distribution: 
   \( r_b = \) constant. I.Popescu and L.Brandus have shown  
   (for thermionic converters with plane electrodes) 
   that no appreciable reduction of the space charge 
   is to be expected for ion density to electron density ratios < 1/10 [Bra 64].

ii) the ions are trapped near the radius \( r_b \) after thermalization in the parent vapour,

iii) the ions disappear (neutralize) after a mean trapping time \( \tau \),

we can write for the ion current through the diode:

\[
(2.5.1) \quad i_i = 2\pi r_b \Delta z \frac{\Delta r_b}{\tau} \rho_{ion}(r_b)
\]

There, \( \Delta r_b \) is the thickness of the "ion trapping volume" of the diode and \( \Delta r_b /\tau \) can be interpreted as a mean ion velocity through the hollow cylinder shaped trapping volume. The value of \( \Delta r_b \) can be estimated by comparing the potential energy of the ions and the mean kinetic energy of the vapour atoms \( \frac{3}{2} k_B T_v \), as indicated in Fig. 2.8 (note log scaled r axis†). The anode current can be written as

\[
(2.5.2) \quad i_a = 2\pi r_b \Delta z \langle v_o \rangle \rho(r_b)
\]

by using equations (2.2.14'') and (2.2.17).

In Chapter I.3 the steady state diode gain was defined as

\[
G_s = \frac{\Delta i_a}{\Delta i_e}
\]
Fig. 2.8: Calculated potential distribution (Fig. 2.6a) with the mean kinetic energy of the vapour atoms (ions) indicated.

For a crude estimate the charge density \( \rho(r_b) \) is assumed to be constant within \( \Delta r_b \). Then the space charge reduction can formally be written as

\[
\Delta \rho(r_b) = \Delta \rho_{\text{ion}}(r_b)
\]

From this we get an expression for the maximum possible gain of the thermionic diode, a ratio of two mean velocities:

\[
G_a \leq \langle v_o \rangle \cdot \frac{\tau}{\Delta r_b}
\]

For our diode at the same typical working conditions as in II.4 we find from the measured cathode temperature \( T_c = 825 \) K and trapping time \( \tau = 23 \) sec (Fig. 1.6), and the calculated
thickness of the ion trapping volume \( \Delta r_b = 0.228 \) cm (Fig. 2.8):

\[
G_e < 9 \cdot 10^8
\]

This crude estimate can be refined by considering the following facts:

i) The space charge reduction is most effective inside the potential minimum, i.e. the "active volume" is about 1/3 of the ion trapping volume (see Fig. 2.8). Hence, \( G_e \) has to be reduced by a factor of 3 since the ions are "active" only 1/3 of their trapping time.

ii) The electron charge density in the active volume varies by a factor of 3 (see Fig. 2.6c), i.e. is not constant as assumed for the crude estimate. From this the effective space charge reduction is assumed to be about 1.5 times lower.

With these two factors we get a better estimation of the diode gain:

\[
G_e \leq 2 \cdot 10^8
\]

The gain from this rough estimation has to be compared with the measured \( G_e \) from Chapter I.3 by considering the ten times lower ion trapping time there:

<table>
<thead>
<tr>
<th>estimation from calculation</th>
<th>measured filament on - off</th>
</tr>
</thead>
<tbody>
<tr>
<td>( G_e ) \leq 2 \cdot 10^8</td>
<td>( = 1.2(0.2) \cdot 10^8 )</td>
</tr>
</tbody>
</table>

This agreement within a factor of two is leading us to the conclusion that the behaviour of the thermionic diode now is understood quantitatively.
II.6 Thermionic diode in an homogenous magnetic field

The measured behaviour of the diode in axial magnetic fields (see paragraph I.5) is analyzed here. The equation of motion of electrons in a magnetic field is

\[ m_{\text{a}} \overset{\mathbf{\ddot{a}}}{=} \overset{\mathbf{\dot{a}}}{fa} + \frac{q}{c} \overset{\mathbf{\dot{v}}}{v_{\text{a}}} \wedge \overset{\mathbf{\mathbf{B}}}{B}, \]

where we have chosen the index a for the absolute frame (inertial frame). By describing the motion in a rotating frame and choosing the rotation frequency as

\[ \overset{\mathbf{\dot{\omega}}}{\omega} = -\overset{\mathbf{\dot{\omega}}}{\omega_{L}} = -\frac{q}{2mc} \overset{\mathbf{\mathbf{B}}}{B}, \]

\( \omega_{L} \) = Larmor frequency, the Lorentz force is compensated by the Coriolis force and only the centrifugal force is left:

\[ m_{\text{a}} \overset{\mathbf{\ddot{a}}}{=} \overset{\mathbf{\mathbf{F}}_{\text{a}}}{F_{\text{a}}} + m\overset{\mathbf{\dot{\omega}}}{\omega_{L}} \wedge (\overset{\mathbf{\dot{\omega}}}{\omega_{L}} \wedge \overset{\mathbf{v}}{v}) \]

For cylindrical geometry with an axial magnetic field we get the equations of motion in the rotating frame:

\[ m(r - r\phi) = -q \frac{d}{dr} \Phi(r;B) - mr\omega_{L}^{2} \]

\[ m \frac{1}{r} \overset{\mathbf{\ddot{z}}}{z} = 0 \]

\[ mz = 0 \]

i.e. \( p_{z} \) = const., \( p_{\phi} = mr^{2} \phi = \text{const.} \), and the energy conservation \( E(r_{0};B) = E(r;B) \) reads

\[ E(r;B) = \frac{m}{2} \overset{\mathbf{\mathbf{z}}}{z}^{2} + \frac{p_{\phi}^{2}}{2m} \frac{1}{r^{2}} + \frac{p_{z}^{2}}{2m} + q\Phi(r;B) + \frac{m}{2} \overset{\mathbf{\mathbf{z}}}{z}^{2} \omega_{L}^{2} \]
Further, current conservation must be valid also for the stationary charge distribution $\rho(r;B)$:

$$\rho(r;B) = \text{div} \ j(r;B) = 0$$

From this and the calculations in II.2.2 and II.2.3 we see that the charge density $\rho(r;B)$ and the current density $j(r;B)$ in the rotating frame may be calculated following the same arguments as without magnetic field by treating the term $mr^2 \omega_L^2/2$ in (2.6.5) as an additional potential (centrifugal potential). We define an effective potential by setting

$$q^\ast_{\text{eff}}(r;B) = q^\ast(r;B) + mr^2 \omega_L^2/2.$$  

Because of the axial symmetry the radial current and charge densities in the rotating frame and in the absolute frame are identical.

In the case of $\Phi_{\text{eff}}$ having the minimum at the anode, i.e. for large negative anode potentials and (or) large magnetic fields ($q^\ast_{\text{eff}}(r_a;B) > 10 /\beta$), we get for the anode current:

$$i_a(B) = 2\pi r_c \Delta z \ j(r_c;B,r_a)$$

with the current density at the cathode ($\Phi(r_c;B) = 0$):

$$j(r_c;B,r_a) = j_o \left[ \exp(-\beta q^\ast_{\text{eff}}(r_c;B)) \ \text{erfc}(U_{ac}) ight.$$

$$\left. + x_a \exp(-\beta q^\ast_{\text{eff}}(r_a;B)) \ \text{erf}(U_{ac}/x_a) \right]$$

and $U_{ac} = \frac{\beta q^\ast_{\text{eff}}(r_a;B) - \Phi_{\text{eff}}(r_c;B)}{1 - 1/x_a^2}$

If however the minimum of $\Phi_{\text{eff}}(r;B)$ is at some radius $r_b$ with $r_c \leq r_b \leq r_a$, then the current density at the cathode reads:
The charge density from the forward current at some radius \( r \) with \( r_c \leq r \leq r_b \) can again be calculated from the forward current and yields:

\[
(2.6.10) \quad \rho^+(r;B) = \frac{j_o}{<v_o>} \exp(-\beta q \phi_{\text{eff}}(r;B))
\]

And for the current density we get:

\[
(2.6.11) \quad j(r;B) = j_o \left[ \exp(-\beta q \phi_{\text{eff}}(r;B)) \text{erfc}(U_{br}) + \frac{x_b}{x} \exp(-\beta q \phi_{\text{eff}}(r_b;B)) \text{erf} \left( \frac{x}{x_b} U_{br} \right) \right]
\]

and \( U_{br} = \sqrt{\frac{8q(\phi_{\text{eff}}(r_b;B) - \phi_{\text{eff}}(r;B))}{1 - x^2/x_b^2}} \)

The anode current is again calculated by using (2.6.11) in the limiting case \( r=r_b \):

\[
(2.6.11') \quad i_a(B) = 2\pi r_b A z j_o \exp(-\beta q \phi_{\text{eff}}(r_b;B))
\]

If \( r_b = r_a \), i.e. the minimum of the effective potential is at the anode, the anode current is determined only by geometry, the anode potential and the magnetic field. In that case, there is no need to solve the Poisson equation to get the current-voltage or current-B-field characteristics:
(2.6.11) \[ i_a(B) = 2\pi r_a \Delta z j / \exp(-8q/(r_a;B) - \delta r_a^2/\omega_a^2) \]

For \( r_c < r < r_b < r_a \) the stationary charge distribution is found to be

\[
(2.6.12) \quad \rho_i(r;B) = \rho^+(r;B) \left[ 1 + \frac{2}{\sqrt{\pi}} \int_0^{y^*} dy \left( e^{-y^2} \text{erf}(S) \right) \right]
\]

with \( y^* = \sqrt{\frac{8q(\delta_{\text{eff}}(r_b;B) - \delta_{\text{eff}}(r;B))}{1-x^2/x_b^2}} \)

and \( S = \sqrt{8q(\delta_{\text{eff}}(r_b;B) - \delta_{\text{eff}}(r;B)) - y^2(1-x^2/x_b^2)} \)

For \( r > r_b \) equation (2.6.12) again leads to the dynamic boundary conditions at the virtual cathode (\( \rho_i(r_b;B) = \rho^+(r_b;B) \)):

\[
(2.6.13) \quad \frac{d\rho_i(r_b;B)}{dv} = \rho_i(r_b;B) \frac{d\mu(r_b;B)}{dv} \]

To calculate the charge density outside the potential minimum we follow the steps in section II.2.3 and find

\[
(2.6.14) \quad \rho_z(r;B) = \rho_i(r_b;B) \frac{2}{\sqrt{\pi}} \int_0^\infty Y(y) \, dy
\]

\[
Y(y) = \frac{e^{-y^2} \text{erf} \left[ y/\sqrt{x^2/x_b^2 - 1} \right]}{\left[ 1 + \frac{8q}{y^2} \delta_{\text{eff}}(r_b;B) - \delta_{\text{eff}}(r;B) \right]}
\]
II.7 Comparison of calculated and measured behaviour of the thermionic diode in axial magnetic fields

First the diode is set to conditions where equation (2.6.11") holds, i.e. low anode-cathode bias voltage to keep the potential minimum at the anode, even without magnetic field. Fig. 2.9 shows the measured current-B-field characteristic.

From equation (2.6.11") we get for the turning-point and the half width at half maximum value of the characteristic:

\[
\omega_L^{2}(t_p) = \frac{k_B T_c}{mr_a^2}
\]

\[
\omega_L^{2}(hwhm) = 2 \ln 2 \omega_L^{2}(t_p)
\]

and from this

\[
(2.7.1) \quad B_z(t_p) = b \sqrt{\frac{T_c}{r_a}}
\]

where

\[
b = \frac{\hbar c}{\mu_B} \sqrt{\frac{k_B}{mc^2}}
\]

\[
= 4.427 \times 10^{-2} \text{ gauss cm/}\sqrt{\text{K}}
\]

Since our diodes don't have a perfect cylindrical geometry we have to use the minimum anode radius, i.e. the anode radius seen by the electrons in the rotating frame (rotation frequency = 1.4 MHz/gauss). The cathode temperature \( T_c = 690(20) \text{ K} \) is read from the calibration characteristic (Appendix D), and \( r_a(\text{min}) = 0.85 \text{ cm} \) (see Appendix B).

\[
\begin{array}{|c|c|c|}
\hline
& \text{Theory} & \text{Experiment} \\
\hline
B_z(t_p) & 1,368(20) \text{ gauss} & 1,2(1) \text{ gauss} \\
\hline
B_z(hwhm) & 1,611(22) \text{ gauss} & 1,7(1) \text{ gauss} \\
\hline
\end{array}
\]
b) Measured current vs. B-field characteristic at:
- Cs pressure ≈ 0.05 mTorr
- cathode temperature = 690(20) K
- anode-cathode voltage = -0.5 V

a) $\frac{dI_a}{dB}(B)$ corresponding to b).
i.e. a rather good agreement. To test the theory accurately, one should work with a perfectly cylindrical diode, one could then decide whether this is a good wireless method for non-optical measurement of surface temperatures of electron emitting materials.
CHAPTER III

Polarization correlations in the field-free two-photon transition $6S_{1/2} \rightarrow 7S_{1/2} \rightarrow 15P_{3/2}$ in Cesium

The prime purpose of the polarization-dependent (but parity-conserving) asymmetry measurements presented here is

i) to investigate those systematic effects arising in asymmetry measurements with our detection scheme that could mimic parity violation,

ii) to optimize the operating conditions of the experimental setup,

iii) to get information on the signal/noise ratio for the measuring time of the parity experiment proper and

iv) to decide whether or not collisional population redistribution among the $15P_{3/2}$ sublevels (quasi elastic scattering) has to be taken into account.

Having shown that the residual electric [Wei 84] and magnetic fields [Che 87], [Che 89], [Chapter I, this paper] can be determined in situ and controlled at the required low level, the uncertainties in light intensity and state of polarization within the diode volume, due to the imperfections of its windows, still required detailed investigation.
III.1 Calculation of parity conserving polarization dependent asymmetries in the $6S_{1/2} \rightarrow 7S_{1/2} \rightarrow 15P_{3/2}$ two photon transition

The model of the processes occurring in our two photon experiments presented in this paragraph was calculated by J. Hoffnagle [Hof 85]. It permits quantitative predictions of the rates of Rydberg atoms produced in the thermionic diode. The levels and transitions involved are shown schematically in Fig. 3.1. The calculations are based on the following simplifying assumptions:

i) The coherence between the atomic states coupled by the electromagnetic transitions is neglected, i.e. "rate equations" are used to describe the evolution of the populations of the levels.

ii) Collision effects in the $6S$ and $7S$ states are neglected since the collisional depopulation cross-section is proportional to $(n^*)^4$ [Dee 77] and the depolarization cross-section [Hof 81;2] contributes only about $1/10$ at a pressure of 10 mTorr to the radiative lifetime.

iii) The experimentally unresolved h.f. splitting of the $P_{3/2}$ state is neglected (see later).

iv) Collisional population redistribution among the sublevels of the $P_{3/2}$ state by scattering processes (quasi-elastic scattering) is included. The rates are assumed to be independent of the h.f. and magnetic sublevels and their magnitude is estimated from experimental broadening data.
Fig. 3.1: Scheme of the Cs terms and processes involved in the resonant two-photon transition $6S_{1/2} \rightarrow 7S_{1/2} \rightarrow nP_{3/2}$ (term energies are not to scale), and definition of the symbols used in the calculations:

- $A_\alpha$, $B_\beta$, $C_\gamma$ populations of the h.f. Zeeman sublevels involved.
- $R^{(s)}_{\beta\alpha}$ the M1 transition rates.
- $R^{(o)}_{\gamma\beta}$ the E1 (saturated) transition rates.
- $T_{\gamma\delta}$ rates of collisional population redistribution among the sublevels of the $nP_{3/2}$ state (quasi-elastic scattering).
- $\lambda_1$, $q$; $\lambda_2$, $q'$ the wavelengths and polarization states of the two laser beams ($q$, $q' = \pm 1, 0$ for circular and linear polarizations respectively).
To estimate the linewidth of the two photon transitions we use the same definitions as [Shi 76]:

- Longitudinal relaxation rates: \( \tau_i = \tau_i^{-1} + \tau_i^{\text{coll}} \)
  where \( \tau_i \) is the radiative lifetime and \( \tau_i^{\text{coll}} \) the collisional depopulation rate of level \( i \), i.e. \( \tau_i \)'s are fwhm's

- Longitudinal relaxation time:
  \[ 2 \tau = \tau_B^{-1} + \tau_C^{-1} \]

- Transverse relaxation rate:
  \[ 2 \gamma = \tau_B + \tau_C + \Delta \omega^{\text{coll}} \]
  with \( \Delta \omega^{\text{coll}} \) the fwhm broadening of the line due to phase changing collisions.

The radiative 7S lifetime is \( \tau_{7S} \approx 53 \) nsec [Hof 81;2], [Bou 84]; the collisional depopulation rate \( \tau_{7S}^{\text{coll}} \) is negligible.

From the measured lifetimes of Rydberg S and D-states [Dee 77] the lifetime of the 15P state is estimated to be of the order of 1 \( \mu \)sec. According to the same reference, the inelastic depopulation cross section for low-lying Rydberg states is well approximated by

\[
\sigma_i = \pi <r^2> = \pi \left( \frac{5}{2} n^*^4 - \frac{3}{2} n^*^2 l(l+1) + \frac{n^*^2}{2} \right)
\]

for \( n^* \leq 11.5 \) (corresponds to 15P). Using this inelastic cross section, the collisional depopulation rate of level \( i \) is

\[
(3.1.1) \quad \tau_i^{\text{coll}} = \rho \sigma_i(i) \bar{v},
\]

where \( \rho \) is the density of the ground state atoms and \( \bar{v} \) the mean relative velocity of the collision partners:

\[
\bar{v} = 4 \, c \, \sqrt{\frac{kT}{M c^2}}.
\]
From these data we estimate the longitudinal relaxation rates at the Cs pressure of 10 mTorr:

\[ \tau_{7S} \approx 2\pi \cdot 3.0 \text{ MHz} \quad \text{and} \quad \tau_{15P} \approx 2\pi \cdot 9.2 \text{ MHz} . \]

For the longitudinal relaxation time we obtain

\[ \tau (7S \rightarrow 15P) \approx 35 \text{ nsec}. \]

The collisional broadening data for the Rydberg D states are taken from [Hei 83] (measured impact broadening parameters of $6S \rightarrow nS$ and $6S \rightarrow nD$ Doppler-free non-resonant two-photon transitions). Since the energy difference of the $14D$ ($n^* \approx 11.53$) and the $15P$ ($n^* \approx 11.44$) states is only $\approx 14 \text{ cm}^{-1}$ [Lor 84], [Wei 84], we take the $14D$ value (corrected to a 1-photon transition) for the estimate:

\[ \Delta \omega_{\text{coll}}^{\text{tot}} (6S \rightarrow 15P) \approx 2\pi \cdot 15 \text{ MHz} \]

This gives us finally an estimate of the homogenous fwhm of the $6S \rightarrow 15P$ two photon transition:

\[ 2\pi (6S \rightarrow 15P) \approx 2\pi \cdot 18 \text{ MHz} \]

From these data we can also estimate a maximum quasi-elastic scattering rate for the $15P_{3/2}$ state in Cs at 10 mTorr. The total collisional width (3.1.4) of the transition includes contributions from:

- inelastic scattering processes
  \[ \tau_{\text{inel}}, \text{i.e. } \Delta n \neq 0 \text{ or } \Delta l \neq 0 \]
- quasi elastic scattering processes
  \[ \tau_{\text{qe}}, \text{i.e. } \Delta n = 0, \Delta l = 0, \Delta J = 0 \text{ and } \Delta F \neq 0 \text{ and (or) } \Delta M \neq 0 \]
- elastic (phase changing) scattering processes.
A reasonable upper bound for the quasi-elastic scattering rate can be estimated by

\[ \tau_{qe}(F'_i M'_i \to F''_i M''_i) \leq \frac{\tau_{tot} - \tau_I}{g_R} \]

with \( \tau_{tot} \) the total collisional width (3.1.4), \( \tau_I \) the pure inelastic width (3.1.2), \( g_R = 32 \) \( I=7/2, F''=2,3,4,5 \) the degeneracy of the \( P_{7/2} \) Rydberg state and the assumption of all the \( \tau_{qe} \)'s being equal.

(3.1.6) \( \tau_{qe} \leq 2\pi \cdot 0.19 \text{ MHz} \)
III.1.1 Effect of hyperfine structure in the 6S_{1/2} \rightarrow 7S_{1/2} \rightarrow \text{nP}_{3/2} two photon transition without saturation

In our experiment the populations of the completely resolved 7S_{1/2} hyperfine levels are detected via a 7S_{1/2} \rightarrow \text{nP}_{3/2} transition with unresolved hyperfine splitting in the final state. The rate equation approximation implicitly assumes no phase correlations among the 7S h.f. Zeeman sublevels which are "pumped" by the second laser to a Rydberg state. In this approximation the h.f. splitting of the nP state is neglected since the expected transition linewidth is about five times larger than the frequency difference between the F'' = 2 and F'' = 5 components (^{133}\text{Cs} : I=7/2), of the 15P_{3/2} state:

\begin{equation}
\Delta \epsilon (F''=5.2) \approx 5.4 \text{ MHz} \quad A(15P_{3/2}) \approx 0.45 \text{ MHz}
\end{equation}

where the interval factor A was obtained by an extrapolation of the data of Arimondo [Ari 77].

The transition probabilities \( W_q(7,F',M') \) to the \(|7S_{1/2} F'M'>\) states are proportional to \(|<7S_{1/2} F'M'|u||6S_{1/2} FM>|^2\) with the notations defined in Fig. 3.2 (we follow the convention of Sobelman [Sob 79]):

\[
W_q(7,F',M') = |<7S_{1/2}|u||6S_{1/2}>|^2 \cdot (2F+1)(2F'+1) \begin{vmatrix} 1/2 & 7/2 \end{vmatrix}^2 \begin{vmatrix} F & F' & 1 \end{vmatrix}^2 \begin{vmatrix} M & -M' & q \end{vmatrix}^2
\]

where the reduced matrixelement corresponds to the M1 transition \(|M1|^2 = |M+M_{hf}(F-F')|^2\) as defined in [Bou 74] [Hof 82]. After the second transition (E1) the probabilities of populating the \text{nP}_{3/2} h.f. Zeeman sublevels are found to be
\[ W_{q'q}(n,F''',M''',F''',M') = W_q(7,F'',M'') \cdot |<nP_{3/2}|1S_{1/2}>|^2 \cdot \]
\[ \cdot (2F''+1)(2F'''+1) \left\{ \begin{array}{ccc} 3/2 & F'' & 7/2 \\ F' & 1/2 & 1 \end{array} \right\} \left[ \begin{array}{cc} F' & F'' \\ M'-M'' & q' \end{array} \right]^2 \]

We have to sum over the magnetic quantum numbers, and since the h.f. levels of the nP_{3/2} state are not resolved, over F'':

\[ W_{q'q}(n,F'F) = \sum_{M'} \sum_{F''} W_{q'q}(n,F''',M''',F''',M') \]

As long as either q or q' = 0, i.e. at least one of the lasers is linearly polarized, no difference is found among the 6S_{1/2} \rightarrow 7S_{1/2} h.f. components in their transition probabilities to the nP_{3/2} states. On the other hand, if both lasers are circularly polarized one finds a dependence ("analyzing power") for the various h.f. components. In Table III.1, this "analyzing powers" are given and ss and os denote there the same and opposite sense of the rotating electric field vectors respectively (ss if q·q' = +1; os if q·q' = -1). \( A_{cp}(0) \) is the circular polarization asymmetry in the limit of zero saturation (G = 0):

\[ A_{cp} = \frac{ss - os}{ss + os} \]

Table III.1: Circular polarization asymmetries in the two photon transition: 6S_{1/2} \rightarrow 7S_{1/2} \rightarrow nP_{3/2}. (G = 0)

<table>
<thead>
<tr>
<th>F → F'</th>
<th>( W_{q'q}(n,F',F) )</th>
<th>( A_{cp}(G = 0) )</th>
</tr>
</thead>
<tbody>
<tr>
<td>ss</td>
<td>os</td>
<td></td>
</tr>
<tr>
<td>3 → 3</td>
<td>15/16</td>
<td>17/16</td>
</tr>
<tr>
<td>4 → 3</td>
<td>19/16</td>
<td>13/16</td>
</tr>
<tr>
<td>3 → 4</td>
<td>21/16</td>
<td>11/16</td>
</tr>
<tr>
<td>4 → 4</td>
<td>17/16</td>
<td>15/16</td>
</tr>
</tbody>
</table>
III.1.2 Rate equation model with saturation

In a simple 3-level system with an unsaturated transition (A→B) and the atoms at rest, the two-photon transition rate is equal to

$$ (3.1.8) \quad R(A \rightarrow B \rightarrow C) = R(A \rightarrow B) \frac{\gamma_C}{\gamma_B + \gamma_C} \frac{G_0}{1 + G_0}, $$

which exhibits the saturation behaviour in its simplest form. Allowing for level degeneracies, one has to solve a more complex system. The rate equations (with the mentioned simplifying assumptions) used in these calculations are as follows (see Fig. 3.1):

$$ (3.1.9) \quad B_{\gamma} = R_{\beta \alpha}^{(\gamma)} - \gamma_B \sum \sigma_\gamma C_{\gamma} \cdot \frac{1}{\gamma_B + \gamma_C} \left( C_{\gamma} B_{\gamma} + \sum \sigma_\gamma C_{\gamma} (C_{\beta} - B_{\gamma}) \right) $$

$$ C_{\gamma} = \frac{-\gamma_C}{\sum \sigma_\gamma} C_{\gamma} + \sum \sigma_\gamma (B_{\beta} - C_{\gamma}) + \sum \sigma_\gamma (C_{\delta} - C_{\gamma}), $$

where $R_{\beta \alpha}^{(\gamma)}$ is the M1 transition rate. This transition is so weak that the effect on the populations $A_{\alpha}$ (Fig. 3.1) can be neglected. The other used symbols denote:

- $\gamma_B$, $\gamma_C$ the longitudinal relaxation rates $\{\gamma = \tau^{-1} + \tau^{coll.}\}$
- $\Phi_2$ the flux of photons with wavelength $\lambda_2$
- $\sigma_\beta$ the $B_{\beta} \rightarrow C_{\gamma}$ absorption cross section and
- $\sigma_\delta$ the quasi-elastic scattering transfer rates as estimated before.

The saturation parameters for the single Zeeman components are defined as:

$$ G_\delta = \Phi_2 \sigma_\beta \left( \gamma_B^{-1} + \gamma_C^{-1} \right) $$

or just for convenience in terms of the intensity proportional and $F$, $M$ independent factor $G_0$:
Equation (3.1.9.) defines the saturation intensity $I_{\text{sat}}$. The populations $C_\sigma$, in the steady state, are obtained by solving the linear system of equations (3.1.9) setting $B_\sigma = 0$, $C_\sigma = 0$. In analogy to equation (3.1.8) we define for our case:

$$R(6S \rightarrow 15P) = R(6S \rightarrow 7S) \frac{\varrho_{15P}}{\varrho_{7S}+\varrho_{15P}} G_\sigma \cdot f(G_\sigma)$$

where all the additional information obtained by solving the rate equations is contained in the function $f(G_\sigma)$. By a numerical calculation of this function [Hof 85], one finds that $f(G_\sigma)$ is very well approximated by

$$(3.1.11) \quad f(G_\sigma) = \frac{1}{1 + a G_\sigma}$$

where $a$ depends on the h.f. components and the laser polarizations, i.e. $f(G_\sigma) = f(G_\sigma, F, F', q, q')$; as an example the case $q=q'=1$ (ss) is shown in Fig. 3.2 for plane waves. This is a very important result because it leads to a simple form for the total transition rate integrated over the Gaussian intensity distribution of the beams. The effect of this two Gaussian profiles is to modify the two-photon rates by an additional multiplicative factor which depends on the geometry of the beams and the degree of Saturation, $aG_\sigma$. This factor is included in the calculations. The "analyzing power" of the transitions observed with circularly polarized beams also depends on saturation, see Fig. 3.3.
Fig.3.2: Calculated function $f(G_0)$ for plane waves and $q = q' = 1$ (same sense of circular polarizations).

Fig.3.3: Calculated saturation dependence of the circular polarization asymmetries for Gaussian beams with equal waists.
A further very important prediction of this rate equation model is the existence of a saturation-dependent, so called "quadrupole" asymmetry $A_{E2}$ between the transition rates observed with parallel and perpendicular linear polarizations of the two beams (Fig. 3.4). This asymmetry, although generally less than 1% (except for the $3 \rightarrow 3$ transition, where the maximum effect is $\approx 2\%$), could constitute an important systematic problem for the PNC experiment if the circular polarization of the $\lambda_1$ and the linear polarization of the $\lambda_2$ lasers were not perfect.

Fig. 3.4: Calculated saturation dependence of the quadrupole asymmetries for plane waves.
Fig. 3.5: Schematic summary of the asymmetries occuring in $6S_{1/2} \rightarrow 7S_{1/2} \rightarrow nP_{3/2}$ transitions.

The asymmetries of interest are illustrated schematically in Fig. 3.5. The most convenient way for testing the rate equation model presented so far is to measure the circular polarization asymmetries $A_{CP}$, especially for the $3 \rightarrow 4$ transition, since that one is the most sensitive to quasi-elastic scattering processes in the $15P_{3/2}$ state. The results of such measurements are presented in Section 3 of this Chapter, following the description of the experimental setup in Section 2.
A block diagram of the experimental setup for the asymmetry measurements is shown in Fig. 3.6. Two argon-ion lasers (both Spectra-Physics Model 171) operated in all-line mode pump two conventional single-mode ring dye-lasers. After splitting off 20 to 30 percents of the laser beams for monitoring- and control devices the remaining beams are brought to overlap within a thermionic diode.

Fig.3.6: Block diagram of the experimental set-up for the asymmetry measurements.

This ion detector is described in detail in the Chapters I and II. In order to discriminate against ions from the two-photon ($2\lambda_1$) processes (see I.4), the infrared laser beam is
chopped and the diode signal demodulated by a lock-in amplifier.

III.2.1 RH 110 Dye Laser

This home-built ring dye laser is operated with the Rhodamine 110 dye. It is described in detail in [HOF 82], [Wei 84], except for a few improvements:

i) The entire laser is enclosed in a box with vertical laminar air flow. The air passes a micro filter before entering the box. This flow box was installed to keep the optical elements clean and to block the breeze from the room air conditioning and the cooling fans of the electronics.

ii) The laser frequency is locked to an external confocal Fabry-Perot etalon (FSR=1.5 GHz, finesse F=6) [WEI 84]. This etalon was used in two scanning modes:
   - smooth scanning, where air is admitted continuously through a needle valve.
   - in stepwise mode where two magnetically controlled valves with a small buffer volume in the interconnecting tube operate alternately. One valve controls the filling (or evacuating) of this volume while by opening the other valve the buffer volume is admitted to the etalon. By adjusting the volume between the two valves and adding a buffer volume of ≈ 8 dm³, discrete frequency steps down to a few MHz can be achieved [Wei 88;1]. Furthermore the frequency steps can be controlled by the data acquisition system.

III.2.2 Atomic Reference 6S₁/₂ → 7S₁/₂

The atomic reference apparatus described in [HOF 82] is still being used to set the dye laser close to the desired 6S₁/₂ → 7S₁/₂ transition frequency. The 0.85 μm fluorescence from
the Stark-enhanced M1 transition in a vapour cell is detected by a Ga-As photomultiplier (Doppler-width of the lines $\approx 750$ MHz).

An improvement of the atomic Cs-beam apparatus [WEI 84] is in progress. To increase the signal to noise ratio the cross section of the Cs beam has been changed to a rectangular shape ($14 \times 1 \text{ mm}^2$). This is done by inserting an overheated multichannel plate (collimation ratio $\approx 40$) in the exit hole on the oven surface [Gil 86], [Zha 88]. This apparatus is designed as atomic reference to stabilize the RH 110 dye laser during the asymmetry measurements.

III.2.3 Frequency Marker

A 1.25 m long confocal etalon (operated in non-mode-matched configuration, finesse $\approx 2.5$) is used as frequency marker. The FSR of $61,91 (36)$ MHz was calibrated by the Cs 7S h.f. splitting in the atomic beam. The mirrors, held by stainless steel mounts, are sealed with a thermal expansion compensation technique to an evacuated fused silica tube. The transmission fringes of this etalon are recorded during frequency scans simultaneously with the other signals to provide a continuous frequency calibration [Hof 84].

III.2.4 Styryl 9 Dye Laser

The infrared radiation is generated by a commercial ring laser operated with the Styryl 9 dye (Spectra-Physics Model 380 D with Stabilock 388 and astigmatism compensator installed). This laser delivered up to 320 mW for 9 W pump power at the required wavelength of 831 nm. As the Rh 110 dye laser this laser is enclosed in a box with laminar air flow of the same type.
III.2.5 Atomic Reference $7S_{1/2} \rightarrow 15P_{3/2}$

The IR-laser, as a part of the detection apparatus for the $7S$ atoms, requires an absolute atomic reference to be locked to. This reference device was developed by A. Weis [Wei 84], [Der 87], [Wei 88;2]. A high density of atoms in the $7S$ state is needed to observe the Lamb dip on the Doppler broadened $7S \rightarrow 15P$ transition with the requisite S/N ratio. The $7S$ atoms are produced by absorbing the light of a RF excited Cs discharge lamp in a glass cell containing Cesium vapour. The IR beam is focused into the cell and drives the $7S \rightarrow 15P$ transition. The Rydberg states are easily ionized by thermal collisions and the resulting charges are collected in a low electric field of $\sim 2$ V/cm on two plane electrodes.

III.2.6 "Lambda Meter"

The absolute wavelengths of both dye lasers are monitored with a scanning Michelson interferometer (Lambda meter). One of the dye lasers and a frequency stabilized He-Ne laser (Spectra-Physics Model 117, absolute wavelength uncertainty $\pm 200$ MHz) are coupled along parallel paths into the interferometer. The optical setup is essentially that described in [Hal 76], [Her 84]: two corner-cube reflectors defining the arms of the interferometer slide over a length of 70 cm on air bearings to generate fringes. The vernier technique [Kah 83] is used to determine the ratio of the counted fringes to an accuracy of $\pm 10^{-8}$. After a measuring time of a few seconds, the wavelength is displayed on a microcomputer controlled read-out. To identify the spectral lines clearly, the wavelength uncorrected for dispersion is accurate enough. If increased accuracy is required, this Lambda meter may be operated in a helium atmosphere [Her 85].
III.2.7 Vibrating Crosshair Scanner

To get optimal and reproducible signals from the diode, the two counterpropagating laser beams must overlap perfectly in the interaction volume. To facilitate the alignment and to monitor the overlap a novel device has been designed in our laboratory [Wei 83]. Two crossed 20 μm thick "hairs" (fused silica fibers) mounted on a vibrating support are so oriented (45°) that they intersect the beams subsequently. The diffusely scattered light is detected by two corresponding photodiodes. Displaying the photodiode currents on an oscilloscope synchronized to the crosshair vibration shows the relative horizontal and vertical positions of the beams in the plane perpendicular to them. Two such devices, one at each end of the diode, allow a perfect adjustment of the overlap. After the alignment, the scanners are moved out of the beams.

III.2.8 Power Stabilization

Although we eliminate the signal from the two-photon (2λ₁) process by chopping the λ₂ beam, fluctuations of the corresponding ion current contribute, proportionally to the intensity of the λ₁ beam, to the noise in the λ₁ - λ₂ two photon signal. Hence this contribution can be reduced by stabilizing the power of the λ₁ (Rh 110) laser. A Pockels cell in the stabilization loop is driven by two home-built cascadable high-voltage amplifiers (0 - 300 V, band width ≈ 8 kHz). The stabilization electronics gets its input signal from one of the power meters in the Stokes Monitor (described below). The power stability of the beam entering the diode, ΔP/P was 1.5 · 10⁻⁵/√Hz at 8 Hz, the chopping frequency used in the measurements.
III.2.9 Stokes Monitor

The optical windows (see Fig. 1.4) can affect the intensity within the diode by the etalon effect, and possibly the other Stokes parameters $S_j$ of the incident beam by stress-induced birefringence. To measure the effect of such imperfections in situ, the polarization of the reflected beam from the diode window is analyzed as proposed by P.P. Herrmann [Her 86;3], Fig. 3.7.

---

![Diagram](image)

---

Fig.3.7: Stokes Monitor set-up.

The incoming beam (from the Pockels cell of the power stabilization) is vertically polarized by a Glan prism. A Brewster plate splits off some light for stabilization and normalization purposes (power meter P). For perfectly circular polarization (ideal $\lambda/4$ plate) incident upon a dielectric surface the reflected light after the same $\lambda/4$ plate is purely horizontally polarized. Since our windows constitute etalons (with a reflectivity given by the index of refraction of the glass), the intensity of the reflected beam depends on
the laser wavelength and the window thickness, i.e. temperature. A purely horizontally polarized beam is not affected by the Brewster plate and is completely transmitted by the Glan prism to the G power meter. From this signal the intensity inside the diode (Stokes parameter $S_{T0}$) is determined.

In the case of an imperfect $\lambda/4$ plate (retardation = $\pi/2 \pm \varepsilon$) and reflection at an ideal dielectric surface, the reflected light is slightly elliptically polarized in front of the Brewster plate [retardation = $2(\pi/2+\varepsilon)$] with the main axis of the ellipse oriented horizontally (Fig. 3.8). If there is some birefringence of the window, the main axis of the ellipse will be rotated by a small angle.

![Diagram](image)

Fig. 3.8: Detail of the Stokes Monitor.

The Brewster plate simply acts as a vertical analyser and is hence not able to discriminate between the two cases of imperfection, i.e. the corresponding power meter B measures them simultaneously and monitors in this way the deviation from the perfect circular polarization inside the diode

$$S_{Ty}^2 = 1 - S_{Tx}^2 - S_{Tz}^2.$$
This setup was carefully tested with a dummy window (i.e. one not mounted on the diode) to check that we are able to monitor the transmitted intensity. The residual birefringence of single windows (with their support tubes; from the same batch as the diode windows), is \( \approx 10 \) m rad. We found that the intensity behind the window, averaged over one second, can be determined with an accuracy of \( 3 \cdot 10^{-3} \) by using the measurements on the reflected beam only. The Stokes Monitor needs of course a careful calibration to achieve this accuracy. This is done as follows:

The dummy window is mounted in a heatable aluminium tube. The thickness of the window (\( d \approx 2.5 \) mm; FSR = \( \frac{c}{2 \nu} \approx 41 \) GHz) is slowly changed by raising or lowering its temperature. During these temperature scans (with locked laser frequency) over two FSR's (\( \approx 8K/\)FSR) the signals of the four power meters are recorded simultaneously. The calibration is done off line by analysing the modulation of the signals. The measured modulation depth \( M \) of the transmitted intensity agreed very well with the \( M_{th} \) obtained from etalon theory:

\[
M_{th} = \frac{F}{(F+1)} \text{ where } F = \frac{4R}{(1-R)^2} \text{ and } R = \frac{(n-1)^2}{(n+1)^2}
\]

\[
\frac{(M-M_{th})}{M_{th}} \leq 5 \cdot 10^{-3}
\]

The refractive index of Pyrex (Duran) at the 6S \( \rightarrow 7S \) wavelength in air is taken from \([\text{Hof 82;1}]\), \( n = 1.47429(14) \). For the 7S \( \rightarrow 15P \) wavelength we use \( n = 1.4651 \) \([\text{Sch 88}]\). Since the agreement of \( M \) and \( M_{th} \) is very good, \( M_{th} \) can be used to calibrate the power meters for the reflected beam.

In situ calibration, i.e. with the diode in the beam path and the windows kept at the constant oven temperature (\( \approx 200 \) °C), is done by scanning the laser frequency over a maximum and a minimum of the fringes near the 6S \( \rightarrow 7S \) and 7S \( \rightarrow 15P \) wavelengths respectively.
The retardation error $\Delta \phi$ of the $\lambda/4$ plate was measured with high accuracy with the same setup of the Stokes meter but with the window replaced by a single reflecting dielectric interface (plane convex lens or prism) and the $\lambda/4$ plate once in neutral and once in 45° position:

$$\Delta \phi(\lambda/4) = 1^\circ 30' 5" (29'')$$

The position of the $\lambda/4$ plate is controlled in a feedback loop with an accuracy of ± $2\pi/4000$ using an encoder with 1000 cycles, quadrature outputs and index marker (Litton Model 820A). Four pre-programmable positions of the $\lambda/4$ plate can be selected automatically by the data acquisition system. Once the required angle of the $\lambda/4$ plate is attained, the positioning electronics generates a ready signal which is fed into the powermeter electronics of the Stokes Monitor. From there the ready signal is sent to the data acquisition system only if the intensity signal of the "birefringence meter" $B$ is below a preset low level. Since there are electronic offsets and also the optical imperfections mentioned above this level cannot be set to zero.

The setup of the infrared Stokes Monitor is identical but has no "$\lambda/4$ positioner".
A CAMAC multiscaling system controlled by a Personal Computer (Olivetti M24) is used to record and store the signals of interest, digitized by voltage-to-frequency converters followed by fast scalers. The scalers are activated during a presetable dwell-time (multiples of the PC internal timer interval of 53 msec). Then the counts accumulated in the individual scalers are stored in the PC memory. The time critical routines were written in assembler to minimize the system dead-time (≈ 35 μsec were achieved for the data handling of one single channel). A new accumulation cycle is started after incrementing the channel number. The maximal number of channels of a single run can be chosen to be 1, 2, 3 or 4 times 128. In the present setup, up to 6 spectra can be recorded simultaneously:

- the signal from the thermionic diode
- the three power meters from the \( \lambda_1 \) Stokes Monitor
- two power meters from the \( \lambda_2 \) Stokes Monitor
  ("birefringence meter" not connected)
- during frequency scans the signal from the frequency marker
  is connected instead of the Rh110 "birefringence meter".

During data accumulation up to four spectra selected out of the six are displayed on the PC-monitor. After completion of a single run, the spectra are written on a floppy disk for further processing.
III.3 Measurements of the parity-conserving circular polarization asymmetry

These measurements were performed with the setup described in Section III.2, with the infrared laser ($\lambda_2$) stabilized either on the $7S_{1/2}(F'=3) \rightarrow 15P_{3/2}$ or the $7S_{1/2}(F'=4) \rightarrow 15P_{3/2}$ transition, while scanning the green laser ($\lambda_1$) in discrete steps over one of the four h.f. components $6S_{1/2}(F=3,4) \rightarrow 7S_{1/2}(F'=3,4)$. During these scans six signals were recorded simultaneously, viz.

- the demodulated diode signal,
- the power meters of the $\lambda_1$ and $\lambda_2$ Stokes Monitors:
  - $P_1$, $P_2$ monitor the intensities of the $\lambda_1$ and $\lambda_2$ beams incident on the diode,
  - $G_1$, $G_2$, the corresponding "Glan" power meters (Fig. 3-8), monitor reflected intensity (dependent on the wavelengths and the window thickness) for offline intensity corrections,
- the $\lambda_1$ frequency marker.

The "birefringence" meter $B_1$ was used to start every single measurement step (after a change of $\lambda_1$ polarization) only if its signal was below a preset level. After each frequency step, the six signals were integrated during a preset dwell time for either circular polarization of the $\lambda_1$ beam, while the circular polarization of the $\lambda_2$ beam was not changed. The "birefringence meter" $B_2$ was not used.

After every readjustment of the setup and at least once a day calibration data for the two Stokes Monitors were recorded. These data are obtained by scanning the lasers ($\lambda_1$, $\lambda_2$) over the nearest maximum and minimum of the etalon fringes of the corresponding diode window. In a further run the offsets of all power meters were recorded as well. Fig. 3.9a shows a typical diode asymmetry spectrum of the $6S_{1/2}(F=3) \rightarrow 7S_{1/2}(F'=4) \rightarrow 15P_{3/2}$ transition, whereas Fig. 3.9b shows the corresponding signal from the frequency marker. The observed
relative line shift, line widths and background will be discussed later.

Fig. 3.9a): Typical circular polarization asymmetry spectrum recorded with two seconds integration time per channel.

(6S_{1/2}(F=3) \rightarrow 7S_{1/2}(F'=4) \rightarrow 15P_{3/2} \text{ transition})

Fig. 3.9b): Corresponding frequency marker signal.
III.3.1 Calibration of the Stokes Monitor and determination of the correction factors for the intensities within the diode

The first step in this calibration procedure consists in determining the fringe maxima and minima by fitting parabolas to the data $G(\lambda)$ and constants to the offset data of the $P$ and $G$ meters. The values $G_{\text{min}}$ and $G_{\text{max}}$ are then corrected for the offsets $G_0$, and normalized to the incident intensities:

$$\tilde{G} = \frac{G - G_0}{P - P_0}$$

The actual position on the transmission fringe during any measurement run is given by $\beta$ from the following definition (note that the $G$ meters measure the reflected light from the windows):

$$\beta = \frac{\tilde{G}_{\text{run}} - \tilde{G}_{\text{min}}}{\tilde{G}_{\text{max}} - \tilde{G}_{\text{min}}} \quad 0 \leq \beta \leq 1$$

\[
\begin{align*}
\tilde{G}_{\text{max}} & \quad \tilde{G}_{\text{run}} \\
\tilde{G}_{\text{min}} & \quad \beta
\end{align*}
\]
While $\tilde{G}_{\text{min}}$ and $\tilde{G}_{\text{max}}$ are obtained from the calibration data, $\tilde{G}_{\text{run}}$ is taken from the individual runs. With $\beta$ defined in this way, we can write for the intensity correction factor $\varepsilon_x$ that compensates for the etalon effect of the windows:

$$\varepsilon_x = 1 - \beta_x M_x$$

The index $x$ specifies the two laser beams ($\lambda_1, \lambda_2$) and their circular polarizations.

The fringe position signal $\tilde{G}_{\text{run}}$ for frequency scans is obtained by first fitting a Lorentzian line shape to the spectral line, and then an appropriate polynomial to the $G$ signal at the position of the spectral line. Next $\tilde{G}_{\text{run}}$ is calculated with the coefficients of the polynomial at the center of the line.

**Estimation of maximum possible errors in measured asymmetries induced by uncertainties in the light intensities:**

The asymmetries are obtained by calculating the ratio

$$A = \frac{S_+ - S_-}{S_+ + S_-},$$

where $S_\pm$ denotes the measured signals for the two orthogonal laser polarizations ($6S \rightarrow 7S$ transition). Assuming a linear intensity dependence, we write

$$S_\pm = \alpha_\pm \cdot T_\pm$$

with $\alpha_\pm$ the corresponding absorption coefficients and $T_\pm$ the intensities in the interaction volume. In the ideal case ($T_+ = T_-$) we obtain
\[ A \approx \frac{A_0}{2} \quad \text{with} \quad 1 + A_0 = \alpha_+ / \alpha_- \]

The intensity inside the diode, determined from the power meter signals of the calibrated Stokes Monitor is

\[ T_+ \sim 1 - B \cdot M_{th}, \quad \text{with} \quad 0 \leq B \leq 1, \]

where the \( B \)'s measure the actual position on the window transmission fringe. In this case the measured asymmetry is given by

\[ A \approx \frac{A_0 - \Delta B \cdot M_{th}}{1 + B M_{th}} \quad \text{with} \quad \Delta B = (B_+ - B_-)/2 \quad \text{and} \quad B = (B_+ + B_-)/2 \]

Although the temperature of the oven containing the diode is stabilized, there are temperature fluctuations with small amplitudes: \( \Delta T_{pp} = \pm 0.9 \text{ K} \) from a measurement over two days (the index pp denotes peak to peak amplitude). The fastest temperature change observed in this period was \( \approx 46 \text{ mK/min} \). From these measurements and the temperature transmission fringe "spacing" of \( \approx 8 \text{ K} \) we can give a worst case estimate of the error term in the asymmetry in function of the time interval \( \Delta t \) of two subsequent measurements:

\[ (3.3.2) \quad \left( \frac{\Delta B \cdot M_{th}}{1 - B M_{th}} \right)_{\text{max}} \leq 2 \cdot 10^{-8} \Delta t / \text{sec} \]

This figure implies that for the measurement of asymmetries of the order of \( 10^{-4} \) and cycle times longer than one second the diode (window) oven temperature has to be monitored simultaneously with the other signals. If the temperature change is faster than the allowed one, the corresponding measurements have to be rejected.
III.3.2 Calibration of the experimentally determined saturation parameter $G_\text{exp}$ to the saturation parameter $G_0$ used in the theoretical calculations

First we give an estimate of the saturation intensity as defined by equation (3.1.10):

$$I_{\text{sat}}^{-1} = \frac{8\pi a}{\hbar} \frac{\pi}{\hbar} \frac{|<15P_3/2||r||7S_1/2>|^2}{6}$$

The value for the reduced dipole matrix element is taken from a model calculation [Hof 85]:

$$|<15P_3/2||r||7S_1/2>| = 0.1037 \text{ a}_0$$

together with (3.1.3) and (3.1.5) we get:

$$I_{\text{sat}} = 27.9 \text{ W/cm}^2$$

The intensity calibration is done with the strongest line, i.e. the $6S_{1/2}(F=3) \rightarrow 7S_{1/2}(F'=4) \rightarrow 15P_{3/2}$ transition with the two circular polarizations equal (ss). The rate equation model calculation gives for the parameter $a$ in equation (3.1.11):

$$a(G_0,3,4,1,1) = 0.83$$

The saturation parameter $G_0^{\text{exp}}$ is determined from a fit of $G_0^{\text{exp}}/(1+G_0^{\text{exp}})$ to the measured line intensities at different $\lambda_2$ intensities $I_2(G_0^{\text{exp}}= bI_2)$. The parameter $b$ has to be determined each time the beam geometry is changed. Thus the calibration to the $G_0$ parameter is given by

$$(3.3.3) \quad G_0 = \left[ \frac{b}{a} \right] I_2$$

The waists of the two beams were brought to overlap in the center of the interaction volume and given the same size (by
adjusting two lenses in each beam path). The waists and their positions were determined by measuring the beam diameters at four positions along the virtual interaction volume in horizontal and vertical direction with a diode array.

With the measured saturation power $P_{\lambda_2}^{\text{exp}} = 30(3) \text{ mW}$ and a waist of $255(20) \text{ \mu m}$, we find an experimental saturation intensity (for Gaussian beams)

$$I_{\text{sat}}^{\text{exp}} = 30(6) \text{ W/cm}^2$$

in good agreement with the theoretical estimate.

Fig.3.10 shows the saturation behaviour of the $6S_{1/2}(F=3) \rightarrow 7S_{1/2}(F'=4) \rightarrow 15S_{3/2}$ transition. The solid line is the fitted saturation curve, the dashed line shows the slope for $G_0^{\text{exp}} \rightarrow 0$.

![Graph](image)

Fig.3.10: Saturation of the F=3 $\rightarrow$ F'=4 transition with circular polarizations of both beams in same sense (ss).
The statistical error bars from the fitted linestrengths lay within the drawn points. We assume that the scatter of the measurements is due to the not stabilized $I_z$ intensity.

### III.3.3 Measured line widths of the $6S_{1/2}(F=3) \rightarrow 7S_{1/2}(F'=4) \rightarrow 15P_{3/2}$ transition

The widths of the spectral lines obtained from the asymmetry runs for equal (ss) and opposite (os) circular polarizations of the two lasers respectively were determined by fitting a Lorentzian lineshape to these lines with a frequency axis calibrated by the frequency marker. These widths vs. $\lambda_z$ intensities were fitted to the saturation broadening formula ($\sim \sqrt{1+G}$). The homogenous line widths obtained in this way are:

<table>
<thead>
<tr>
<th></th>
<th>ss</th>
<th>os</th>
<th>mean</th>
</tr>
</thead>
<tbody>
<tr>
<td>$2\gamma$</td>
<td>$2\pi \cdot 25.2 (2.0) \text{ MHz}$</td>
<td>$2\pi \cdot 24.9 (2.9) \text{ MHz}$</td>
<td>$2\pi \cdot 25.1 (1.7) \text{ MHz}$</td>
</tr>
</tbody>
</table>

This agrees within a factor 1.4 with the estimated width $2\gamma \approx 2\pi \cdot 18 \text{ MHz}$, eq. (3.1.5).

The observed (but unexpected) relative shift ($\Delta \nu = \nu(\text{ss}) - \nu(\text{os})$) of the two lines as seen in Fig. 3.9 was found to be:

$$\Delta \omega \approx 2\pi \cdot 3.5 (1.0) \text{ MHz}.$$ 

This shift can be explained by the unequal populations of the $15P_{3/2}$ h.f. levels resulting in the two cases. A theoretical estimation of this effect (geometrical factors and weighted means: $\Delta \omega_{\text{th}} = \langle \Delta \omega(\text{ss}) \rangle - \langle \Delta \omega(\text{os}) \rangle$) for the $6S_{1/2}(F=3) \rightarrow 7S_{1/2}(F'=4) \rightarrow 15P_{3/2}$ transition with the interval factor $A(15P_{3/2}) \approx 0.45 \text{ MHz}$ from (3.1.7), predicts
\[ \Delta \omega_{th} = 2\pi \cdot 1.7 \text{ MHz}, \]
i.e. a shift with the observed sign and order of magnitude.

### III.3.4 Intensity dependence of the background

The two photon Doppler-free transitions investigated here are sitting on a Doppler broadened pedestal. The signal to background ratio (S/B) for the $6S_{1/2}(F=3) \rightarrow 7S_{1/2}(F'=4) \rightarrow 15P_{3/2}$ transition shows no dependence on the $\lambda_1$ intensity (varied by a factor 2):

\[ \frac{S}{B(3\rightarrow 4)} = 1.015(10) \]

The $\lambda_1$ as well as the $\lambda_2$ intensity dependence of the background showed a small saturation effect. The fitted saturation parameters in both cases differed from zero by two standard deviations. These effects are not yet fully understood and will need more investigation for the final PNC measurements.

The nearly linear dependence of the background on both intensities suggests that there exists a process for the $\lambda_1$ beam that leads to $7S_{1/2}(F'=3,4)$ atoms with an isotropic velocity distribution. These atoms are then pumped to the $15P_{3/2}$ state by the $\lambda_2$ beam. The most probable candidate to produce this background is the Cs$_2$ molecule (see Fig. 3.11, the Cs$_2$ potential curves are taken from [Col 81]). The $\lambda_1$ energy and an assumed repulsive state $Y$ leading to the isotropic $7S_{1/2}$ atoms are indicated with dashed lines.
Fig. 3.11: Cs₂ potential curves taken from [Col 81].
III.3.5 Measured circular polarization asymmetries $A_{cp}$:

The measured spectra at various $\lambda_2$ intensities (example: Fig. 3.9) were analysed as follows:

- Each scan ("spectrum") consisted of 128 discrete frequency steps, i.e. 256 data points with the two $\lambda_1$ polarizations. The integration time for one single data point varied for different spectra between 1 and 4 seconds.

- Each data point consisted, besides the diode output, of the 5 auxiliary quantities listed above.

- The intensities of the M1 line (and their statistical errors) in the diode spectra were determined by fits to Lorentzians. These intensities are denoted by $I(ss)$, $I(os)$ and correspond to the relative circular polarizations of the two beams.

- The intensities at the line centers of the $\lambda_1$ ($\lambda_2$) beam inside the diode (i.e. corrected for the etalon effect) differed by less than $5 \cdot 10^{-4}$ ($5 \cdot 10^{-3}$) for the two polarizations, i.e. the $\lambda_1$ ($\lambda_2$) intensities were equal for $I(ss)$ and $I(os)$. The ten-times worse value for the $\lambda_2$ beam is attributed to the not stabilized intensity of that laser.

The circular polarization asymmetry from one single run is obtained from:

$$\{3.3.4\} \quad A_{cp} = \frac{I(ss) - I(os)}{I(ss) + I(os)}.$$

The $\lambda_2$ intensity calibration factors $b/a$ from equation (3.3.3) were determined with an accuracy of $\approx 10\%$. This uncertainty is due to the statistical errors in the fitted
A few test runs with linear $\lambda_2$ polarization and reversed circular $\lambda_2$ polarization were performed. The asymmetries obtained from these runs agreed within the error bars with the expected values 0 and $-A_{cp}$ respectively. Measured asymmetries at the same $\lambda_2$ intensity were averaged (weighted means) and the resulting error was determined as $\max \{ s_m(\text{internal}), s_m(\text{external}) \}$ [Grä 78].

Fig. 3.12 shows the measured circular polarization asymmetries together with the predictions of the rate equation model. Note that the asymmetries at $G_0 = 0$ (see Table III.1) are given by geometric coefficients, i.e. are independent of dynamics. They can serve as a check on systematics. All but the measurements on the $F=4 \rightarrow F'=3$ transition agree with the predicted saturation behaviour of the asymmetries. This leads
us to the conclusion that the rate equation model (3.1.9) gives an adequate description of the measured circular polarization asymmetries. The disagreement observed for the $F=4 \rightarrow F'=3$ transition suggests that corresponding model calculations be checked for possible computational error.

Fig. 3.13: Detail of Fig. 3.12 showing only the data for the $F=3 \rightarrow F'=4$ transition.

Fig. 3.13 shows the subset of data concerning the $F=3 \rightarrow F'=4$ transition. The measured data show clearly that population redistribution among the h.f. sublevels of the $15P_{\frac{3}{2}}$ state is significant, or putting more strongly: The maximum quasi elastic scattering model is clearly favoured.

The measurements of the quadrupole asymmetries $A_{E2}$ are in progress but results are not yet available to be presented here.
III.3.6 **Signal to noise ratio considerations:**

We analyse the signal to noise (S/N) ratio in function of the $\lambda_2$ intensity $I_2$. The signal in this analysis is taken to be the mean $\langle I(\text{ss}) + I(\text{os}) \rangle$ of the measured intensities. The saturation parameter $G_0^{\text{exp}}$ is determined and the signal parameterized as follows:

$$S(G_0^{\text{exp}}) = S_\infty \cdot \frac{G_0^{\text{exp}}}{1 + G_0^{\text{exp}}}.$$

We define the noise for each spectrum $N(G_0^{\text{exp}})$ as the rms fluctuations of the background near the M1 line, assuming no additional noise from the M1 signal proper. Furthermore a constant noise $N_0$ originating from the intrinsic diode noise and the 2-$\lambda_1$ background processes (see Section 1.4) is assumed since the $\lambda_1$ intensity was stabilized during this measurements (both noise sources were measured to contribute about equal amounts to $N_0$).

The obtained S/N ratios are analysed following this simple model:

$$\frac{S}{N} (G_0^{\text{exp}}) = \frac{S_\infty}{N_0} \cdot \frac{G_0^{\text{exp}}}{1 + G_0^{\text{exp}}}.$$

For the S/N ratios ($F=3 \rightarrow F'=4$ transition, 2 sec integration time) shown in Fig.3.14 we find by a least squares fit:

$$\frac{S_\infty}{N_0} = 48 \ (3)$$

The scatter of the measured points on Fig. 3.14 is attributed to the not stabilized $\lambda_2$ intensity. Unfortunately a more precise analysis of the relative $I_2$ fluctuations of the individual runs was not possible since the signal from the power meter $P_2$ showed a small beat structure due to chopping frequency and integration time interference.
Fig.3.14: Signal to noise ratios vs. $I_2$ for the $F=3 \rightarrow F'=4$ transition.

These facts call for the following changes in the experimental set-up (Fig. 3.6):

- The intensity of the $\lambda_2$ beam has also to be stabilized.

- The chopper has to be moved in front of the diode in order that only the $G_2$ and $B_2$ signals will show the small beat structure (The $G_2$ signal can induce a maximum $I_2$ correction of $\approx 13\%$).

From the equal contributions of the intrinsic diode noise and the $2\lambda_1$ background noise to $N_0$ we anticipate a maximum possible increase in the signal to noise ratio of $\sqrt{2}$ by increasing the beam diameters.
In this thesis we have investigated experimentally and theoretically the properties of thermionic diodes of novel design and exploited their power as sensitive spectrometers.

The construction of long diodes with electrostatically shielded interaction volume and large solid angle for ion collection is presented in detail. We also present the first quantitative theory of the diode's current vs. voltage characteristic; it is shown to agree very well with the experimental results.

From the calculated potential distributions and the measured ion trapping times we can estimate the possible diode gain and its (chopping) frequency dependence. We used two distinct methods to determine the diode gain experimentally and found an order-of-magnitude agreement with our estimate.

The behaviour of the diode in axial magnetic fields was also investigated. The gross features can readily be explained on the basis of an extension of the theory. We also observed an additional "fine structure" which is as yet not understood, but can nevertheless be exploited to measure magnetic fields down to ≈ 5 mGauss within the diode proper.

As an application of the diode as an ultra-sensitive spectrometer we investigated the circular polarization correlations in the resonant 2-photon absorption $6S \rightarrow 7S \rightarrow 15P$ in Cs, where the first step is an extremely week M1 transition. We compared our result with a rate-equation model including saturation and found that collisional population redistribution (quasi-elastic scattering) among the hyperfine sublevels of the final 15P state occurs.
Possible extensions of this work are:
- By including the image charge potential for the emitted electrons (Schottky effect) into the theory of the current-voltage characteristics of the diode, also the saturation behaviour of the anode current could be calculated.
- Further investigation of the as yet not understood "fine structure" in the anode-current vs. B-field characteristics of the diode.
- A check of the rate-equation model calculations since we found a disagreement for one specific h.f. component ($F=4 \rightarrow F'=3$) of the M1 transition.
ACKNOWLEDGEMENTS

This work was suggested by Professor V.L. Telegdi and carried out under his guidance. I am deeply indebted to him for the many enlightening discussions and his constant interest. His way of looking at physics has substantially contributed to develop my understanding and appreciation of that subject. His critical reading of the manuscript is also greatly appreciated.

I am grateful to Professor K. Niemax for kindly accepting to be co-referee of this thesis.

Further my thanks go to my colleagues Dr. A. Weis, Dr. D. Zevgolis and L. Zhao for their help throughout this work, and to Dr. P.P. Herrmann and Dr. J. Hoffnagle for their contributions in its earlier stages, specifically:
- Paul Herrmann made decisive contributions to the design and assembly of the "final" thermionic diodes, as well as to the Stokes Monitor.
- John Hoffnagle provided the rate-equation model calculations presented in chapter III.
- Tun Weis for many stimulating discussions and the designing and programming of the data-analysis software.

The untiring help of Tun, Liang and Dimitrios in the final moments of this work is gratefully acknowledged.

Among the many other people who contributed to the success of this work I want to thank particularly:
- Prof. W. Franzen (Boston University) for helpful discussions in the early phases of the diode theory;
- Messrs. I. Fritschi and P. Rettich for building most of the electronics, Messrs. R. Denner and S. Flühler for technical support.
- the glassblower Mr. W. Widmer for the careful work in the final assembly of the diodes.
- Mrs. G. Kägi and Mr. M. Markwalder for the illustrations.
- Mrs. E. Redard for careful typing of the manuscript.
- my wife Irene and Patrick for their patience.
Appendix A

Details of diode construction

For the schematic cross section of the diode we recall Fig. 1.2. A completely assembled double diode unit (active length 10.8 cm) ready to be vapour cleaned is shown in Fig.A.1 (the thick V-shaped wire is only for suspension during the cleaning procedure). This cleaning is done to remove possible residual organic compounds, although the diodes were assembled on a clean bench with laminar air flow. The anodes, consisting of a tantalum support with a spot welded tungsten sheet (25μm thick) to it, were mounted after the cleaning by means of tungsten springs (S) and molybdenum bolts (B) to the diode body (the same mounting technique is used for the diode body proper, Fig.A.3).

Fig.A.2 shows our longest diode on its mounting bench. Note the symmetric mounting of the tungsten rods (diam. 1 mm) on the upper and the lower side of the diodes. All electrical interconnections were made with platinum wires since that material has very good spot welding properties and is highly ductile.

Fig.A.3 shows the details of a feed-through part. The interaction volume is here filled by a polished alumina rod which is removed after the glassblower has connected a Pyrex envelope to the two feed-through parts. After connecting the support tubes carrying the optical windows (see Fig.1.4) and a side arm to the diode envelope, some tenth of a gramm of Cesium is distilled into the side arm following evacuation.
Fig. A.1: Photograph of an assembled double diode unit (active length 10.8 cm) with the anodes removed.
Fig. A.2: Photograph of our "30 cm diode" on its mounting bench, see text.
Fig.A.3: Photograph of feed-through detail of the diode.
Appendix B

Calculation of mean anode radius

The geometry of our diodes is defined by the two parameters $a$ and $r$ (see Fig. B.1). Calculating the mean $n^{th}$ power of the radius for the outer and the inner area of the diode yields:

$$<r_a^n>_{\text{out}} = \frac{(r^2+m^2)^{n/2}}{\varphi} \int_{\theta_1}^{\theta_2} \left[ 1 - \frac{2rm}{r^2+m^2} \cos \varphi' \right]^{n/2} d\varphi'$$

where $m = r - a$, $\varphi = \arccos \left(1 - \frac{2a}{r}\right)$

$$<r_a^n>_{\text{in}} = \frac{a^n}{\delta} \int_{\delta}^{\delta'} \left( \frac{1}{\cos \delta} \right)^n d\delta'$$

where $\delta = \arctan \left( \frac{r \sin \varphi}{a} \right)$

and from this we get for the mean $n^{th}$ power anode radius:

$$<r_a^n> = (1 - \frac{\delta}{a}) <r_a^n>_{\text{out}} + \frac{\delta}{a} <r_a^n>_{\text{in}}$$

For $n = 1$ we get for the outer area:

$$<r_a>_{\text{out}} = \frac{\sqrt{r^2+m^2}}{\varphi} \left[ 2\sqrt{1+b} \ E(\Theta,k) - \frac{2b \sin \Theta}{\sqrt{1-b+2ab/r}} \right]$$

where $b = \frac{2rm}{r^2+m^2}$, $k = \frac{2b}{1+b}$, $\Theta = \arcsin \left( \frac{(1+b)2a/r}{\sqrt{2(1-b+2ab/r)}} \right)$

and $E(\Theta,k)$ the elliptic integral of the second kind,
and for the inner area we find:

\[
<r_a>_{in} = \frac{a}{\theta} \ln \frac{1 + \sin \theta}{1 - \sin \theta}
\]

Fig.B.1: Schematic cross section through the diode: 

θ defines the border between inner and outer area.

Two types of diodes with slightly different geometries were used:
- the prototype double diode (pdd) and
- the Macor double-diodes (Mdd).

<table>
<thead>
<tr>
<th></th>
<th>r</th>
<th>a</th>
<th>&lt;r_a&gt;</th>
<th>√&lt;r_a^2&gt;</th>
</tr>
</thead>
<tbody>
<tr>
<td>pdd</td>
<td>24.0 mm</td>
<td>9.5 mm</td>
<td>15.02 mm</td>
<td>15.78 mm</td>
</tr>
<tr>
<td>Mdd</td>
<td>23.5 mm</td>
<td>8.5 mm</td>
<td>13.92 mm</td>
<td>14.70 mm</td>
</tr>
</tbody>
</table>
Appendix C

Cathode heating and anode-cathode circuit:

Fig.C.1 shows the circuit schematics of our current to voltage converter and the cathode heating. The symbol $\|\|$ denotes the protective ground and $\perp\|$ the signal ground in this completely floating design. The two grounds are connected together only at the input of the lock-in amplifier in order to avoid possible ground loops (the 50 Hz noise on the output signal was considerably reduced by this interconnection scheme).

The low / high pass filter option was not used working with the lock-in amplifier (diagram a). The diagrams b) and c) show first order and second order low pass filters respectively (the internal resistors in the instrumentation amplifier INA 104 could be used in conjunction with the external elements shown).

Fig.C.1: Circuit schematics of cathode heating and biased current to voltage converter.
Appendix D

Cathode surface temperature calibration:

We use the measured resistivity of the cathode heating wires to determine the cathode surface temperature. The resistivity vs. temperature characteristic for tungsten is taken from the data of Johns and Langmuir [CRC 75]. The resistivity of the heating wires is determined from current and voltage measurements by taking into account the resistance of the connecting wires.

The temperature calibration of the cathode surface was performed with a two wavelength optical pyrometer (\(\lambda\)'s at 0.95 \(\mu\)m and 1.05 \(\mu\)m, calibrated for measurements down to 600°C and statistical errors over 5 measurements < ± 5 K). The cathode surface was observed through a hole (diam. 10 mm) in the oven and the anode (Tungsten foil) was replaced with a fine tungsten mesh (transmission ≈ 85%). The calibration characteristic at a Cs pressure of 5.8 mTorr is shown in Fig. D.1.

For the determination of \(T_c\) from the measured resistivity of the heating wires at the current working conditions we use a linear extrapolation of the measured characteristic to lower temperatures and estimate there an uncertainty of < ± 20 K.
Fig. D.1: Cathode surface temperature calibration characteristic for the indirectly heated cathodes.
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