Excitation of He⁺ in a hollow cathode arc discharge

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EXCITATION OF He^+ IN A HOLLOW CATHODE ARC DISCHARGE

W. KOHSIEK

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IN A HOLLOW CATHODE ARC DISCHARGE

PROEFSCHRIFT

ter verkrijging van de graad van doctor in de technische wetenschappen aan de Technische Hogeschool Eindhoven, op gezag van de rector magnificus, Prof. dr. ir. G. Vossers, voor een commissie aangewezen door het college van dekanen in het openbaar te verdedigen op dinsdag 25 juni 1974 te 16.00 uur

door

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geboren te Amsterdam

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en

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SURVEY

The subject of this thesis is the relation between the population densities of excited He^+ levels in a plasma and the relevant densities and temperatures of the plasma particles.

Chapter I shows how the population densities can be calculated in the frame of the collisional-radiative model.

Chapter II offers a description of the plasma device, the hollow cathode arc discharge. Chapter III describes the diagnostic methods by which the population densities and the quantities needed in the collisional-radiative model were measured.

Some restrictions of the collisional-radiative model in relation to our plasma are discussed in Chapter IV. In Chapter V the experimental results are given and compared with those of the collisional-radiative model. In connection with this comparison the cross-section for electronic excitation is discussed.

INTRODUCTION

If a plasma is in completely thermodynamic equilibrium, the levels of the excited particles are populated according to the Boltzmann factor $g_n \exp \{-E_n/kT\}$, where g_n is the weight and E_n the excitation energy of level n , and where T is the temperature of the plasma. The plasma is in thermodynamic equilibrium (T.E.) if, and only if, each microscopic process is balanced by its inverse process. This is called the principle of detailed balancing (D.B.).

In practice, no natural nor artificial plasma is in T.E.

In many cases the radiation processes are not in D.B. However, in some of them the population densities still follow the Boltzmann factor. Such a plasma is said to be in Local Thermodynamic Equilibrium (L.T.E.). It is possible to create a laboratory plasma that is in the L.T.E. state. In order to attain this state the plasma has to fulfil several conditions, among which the most important are that its electron density has to be high enough and that spatial density and temperature gradients remain in certain limits. Generally, one can say that a He plasma with an electron density greater than 10^{25} m^{-3} is in L.T.E. This criterion is somewhat different for other gases, and is slightly dependent on the electron temperature.

The plasma which has been the subject of our investigations is a He plasma with an electron density of about $3 \times 10^{19} \text{ m}^{-3}$. This plasma is far from being in L.T.E. In such a case the population densities are complex functions of the densities of all plasma particles and their velocity distributions. This is the domain of the Collisional-Radiative (C.R.) model. The intention of the C.R. model is to calculate the population densities and other quantities like the overall ionization and recombination starting from an as limited as possible number of plasma quantities. In most cases these basic quantities are the electron temperature, the electron density, and the densities of the ions and the atoms. The calculated population densities are functions of the basic quantities and of coefficients for several atomic processes, like transition probabilities and rate coefficients for electronic (de)excitation, recombination and ionization.

The most extensive C.R. calculations have been done in case of a Hydrogen plasma ($e\text{-H-H}^+$), Hydrogen-like plasma's (for instance, $e\text{-He}^+\text{-He}^{++}$) and a Helium plasma.

There are only a few publications permitting a comparison between the C.R. recombination coefficient and the experimentally determined recombination coefficient [1, 2]. The agreement between theoretical and experimental

values is at best within a factor two. Concerning the population densities the situation is not much better. Johnson [3] reports a comparison between calculated and experimentally determined population densities in the case of neutral helium, using a variety of assumed cross-sections for the excitation processes. The agreement is generally better than a factor 2. To our knowledge there are no publications dealing with such a comparison for H or H-like ions.

It was our purpose to check C.R. calculations on an $e\text{-He}^+\text{-He}^{++}$ plasma produced by a hollow cathode arc discharge. This plasma has an electron density in the 10^{19} m^{-3} region and an electron temperature of about 200,000 K. The degree of ionization is over 50%; this makes the plasma very suitable for our purpose because excitation by neutral particle collisions and excitation of He^+ levels directly from the He ground state are negligible as compared to electronic excitation, and excitation from the He^+ ground state respectively.

The reasons for choosing He^+ as subject and not H or neutral He were the following:

- (1) He has a singlet and a triplet level system. The C.R. model has to take into account the processes between these two systems. This, compared to H and H-like ions, introduces additional uncertainties.
- (2) both H and He have the disadvantage that the density of the atoms in the ground state is, in contrast to the density of the He^+ ions, not measurable in our case.

1. Introduction

The purpose of C.R. models is to calculate plasma quantities like population densities of excited states and coefficients for ionization and recombination, starting from a limited number of known quantities, such as the electron temperature and density and the ion densities.

There is not just one C.R. model: various gases, like He and H, require different C.R. models; if there is a strong absorption of radiation this has to be incorporated in the model; the same applies to the case where inelastic neutral particle collisions are abundant.

From this point of view the Corona model may be regarded as a specific C.R. model. Only two kinds of atomic processes are dealt with in the Corona model: electronic excitation from the ground state and spontaneous emission. The Corona model is limited to electron densities less than about 10^{18} m^{-3} . At higher densities more atomic processes have to be considered.

2. Short historical review

This review only gives outline of work done on C.R. calculations since 1960 for H, H-like ions and He. An extensive historical review may be found in [4]. In 1962 Bates et al. [5] published a C.R. model dealing with the calculation of the overall coefficients for ionization and recombination in case of H-like ions. This publication was followed in 1963 by one by McWhirter and Hearn [6], who used the same C.R. model to calculate the population densities of excited levels of H-like ions. This model did not take into account radiative absorption. Absorption of radiation was taken into consideration in a report of Drawin and Emard [7] in 1970, in which the population densities of excited states of H, atomic He and H-like ions were calculated. The model they used in the case of H and H-like ions was -except for radiative absorption- the same as used by Bates et al., but their values of rate coefficients for inelastic electronic collision processes were different. In 1972 this report was followed by a publication of the same authors, dealing with the C.R. ionization and recombination coefficient in an $e\text{-He}^+ \text{-He}^{++}$ plasma [8].

3. The collisional-radiative model

The C.R. model that will be discussed in this section is the one used by McWhirter and Hearn [6] and by Drawin and Emard [7] for H-like ions in the case of an optically thin plasma. This model accounts for all inelastic electronic collision processes and all radiation processes, except absorption of radiation and stimulated emission. We confine ourselves to the $e\text{-He}^+ \text{-He}^{++}$ system.

The following processes are considered (fig. I.1.):

- (1) electronic excitation and deexcitation between the atomic states i and j

- (2) ionization and recombination

- (3) spontaneous emission

- (4) radiative recombination

fig. I.1

Scheme of the C.R. model for He^+ .

Solid lines : electronic collision processes

wavy lines : radiative processes.

The number of processes (1), (2) and (4) per unit time and volume are given by the rate coefficients C_{ij} (in $\text{m}^3 \text{s}^{-1}$), C_{ic} ($\text{m}^3 \text{s}^{-1}$) and β_i ($\text{m}^3 \text{s}^{-1}$) respectively. For example, the number of processes (1) from the right to the

left per unit time and volume is $C_{ij} \cdot n_i \cdot n_e$, where n_i is the particle density of He_i^+ ions (in m^{-3}) and n_e is the electron density (in m^{-3}). The spontaneous emission processes (3) are characterized by the transition probabilities A_{ij} .

The rate coefficients of the inverse processes (1) and (2) are indicated by C_{ji} and C_{ci} . The rate coefficients C_{ij} and C_{ji} are related by the principle of detailed balancing: if a plasma is in thermodynamic equilibrium, the relation

$$C_{ij} \cdot n_e \cdot n_i = C_{ji} \cdot n_e \cdot n_j \quad (\text{I.1})$$

holds. Furthermore, the particle densities n_i and n_j obey the Boltzmann distribution law

$$\frac{n_i}{n_j} = \frac{g_i}{g_j} e^{-\frac{E_{ij}}{kT}} \quad (\text{I.2})$$

where g_i and g_j are the statistical weights of levels i and j , $E_{ij} = E_j - E_i$ is the difference in ionization energy of the levels j and i ($E_i = 54.4/i^2$ (eV)).

From (I.1) and (I.2) follows

$$C_{ji} = C_{ij} \frac{g_i}{g_j} e^{-\frac{E_{ij}}{kT}} \quad (\text{I.3})$$

Because of thermodynamic equilibrium, one can put $T = T_e$, i.e. the electron temperature. Now equation (I.3) provides a relation between two pure atom-electron processes. This relation has to be independent of the plasma state, so that it applies also under non-T.E. circumstances. The only condition is that the electrons have a Maxwellian velocity distribution.

In the same way the relation between C_{ic} and C_{ci} can be derived. C_{ci} is so defined that $n_e \cdot n(\text{He}^{++}) \cdot C_{ci}$ is the number of three particle recombination processes per unit time and volume. Therefore, C_{ci} is directly proportional to the electron density n_e . All the other rate coefficients are functions of the electron temperature only.

The population densities n_i obey the differential equation

$$\begin{aligned}
\frac{\partial n_i}{\partial t} = & -n_i \{ n_e [C_{ic} + \sum_{j \neq i} C_{ij}] + \sum_j A_{ij} \} \\
& \text{ionization (de)excitation} \quad \text{emission} \\
& + n_e \sum_{j \neq i} n_j C_{ij} + \sum_{j \neq i} n_j A_{ji} + n_e n(\text{He}^{++}) (C_{ci} + S_i) \\
& \text{(de)excitation} \quad \text{cascading} \quad \text{recombination} \\
& + \vec{V}(n_i \cdot \vec{v}_i) \quad (i = 1, 2, 3, \dots) \quad (\text{I.4}) \\
& \text{transport}
\end{aligned}$$

where $\vec{v}_i$ is the mean macroscopic velocity of the He_i^+ particles.

The following assumptions are made:

- the transport term $\vec{V}(n_i \cdot \vec{v}_i)$ is small compared to the other terms
- quantities $\frac{\partial n_i}{\partial t}$ are zero. This is the case if the plasma is stationary, or if the n_i 's do not change in the time needed for the establishment of the equilibrium between the positive and negative terms of eq. (I.4).

As we have, at first sight, an infinite number of levels i , the system (I.4) forms a set of an infinite number of equations with an infinite number of unknown quantities n_i . There exists, however, a natural cut-off of this set: because of microelectric fields, the levels above a certain level p_0 are no longer bound ones, but merge into the continuum. Drawin and Emard took up to level 25 into consideration. McWhirter and Hearn assumed that level 20 and above have a Saha-Boltzmann population density. The contribution of these levels to the population densities of the lower levels was included by adding to them the contributions of successive high levels until the contribution of the highest level was less than 0.1% of the sum of the preceeding ones. Since we are only interested in the levels 4 to 10, the population densities are not sensitive to the cut-off procedure. Now, neglecting the transport terms, we have a finite set of equations which is linear in the unknowns n_i

$$a_{i1}n_1 + a_{i2}n_2 + \dots + a_{ip_0}n_{p_0} = b_i n(\text{He}^{++}) \quad i = 1, \dots, p_0 \quad (\text{I.5})$$

where the coefficients a_{ij} and b_i are functions of the electron density, the rate coefficients and the transition probabilities. The rate coefficients are known functions of the electron temperature and the electron density (see next section), therefore the a_{ij} 's and the b_i 's are known functions of n_e and T_e .

The solution of this set is called the steady state solution. The quantities one has to know in order to calculate the n_i 's are:

- the electron density n_e
- the electron temperature T_e
- the density of He^{++} ions $n(\text{He}^{++})$.

In practice one can rarely use the steady state solution because the transport of the He_1^+ particles, i.e. the He^+ ions in the ground state, is not negligible. In order to solve the set of equations (I.4) one has to know the value of the term $\vec{\nabla}(n_1 \vec{v}_1)$. There is, however, another way to treat the set (I.4). If n_1 is a known (measurable) quantity, the set of equations becomes

$$a_{i2}n_2 + \dots + a_{ip_0}n_{p_0} = b_i n(\text{He}^{++}) - a_{i1}n_1 \quad i = 2, \dots, p_0 \quad (\text{I.6})$$

The solution of this linear set may be presented in the form

$$n_i = n_i^{(0)} \frac{n(\text{He}^{++})}{n_e} + g_i^{(1)} n_1 \quad (\text{I.7})$$

where $n_i^{(0)} \frac{n(\text{He}^{++})}{n_e}$ is the solution of set (I.6) if n_1 is put zero, and $g_i^{(1)} n_1$ is the solution if $n(\text{He}^{++}) = 0$. The population densities n_i are often given as fractions of the corresponding Saha-Boltzmann population densities n_i^s

$$\frac{n_i}{n_i^s} = r_i^{(0)} + r_i^{(1)} \frac{n_1}{n_1^s}, \quad (\text{I.8})$$

$$n_i^s = n_e n(\text{He}^{++}) \frac{g_i}{g_e g(\text{He}^{++})} \left(\frac{h^2}{2\pi m_e k T_e} \right)^{3/2} e^{-\frac{E_i}{k T_e}}$$

where: g 's are weight factors, $g_i = 2i^2$, $g_e = 2$, $g(\text{He}^{++}) = 1$.

E_i is the ionization potential of level i ; the other symbols have their usual meanings. The coefficients $r_i^{(0)}$ and $r_i^{(1)}$ are functions of the electron temperature and density. These coefficients are tabulated by McWhirter and Hearn [6] and by Drawin and Emard [7] for a large range of

n_e and T_e . The coefficient $r_i^{(0)}$ represents the contribution by the He^{++} ion to the population density of level i , and $r_i^{(1)}$ gives the contribution by the He^+ ion.

The quantities one has to know are:

1. the electron density, n_e
2. the electron temperature, T_e
3. the density of He^+ ions, n_1 ; in our case one may put $n_1 = n(\text{He}_1^+) = n(\text{He}^+)$, where $n(\text{He}^+)$ is the particle density of all He^+ ions.
4. the density of the He^{++} ions, $n(\text{He}^{++})$.

The quantities n_e , $n(\text{He}^+)$ and $n(\text{He}^{++})$ are related by

$$n_e = n(\text{He}^+) + 2 n(\text{He}^{++}) \quad (\text{I.9})$$

which is required by quasi-neutrality.

Finally, the assumptions underlying solution (I.8) may be summarized:

1. Other atomic processes than stated at the beginning of this section are negligible. For instance: absorption of radiation, inelastic atom-ion collisions.
2. The transport of excited particles can be neglected.
3. The electrons have a Maxwellian velocity distribution.
4. The sublevels of the excited main levels are populated according to their statistical weights. The latter assumption is of importance to the values of the transition probabilities A_{ij} . Each transition probability is the average of the transition probabilities $A_{i_k j_k}$ belonging to sublevel-sublevel transitions:

$$A_{ij} = \frac{\sum_{k,k'} g_{i_k} A_{i_k j_k}}{2i^2} \quad (\text{I.10})$$

where g_{i_k} is the statistical weight of sublevel k , of main level i . Values for A_{ij} may be found in [9].

4. The rate coefficients

a. General remarks

The reason why the first C.R. calculations were done on H and H-like ions is that the level system is simple and the rate coefficients are reasonably

well known, especially as compared to other atoms. Neutral helium is already a more difficult case because of the coupling between the singlet and triplet systems.

In the case of H or H-like ions, the transition probabilities A_{ij} are very well known; they have an accuracy better than 1% [9]. Concerning the rate coefficients for inelastic electronic processes the situation is much less satisfactory. McWhirter and Hearn estimate the accuracy of these coefficients to be of a factor 2, and as a result the calculated population densities may be a factor 2 or 3 in error. In reverse, it is possible to investigate the rate coefficients experimentally by comparing observed population densities with calculated population densities. Chapter V will show that we can investigate the contribution of the He^+ ground state separately from the contribution of the He^{++} ions. For the first, and most interesting contribution, only the rate coefficients C_{ic} , C_{ij} and C_{ji} are of importance. In the following we shall discuss these rate coefficients.

b. The rate coefficients C_{ij}

Both McWhirter and Drawin based these rate coefficients on Burgess' cross-section for the transition $1s - 2p$ in He^+ [10]. This cross-section is shown in fig. I.2 by curve 3 and will be discussed more extensively in Chapter V. Drawin takes the following expression for the cross-section of the electronic excitation processes $i - j$ [11]

$$\sigma(U_{ij}) = 4\pi a_o^2 \left(\frac{E_1^H}{E_{ij}} \right)^2 f_{ij} \begin{cases} 0.302 & 1 \leq U_{ij} < 3.85 \\ \frac{U_{ij}^{-1}}{2} \ln 1.25 U_{ij} & U_{ij} \geq 3.85 \end{cases} \quad (\text{I.11})$$

where: $4\pi a_o^2 = 3.516 \cdot 10^{-20} \text{ m}^2$

E_1^H is the ionization energy of hydrogen

E_{ij} is the energy difference between main levels i and j

f_{ij} the oscillator strength of the transition $i \rightarrow j$

$U_{ij} = \frac{E}{E_{ij}}$, E is the energy of the incident electron.

For $1 - 2$ this corresponds to curve 2 fig. I.2. This curve lies lower than curve 3, the one calculated by Burgess for the transition $1s - 2p$. At thres-

$$C_{ij} = \langle \sigma_{ij} v_e \rangle = 8.69 \cdot 10^{-14} f_{ij} Z^{-3} \frac{u_{ij}^{3/2}}{u_{ij}} \Psi_2(u_{ij}) \quad (\text{m}^3 \text{ s}^{-1}) \quad (\text{I.12})$$

$$u_{ij} = \frac{Z^2 E_{ij}^H}{kT_e}, \quad u_{ij} = \frac{E_{ij}}{kT_e}$$

where: Z is the effective charge of the ion (He^{++} : $Z = 2$)

T_e the electron temperature

$\Psi_2(u_{ij})$ a dimensionless quantity, originating from the shape of the cross-section, and tabulated in [11]

This expression may be transformed into

$$C_{ij} = 4.35 \cdot 10^{-12} \frac{f_{ij}}{(kT_e)^{1/2} E_{ij}} \Psi_2(u_{ij}) \quad (\text{m}^3 \text{ s}^{-1}) \quad (\text{I.13})$$

kT_e in eV, E_{ij} in eV.

McWhirter [6] treats the rate coefficients C_{12} , C_{13} and C_{14} separately from the others. He obtained the rate coefficient C_{12} by numerical integration of curve 3 over the Maxwellian velocity distribution, and adding to the result a correction of 15% for the transition $1s - 2s$. C_{13} and C_{14} were calculated in the same way after fitting curve 3 to the appropriate thresholds. The remaining rate coefficients C_{ij} were calculated with [5]

$$C_{ij} = 6.00 \cdot 10^{-12} \frac{f_{ij}}{(kT_e)^{1/2} E_{ij}} e^{-u_{ij}} \quad (\text{I.14})$$

with the same units as in (I.13).

The factor $e^{-u_{ij}}$ may be considered as an approximation to $\Psi_2(u_{ij})$ (formula I.13). It is correct for values of $u_{ij} \sim 1$, but gives too low values for $u_{ij} < 1$, and too high values for $u_{ij} > 1$.

Table I.1 shows the rate coefficients $C_{1,2}$, $C_{1,3}$, $C_{1,4}$, $C_{1,5}$ to $C_{1,15}$, and $C_{6,7}$, obtained by Drawin and by McWhirter for an electron temperature of $2.56 \cdot 10^5$ K, which is a representative value for our discharge.

Table I.1 Rate coefficients C_{ij} (in $m^3 s^{-1}$) at $T_e = 2.56 \cdot 10^5$ K.

	$C_{1,2}$	$C_{1,3}$	$C_{1,4}$	$C_{1,5}$ to $C_{1,15}$	$C_{6,7}$
McWhirter	$1.50 \cdot 10^{-15}$	$2.19 \cdot 10^{-16}$	$7.37 \cdot 10^{-17}$		$4.5 \cdot 10^{-12}$
Drawin	$7.01 \cdot 10^{-16}$	$7.24 \cdot 10^{-17}$	$2.21 \cdot 10^{-17}$		$1.16 \cdot 10^{-11}$
McW/Dr.	2.1	3.0	3.3	3.3	0.39

For $C_{1,2}$ to $C_{1,5}$, $u_{ij} \approx 2$; for $C_{6,7}$, $u_{ij} \approx 0.02$.

The difference between the two values for C_{1j} ($j = 1, \dots, 15$) is mainly caused by Drawin's cross-section for these processes, which is about twice smaller than McWhirter's. The difference in case of $C_{6,7}$ is due to McWhirter's approximation $\exp(-u_{ij})$, which produces for $u_{ij} = 0.02$ a value which is about 3 times too small.

c. The rate coefficients C_{ic}

Drawin arrives at the following expression for the ionization rate coefficient C_{ic} [11]

$$C_{ic} = 7.96 \cdot 10^{-14} f_i \frac{i^2}{(kT_e)^{\frac{1}{2}}} \Psi_1(u_i) \quad (m^3 s^{-1}) \quad (I.15)$$

where: f_i is the absorption oscillator strength of the transition $i - c$;

$f_1 = 0.665$, $f_2 = 0.71$, $f_3 = 0.81$, $f_4 = 0.95$, f_i ($i \geq 5$) = 1.

$u_i = \frac{E_i}{kT_e}$, E_i is the ionization energy of level i .

$\Psi_1(u_i)$ is a dimensionless quantity, tabulated in [11].

McWhirter calculated these rate coefficients with

$$C_{ic} = 3.24 \cdot 10^{-14} \frac{i^2}{(kT_e)^{\frac{1}{2}}} e^{-u_i} \quad (m^3 s^{-1}) \quad (I.16)$$

Both rate coefficients are, again, based on the same theoretical cross-section [10]. Drawin's expression for C_{ic} contains a factor $\Psi_1(u_i)$, resulting from the shape of the cross-section. In this case, the absolute value of the

cross-section is in accordance with the theoretical value. The factor e^{-u_i} appearing in McWhirter's formula is an approximation to $\Psi_1(u_i)$. For $u_i = 1$, the formulae (I.15) and (I.16) produce the same result (regarding the factor f_i). For $u_i = 0.02$, Drawin's value for C_{1c} is 9 times as high as McWhirter's, for $u_i = 2$ Drawin's formula gives a value that is twice as small as McWhirter's. Table I.2 gives numerical data for C_{ic} at an electron temperature of $2.56 \cdot 10^5$ K.

Table I.2 Rate coefficients C_{ic} (in $m^3 s^{-3}$) at $T_e = 2.56 \cdot 10^5$ K.

	C_{1c}	C_{2c}	C_{3c}	C_{4c}	C_{6c}	C_{8c}
McWhirter	$5.8 \cdot 10^{-16}$	$1.45 \cdot 10^{-14}$	$0.48 \cdot 10^{-13}$	$1.04 \cdot 10^{-13}$	$2.27 \cdot 10^{-13}$	$4.3 \cdot 10^{-13}$
Drawin	$1.67 \cdot 10^{-16}$	$1.92 \cdot 10^{-14}$	$1.17 \cdot 10^{-13}$	$4.0 \cdot 10^{-13}$	$1.32 \cdot 10^{-12}$	$2.9 \cdot 10^{-12}$
McW/Dr	3.5	0.77	0.41	0.26	0.17	0.14

5. Relation between the rate coefficients and the population densities

As we are mostly interested in the contribution of the He^+ ground state to the population densities of the excited levels we can restrict ourselves to the relation between the rate coefficients C_{ij} and C_{ic} and the collisional-radiative coefficients $r_i^{(1)}$ (eq. I.8). In the first place, the coefficients $r_i^{(1)}$ are directly proportional to the rate coefficients C_{1i} . If one multiplies the rate coefficients C_{1i} by a factor p , the population densities change by the same factor. This will be immediately clear for the lowest excited levels; these levels are mainly populated by direct excitation from the ground state (C_{1i}), and depopulated by spontaneous emission. Higher levels are populated by direct excitation from the ground level, depopulated by ionization (C_{ic}) and spontaneous emission, and (de)populated by (de)excitation to neighbouring levels, so that we have here a complex mechanism. But all the populating processes, whether direct from the ground state or via other excited states, start eventually from the ground state. Therefore, the rate of populating an excited level is directly proportional to the direct excitation processes.

As an illustration of the foregoing we show in table I.3 the C.R. coefficients $r_i^{(1)}$ calculated by McWhirter and by Drawin, for an electron temperature of $2.56 \cdot 10^5$ K and an electron density of $1.28 \cdot 10^{20} m^{-3}$.

Table I.3 $r_i^{(1)}$ for $T_e = 2.56 \cdot 10^5$ K and $n_e = 1.28 \cdot 10^{20} \text{ m}^{-3}$.

	i=2	3	4	5	6	7	15
McWhirter	4.4(-5)	1.9 (-5)	1.2(-5)	5.9(-6)	2.8 (-6)	1.3 (-6)	7.9(-9)
Drawin	2.0(-5)	0.63(-5)	0.3(-5)	1.2(-6)	0.47(-6)	0.21(-6)	3.5(-9)
ratio McW/Dr	2.1	3.0	4.0	4.9	6.0	6.2	2.3

The differences at $i=2$ and $i=3$ can be ascribed entirely to the differences between the coefficients C_{12} and C_{13} of the two authors (table I.1). The differences at the levels $i=5-15$ are the result of differences between the rate coefficients C_{1i} , C_{ij} , and C_{ic} . The first type of rate coefficient is responsible for a difference of a factor about 3; the other two types are responsible for an additional factor that depends on the level considered. The latter factor is partly occasioned by different values for the ionization rate coefficient C_{ic} ; this depopulating process is stronger in Drawin's case than in McWhirter's (table I.2), and therefore Drawin's population densities will be lower than McWhirter's. The influence of the rate coefficient C_{ij} for (de)excitation to neighbouring levels upon the population densities is not simply predictable. Processes of this kind can lead to populating, as well as depopulating an excited level, and the balance can be different for different levels.

In the foregoing it is shown that the C.R. coefficients $r_i^{(1)}$ show serious discrepancies. An experimental check seems to be useful, especially because these coefficients are of interest to those who want to deduce the electron temperature and density from the emitted line radiation. This method applies to astrophysical plasmas as well as to laboratory plasmas intended for fusion reactions. In the former case, He^+ ions are abundant, in the latter case a small quantity of He is added to the plasma.

1. Introduction

This type of discharge was frequently used for many different purposes. The investigations can be divided into two classes:

1. the study of the discharge in the hollow cathode [13, 14, 15];
2. the study of the arc. The following subjects may be mentioned:
 - i. The investigations of the plasma properties of the arc. For example: rotation [16, 17, 18, 19], waves and instabilities [20];
 - ii. The investigation of the thermal balance [21, 22];
 - iii. The investigation of the population densities of the excited states and the atomic processes involved [23, 24].

The majority of the experiments were done on Ar plasmas. Hydrogen is also often used, and incidentally He, Ne, Xe or molecular gases, and mixtures. Ar is often used because the HCD runs easily with it and Ar is relatively cheap. The same holds for H. Our HCD does not operate so easily with He as it does with Ar. Also, the discharge is less stable than it is with Ar. The plasma created by HCD has the important property that it is highly ionized (for more than 50 per cent) at intermediate electron densities (10^{19} - 10^{20} m⁻³). The HCD may operate stably for many hours.

2. Description of the HCD

a. Introduction

The main features are shown in fig. II.1. The pyrex discharge tube is 1.60 m long. It has a diameter of 30 cm at the ends, and of 25 cm in the middle. In the middle section 12 pyrex observation tubes have been melted to the main tube. The cathode is a hollow tungsten tube, the anode a tungsten disc. A constant He gas flow is fed through the cathode, and the gas is pumped off by two mercury diffusion pumps.

The discharge is confined to the tube axis by a magnetic field of about 0.1 T, produced by two magnetic coils. The discharge current is adjustable between 0 and 300 A. In the following we shall discuss some details.

fig. II.1 *Hollow cathode arc discharge.*

b. The cathode (fig. II.2)

The hollow cathode is a tungsten tube having an inner diameter of 5 mm, an outer diameter of 9 mm and a length of 8 cm. It is fitted to a molybdenum cylinder, which is held by a water-cooled stainless steel holder. The hol-

fig. II.2
Cathode

der is electrostatically screened by a tantalum shield at floating potential. The cathode is at ground potential. The whole construction is mounted on a flange of the vacuum vessel and is screened by a pyrex tube. The life-time of the cathode is about 150 hours at an arc current of 100 A.

c. The anode (fig. II.3)

The anode is a tungsten disc having a diameter of 7 cm and a thickness of 10 mm. It is mounted on a tungsten cylinder, which is also attached to a stainless steel holder identical to the cathode holder. There is no gas

fig. II.3
Anode

feed through the anode. The anode is "isolated" from its surroundings by a floating tantalum screen. The water-cooled holder is electrically connected to the two water pipes and the gas pipe. The three pipes are carried through a flange joined to the vacuum vessel. The construction of the three insulated feed-throughs is shown in fig. II.4.

fig. II.4
Anode feed-through

d. The pumping system (fig. II.1)

The basis of the pumping system is formed by two mercury diffusion pumps. Each diffusion pump has an unbaffled pumping speed of 2700 l s^{-1} . Above each diffusion pump is fitted a liquid-nitrogen cooled trap. Two butterfly valves can shut off the discharge vessel from the diffusion pumps. If the discharge is running on He, one valve is completely open, the other one only slightly. This is done because at the working pressure of the discharge, $2 \cdot 10^3$ torr, the pumping off is unstable; it means that a small change in the gas feed occasions a great change in the gas pressure in the vessel. By nearly closing one valve, the gas pressure above this pump is lowered, and the pressure in the vessel can be maintained stable. The diffusion pumps are backed by a two-stage rotary oil pump of 660 l/min. The ultimate vacuum is $1\text{--}3 \cdot 10^{-6}$ torr.

e. The gas inlet system (fig. II.5)

Ar gas or He gas, is each led to the hollow cathode via a valve. The Ar-valve is hand-operated, the He valve electrically. Two differential pressure meters are used for measuring the gas flow.

fig. II.5
Gas-supply system

The He gas is from l'Air liquide, purity N 45. The pressure in the discharge vessel is measured by an ionization gauge, which is electrically connected to a pressure control system that keeps the pressure on a constant, preset level by regulating the He valve. The pressure drift is $\pm 2\%$ at a pressure of $2 \cdot 10^{-3}$ torr. The ionization gauge has been calibrated to a McLeod manometer.

f. The observation tubes

There are 12 observation tubes, placed in the middle section of the pyrex discharge tube (fig. II.1). Every observation tube is open-ended. Sealing-off was done by o-rings. The observation tubes are 7 cm long, the inner diameter being 4 to 10 cm.

The set of tubes nearest the cathode is used for Thomson scattering and for line intensity measurements. The next set is used for a vacuum monochromator, a Langmuir probe, and for several short running experiments. The set next to the anode is used for a microwave interferometer and for line profile measurements.

g. Cooling

If the arc is operated at 100 A, the anode-cathode voltage is about 100 V, and a power of 10 kW is dissipated. The greater part of this power is disposed of by radiative emission of the anode and the cathode. This emission would heat the pyrex tube and the o-rings between the pyrex tube and the

stainless steel construction to the diffusion pumps to an intolerable value. Therefore, round each electrode a water cooled, winding brass tube was fitted (fig. II.1). The pyrex discharge tube was cooled by several air blowers. If necessary, constructions fitted to the observation tubes were water cooled as well.

h. Electrical systems

The power supply for the discharge is capable of delivering 300 A at 300 V. It is current-stabilized.

The power supply for the magnetic coils delivers at most 580 A at 110 V. It is current-stabilized, too.

i. The magnetic field

The shape of the magnetic field is shown in fig. II.6. In this figure the positions of the electrodes and the observation ports are indicated.

fig. II.6

Shape of the magnetic field along the discharge axis; *ca* indicates the position of the cathode; *c₁* and *c₂*, magnetic coils; *o₁*, *o₂* and *o₃*, observation ports; *a*, anode. The magnitude of the magnetic field is indicated by the value of the electrical current through the coils.

The magnetic induction on the discharge axis is at the utmost 0.2 T.

j. Starting procedure

The arc is started with Ar gas in the following way: A copper rod, mounted close to the cathode is connected to a transformer that is capable of delivering 3000 V at 0.5 A (50 Hz). This causes pre-ionization. The discharge current is almost immediately taken over by the anode; this small current heats the cathode, the cathode emits more electrons, the current increases, and so on. After about 30 seconds the cathode is at its working temperature (2500 - 3000 K).

Then the Ar flow rate is lowered, and compensated by an additional He flow, until there is He flow only.

3. Discharge characteristics

The discharge parameters that can be changed independently are:

- (a) The gas pressure p from $1.3 \cdot 10^{-3}$ to $3 \cdot 10^{-3}$ torr.
- (b) The magnetic field current I_B from 250 to 500 A.
- (c) The arc current I_{arc} from 30 to 200 A.

The anode-cathode voltage V_{arc} is a dependent variable. The lower limit of the pressure range is determined by a fast rise of V_{arc} . This limit is slightly dependent on I_{arc} . At the upper limit the pressure control system fails to keep the pressure at a constant level. This is caused by the pumping characteristic of the diffusion pumps. The magnetic field current could not be lowered to under 250 A because the discharge would then become unstable; the same holds for the lower and the upper limits of the discharge current. The physical reasons for these limits have not been investigated. We intended to change the independent discharge parameters so as to vary the electron temperature and the electron density of the plasma as much as possible. As an indication of these quantities we observed the line intensity of the He II line at 4686 Å. This line intensity showed only little dependence on the pressure; it was about directly proportional to the magnetic field strength, and increased almost quadratically with the arc current up to $I_{arc} = 100$ A. At larger arc currents this dependence levelled off. From these facts we concluded that it was suitable to change the arc current from 35 to 100 A only, while the pressure was fixed at $2 \cdot 10^{-3}$ torr, and the magnetic field current at 400 A. For these circumstances the anode-cathode voltage as a function of the arc current is shown in fig. II.7.

fig. II.7

Anode-cathode voltage V_{arc} as a function of the discharge current I_{arc} .
 $p = 2 \cdot 10^{-3}$ torr, magnetic field current 400 A.

4. Stability of the discharge

(a) After the arc current had been set at the desired value, it took about one hour until the anode-cathode voltage had reached a stable level and the irregularities in the gas pressure were below $\pm 2\%$. After this warming-up period the discharge ran stably for many hours.

Because the plasma was not completely reproducible from one day to another, we arranged our experiments in such a way that the measurements needed for each subject of investigation were completed in one day.

(b) The periphery of the plasma showed oscillations with a frequency of about 40 kHz. The oscillations were apparent in the line intensity of the He II line at $4686 \text{ \AA}$. Observing the centre of the arc (side-on), the modulation was 10% or less; more off-centre it might increase to 30%. All our experimental data were time-averaged over these oscillations. It will be clear that errors may be made in the interpretation of the data, if one considers them as really time-independent quantities. We, therefore, concentrated our measurements on the centre of the arc.

1. Introduction

The following table shows the quantities observed and the methods used.

	population density	n_e	T_e	$n(\text{He}^+)$	$n(\text{He}^{++})$	others
Line intensities	x				x	
Thomson scattering		x	x			
Microwave interferometer		x				
Langmuir probe			x	x	plus	x
Line profile						$T_{\text{He}}, T_{\text{He}^+},$ line splitting

The microwave interferometer and the Langmuir probe served as supplements to the Thomson scattering. The line profile measurements gave information on the velocity of excited He^+ particles, and on the multiplet structure of the He II emission line at $4686 \text{ \AA}$.

The measurements intended for our main purpose, viz. the experimental investigation of the C.R. model, were carried out by two diagnostic methods: line intensities and Thomson scattering. These two diagnostics, together with the line profile diagnostic will be discussed in this chapter. The microwave interferometer and the Langmuir probe will be treated in Appendix III.

2. Line intensitiesa. Theory

The radiant power emitted by a unit volume filled with n_q particles in the excited state q by spontaneous emission to a lower level p is (neglecting stimulated emission)

$$P_{qp} = n_q A_{qp} h\nu_{qp} \quad (\text{III.1})$$

where A_{qp} is Einstein's transition probability, h is Planck's constant and ν_{qp} is the frequency of the emitted radiation. If the absolute value of P_{qp} is measured, and A_{qp} is known, one can calculate n_q .

b. Experimental

The experimental arrangement is shown in fig. III.1. The plasma is imaged on the slit of a Jarell-Ash 0,5 m grating monochromator by means of a concave mirror (M_1). The image is turned 90° by 2 flat mirrors M_3 and M_4 so that the radiation from a horizontal slice of the plasma passes through the

fig. III.1

Optical arrangement of the line intensity equipment.

PL, plasma; D_3 , D_4 , diaphragms; W_2 fused silica window; R , position of the fused silica reflector, or Al mirror; M_1 to M_5 , mirrors; F , position of filters; L , fused silica lens; C , position of the ribbon lamp, or the carbon arc lamp.

Distances are not on scale.

slit of the monochromator. The magnification plasma-to-slit is 0.8 x. Both the entrance and exit slits of the monochromator are $50\text{ }\mu\text{m}$ wide. The en-

trance slit is 8 mm high, the exit slit 17 mm. The narrowest diaphragm is mirror M_1 (11.8 cm diameter). To obtain an achromatic image, only mirrors are used between the plasma and the entrance slit (except the fused-silica window W_2). The wavelength range is 2000-12000 Å. All mirrors are Al on glass, without an additional quartz coating. Mirror M_1 can be tilted along an axis parallel to the plasma-axis in order to scan the plasma image over the monochromator entrance-slit. The scanning range over the plasma is 3 cm. At position F optical filters can be placed in the light path in order to suppress stray light. Stray light could interfere in two ways:

- a) if a weak He II line close to a strong He I line was measured
- b) during the calibration of the system at UV wavelengths.

Two gratings were used ; one for 2000-6600 Å with 1180 grooves/mm, blazed at 3000 Å, and one for 6600-12000 Å with 590 grooves/mm and blazed at 7500 Å. Two photomultipliers were employed: one for the spectral range of 2000-8300 Å, and one for 8300-12000 Å. The first is an EMI 9659 QB, with spectral response S20, the second an EMI 9684 B, with spectral response S1. Both were cooled in order to lower the dark current; the S20 was cooled to -20°C , and had an anode dark current of $5 \cdot 10^{-11}$ A, the S1 was cooled to -100°C and had a dark current of $2 \cdot 10^{-11}$ A. Cooling was done with liquid nitrogen; photomultiplier housings and temperature regulation (S20 only) were of our own design. A movable mirror M_5 served for reflecting the light to the S_1 photomultiplier. Fused silica lenses (L) were used to fill the photocathodes correctly with light.

Emission lines were recorded by scanning the wavelength with the monochromator over the line and displaying the photomultiplier anode current on a recorder after amplification by a Keithly electrometer.

c. Calibration

We used a tungsten ribbon lamp as an absolute light source for the wavelength interval 2300-12000 Å [25]. The lamp was calibrated by the Fysisch Laboratorium, Rijks Universiteit of Utrecht.

This lamp, located at C (fig. III.1), and a reflector (R) were so positioned that the virtual position of the lamp coincided with the plasma (P). The reflector was either a wedge-shaped fused silica plate or an Al mirror. The front reflection of the fused silica plate was used, the reflection at the back was cut off. The reflection coefficient was calculated from the known index of refraction (prospectus issued by Halle).

At UV wavelengths the spectral radiance of the ribbon lamp is low; here, an Al mirror had to be used. The reflection coefficient of the mirror was measured by comparing it with the fused silica plate, using several gas discharge lamps (Hg, Cd, Zn), located at C, as light sources.

As mentioned in the preceding section, we used optical filters to suppress stray light. Stray light of the ribbon lamp becomes important at wavelengths lower than 3000 Å. Interference filters were used to suppress it. We checked this by scanning the wavelength of the monochromator outside the passband of each interference filter; no radiation was detected.

The following table shows the He II lines for which we measured the absolute intensity, the filters used and some more information about the calibration.

Table III.1 Calibration He II lines

line Å	temperature ribbon lamp K	reflector	filter	purpose filter	error %
2306	2700	Al	IF 235	straylight lamp	6
2385	2700	Al	IF 235	" "	6
2511	2700	Al	IF 256	" "	5
2733	2700	Al	IF 256	" "	5
3203	2700	fus. silica	IF 315	" "	3
4200	2200	"	IF 420	straylight He I	3
4339	2200	"	IF 436	" "	3
4542	2200	"	IF 447	" "	3
4686	2200	"	NG 3	strong line	3
4859	2200	"	-		3
5412	2200	"	-		3
6560	2200	"	-		3
6891	2200	"	RG 665	2 nd order stray-	3
7178	2200	"	RG 665	light lamp	3
7593	2200	"	RG 665	" "	3
8237	2200	"	RG 665	" "	3
9545	2200	"	RG 665	" "	3
10124	2200	"	RG 665	" "	4
11626	2200	"	RG 665	" "	5

IF 235 denotes "interference filter, peaked at $2350 \text{ \AA}$ ", and so on.

NG 3 is a gray filter, transmission about 1% (Schott).

RG 665 is a blocking filter, transmitting $\lambda > 6650 \text{ \AA}$ only.

The filters were also used for the corresponding line measurements.

The calibration was carried out more than once. The differences were never more than a few per cent. We also looked for systematic errors. The following possibilities were investigated:

(1) Influence of the slit height of the monochromator.

Using the ribbon lamp as light source, we varied the slit height in steps from 1 to 8 mm, at selected wavelengths. The observed signal appeared to be directly proportional to the slit height.

Secondly, absolute intensities of most of the He II lines of table III.1 were measured using a slit of 1 mm height and compared with the absolute intensities of the same lines measured with an 8 mm high slit. There were no significant differences.

(2) Influence of the slit width.

The absolute radiances of selected He II lines were independent of the slit widths $2 \times 50 \text{ }\mu\text{m}$ or $2 \times 100 \text{ }\mu\text{m}$.

(3) Linearity.

The photomultipliers were tested for linearity at different wavelengths. No deviation from linearity exceeding the experimental error of $\pm 2\%$ could be detected.

(4) Calibration source.

The tungsten ribbon lamp is the most precise commercially available radiation source. Its accuracy depends on the precision in determining its temperature and that in determining the emission coefficient of the ribbon [25]. The inaccuracy of the temperature measurement is certainly lower than 5 K, leading to an error in the radiance of 3% at $2300 \text{ \AA}$. The measurement of the emission coefficient is better than $\frac{1}{2}\%$. Therefore, the total error is less than 3%.

However, we thought it useful to employ a second calibration source; the graphite arc lamp. This lamp is reliable down to $2700 \text{ \AA}$, and has an inaccuracy of 10% [26, 27]. As the temperature of the emitting surface is about 3800 K, the spectral distribution of the radiation is quite different from that of the tungsten ribbon lamp, which has a temperature of at most 2700 K. We were now able to use the fused silica reflector at all wavelengths.

The graphite arc lamp was equipped with a Noris D, 7 mm $\varnothing$ cathode and a RW II, 6.3 mm $\varnothing$ anode, at right angles to each other, with automatic adjustment

of the two electrodes. The arc current was 7.3 A, the same value as used by Magdeburg [26] in comparable conditions. The voltage over the electrodes was 65 V, somewhat lower than in Magdeburg's case (70 V).

The spectral radiance was obtained from tables of Hattenburg [27] (2300 - 8237 Å) and of Magdeburg [26] (9345 - 11626 Å). These tables have an overlap between 2500 and 8500 Å and are in excellent agreement in it.

For wavelengths between 2500 and 11626 Å the deviations of the graphite arc lamp with respect to the ribbon lamp were in 3 cases ~10%, and for the other wavelengths in this region less than 6%. At 2385 Å the deviation was +20%, at 2306 Å +100%, but in both cases the graphite arc lamp was no longer reliable.

(5) Polarization.

Polarization, if any, of the He II lines is of importance to the calibration. The degree of polarization of He II lines between 4200 Å and 11600 Å was investigated with the help of a Glan-Taylor prism as polarizer. These lines showed no polarization. We leave the question why He II lines should be polarized undiscussed.

(6) Influence of the magnetic field.

The magnetic field of the HCD has some influence on the photomultipliers. Two mu-metal shields for each photomultiplier reduced the influence, but could not eliminate it. Therefore, all calibrations were done in the presence of the magnetic field.

The overall accuracy of the calibration was determined by adding quadratically errors originating from (a) the temperature of the ribbon, leading to errors of 0-3%, depending on the wavelength; (b) the photomultiplier anode current (1%-3%); (c) the coefficient of reflection of the Al reflector (4%); (d) repeated calibrations showed differences of 2%. The resulting error depends on the wavelength, and varies between 3% and 6% (table III.1).

The absolute line intensities have been determined as follows:

(1) the area S of the line profile is measured, the height in unit ampere, the breadth in μm .

(2) The radiance of a spectral line is

$$L_{\text{line}} = \frac{L_{\lambda}^{\text{lamp}}}{I^{\text{lamp}}} \cdot S \quad (\text{III.2})$$

where I^{lamp} is the photomultiplier current for the ribbon lamp. $L_{\lambda}^{\text{lamp}}$ is the

spectral radiance of the ribbon lamp. This quantity may be found in [25]; it is expressed there in units $\text{erg} \cdot \text{cm}^{-2} \cdot \mu\text{m}^{-1} \cdot \text{sr} \cdot \text{s}^{-1}$. If we express $L_{\text{lamp}}^{\text{line}}$ in units $\text{W} \cdot \text{m}^{-2} \cdot \mu\text{m} \cdot \text{sr}^{-1}$, S in $\text{\AA} \cdot \mu\text{m}$, $I_{\text{lamp}}^{\text{line}}$ in $\text{\AA}$, we obtain L^{line} in $\text{Wm}^{\frac{1}{2}} \text{sr}^{-1}$. This quantity is the radiance of a spectral line integrated along the line of sight.

We are mostly interested in the radiant intensity per unit volume, L_v^{line} , expressed in units $\text{Wm}^{-3} \text{sr}^{-1}$. $L_v^{\text{line}}(r)$ is calculated from $L^{\text{line}}(u)$ (fig. III.2) by Abel transformation. Therefore, $L(u)$ has to be measured at different

fig. III.2 Illustration of the quantities $L(u)$, the radiance integrated along the line of sight, and $L_v(r)$, the local radiant intensity per unit volume.

positions u .

The preceding procedure is only correct if there is no absorption of radiation by the plasma. It will be shown that this is indeed the case for the He II lines we measured (see Chap. IV.1.e). It may be remarked that only the He II lines mentioned were measured in absolute value. Other He II lines are too weak, or have a wavelength below $2300 \text{ \AA}$. We began measuring two He II lines in the vacuum UV region, at $1215 \text{ \AA}$ and $1639 \text{ \AA}$, with a McPherson vacuum monochromator of the Seya-Namioka type. The results of these measurements will not be reported in this work, because calibration encounters too much problems yet.

3. Thomson scattering

a. Introduction

Thomson scattering is nowadays a common technique in plasma diagnostics. There are many papers on the theory and on experimental arrangements. It is not our purpose to discuss these papers. For a general review of theory and experiment we refer to the contribution by Kunze in "Plasma Diagnostics", edited by Lochte-Holtgreven [28]. Here we give a description of our specific arrangement and experimental procedure.

Thomson scattering is scattering of radiation by free electrons. The cross-

section for the scattering can be calculated with the help of the classical theory on electro-magnetic waves and is called Thomson cross-section; it is independent of the wavelength, and has a magnitude of $0.665 \cdot 10^{-28} \text{ m}^2$. The spectrum of the scattered light is broadened with respect to that of the incident light as a consequence of the velocity of the electrons. Therefore, the scattered light gives information on the number of scattering electrons and their velocities.

One can distinguish two scattering modes, non-collective and collective. The spectra of these modes are quite different in shape. Whether one has to do with one mode or the other is determined by the parameter

$$\alpha = \frac{1}{k\lambda_D} \quad (\text{III.3})$$

where: $k = |\vec{k}_s - \vec{k}_o|$, $\vec{k}_o$ is the wave vector of the incident radiation,
 $\vec{k}_s$ that of the scattered radiation,
 λ_D is the Debye length.

If $\alpha \ll 1$, there is non-collective scattering, if $\alpha \gtrsim 1$, scattering is collective. In our case $\alpha = 0.03$.

The spectrum of the scattered radiation has, in the case of non-collective scattering of monochromatic radiation, a Gaussian profile if the electrons have a Maxwell-Boltzmann velocity distribution. The area of this profile is directly proportional to the electron density; the half-width is proportional to the square root of the electron temperature. Thomson scattering is one of the few diagnostic techniques that measure plasma quantities locally without disturbing the plasma (see later discussion).

b. Arrangement of components

The experimental arrangement is shown in fig. III.3.

The ruby laser is used in the free mode. It delivers a light pulse of 30 J in 1.5 ms. The beam diameter is 16 mm, and the beam divergence is 10 mrad (manufacturer's rating). The light is horizontally polarized. The wavelength is 6943 Å.

We used a repetition frequency of the laser between 2 shots per minute and 1 shot per two minutes.

The laser light is focussed on the plasma by two fused silica lenses L_1 and L_2 (fig. III.4); the focal spot has a diameter of 2 mm. An intermediate beam stop BS is placed in front of lens L_2 . Behind lens L_2 the light passes

fig. III.3

Optical arrangement of the Thomson scattering equipment. PL, plasma; RL, ruby laser; PM, photomultiplier; L_1 , L_2 , fused silica lenses; F_1 , position of the additional filters; HeNe, position of the HeNe laser; P, fused silica prism; BS, beam stop; W_1 , W_2 , fused silica windows; D_1 to D_4 , diaphragms; DIF, diffusor; PIN, PIN diode; L_3 to L_6 , achromatic lenses; S_1 , S_2 , rectangular diaphragms; F_2 , position of filters; IF, interference filters.

fig. III.4

Focussing of the laser beam in the plasma. For symbols see fig. III.3. The focal distances of lenses L_1 and L_2 are 200 cm and 25 cm respectively.

through a fused silica window W_1 and is then led through a set of diaphragms D_1 which screen off stray light coming from lens L_2 and window W_1 (fig. III.5). Having passed through the plasma the laser light enters a beam dump,

fig. III.5

Construction of the adjustable holder for lens L_2 , beam stop BS, window W_1 and diaphragm tube D_1 . T denotes the pyrex discharge tube.

which consists of a thick plate of black glass placed at the Brewster angle. The transmission of this glass plate is about 0.3%. A diaphragm D_2 screens off the stray light coming back from the glass plate. Behind the glass plate a diffusor and a PIN diode have been placed. The PIN diode detects the remaining laser light, in this way serving as a laser monitor.

The whole light path is surrounded by brass tubes painted black on the inside. This was done to prevent stray light from entering the detection system, and as a safety measure. The set of diaphragms D_1 and the diaphragm D_2 are made of chemically blackened aluminium.

The position of the ruby laser is adjustable. Lens L_2 , window W_1 and diaphragms D_1 are mounted in a holder that is also adjustable (fig. III.5). With the latter adjustment we were able to position the focal spot at the right place in the plasma.

Light scattered at right angles to the incident laser beam is focussed by

lens L_3 (fig. III.3) on a diaphragm S_1 via mirror M. The solid angle of detection is determined by the opening of lens L_3 and measures 0.023 sr. Diaphragm S_1 has a rectangular shape and is 8.0 mm wide and 7.0 mm high. The image of this diaphragm projected by L_3 back into the plasma is slightly wider than the focal spot of the laser beam, so the position of the intersection of the laser and the detection beam is not critical. A set of diaphragms D_3 prevents stray light from entering the detection system. A Rayleigh horn serves as a black background. The horn has double walls with a filling of a solution of CuSO_4 in H_2O that absorbs the laser light very efficiently. In front of the horn a diaphragm D_4 is placed, to prevent stray light from entering the horn.

Diaphragm S_1 is placed in the focal plane of lens L_4 ; lens L_5 focusses the light on a rectangular diaphragm S_2 that is slightly larger than the image of diaphragm S_1 . Between lenses L_4 and L_5 two interference filters are placed, for scanning the spectrum of the scattered light. Finally, diaphragm S_2 is imaged on the photocathode of a photomultiplier by lens L_6 . The spectral scanning and the light measurement will be discussed separately.

All lenses $L_3 - L_6$ are achromats; lens L_3 is adjustable vertically, so that the detected scattering volume can be positioned in the plasma. Lens L_3 and mirror M are mounted on motor-driven carriages, so that they can be moved out of the light path; in this way the observation port can also be used for line intensity measurements.

At location F_2 several optical filters were inserted into the light path, such as a polaroid filter (the scattered laser light is horizontally polarized), a red transmitting blocking filter RG 665, and, if necessary, other filters.

The whole equipment was screened off from stray light. From diaphragm D_3 to diaphragm S_1 this was done with black tubes; the equipment between S_1 and the photomultiplier is placed in a wooden box.

All optical components are placed on optical rails, except for diaphragms D_2 , D_3 and D_4 , window W_2 and the beam dump, which are mounted on the pyrex discharge vessel.

c. Lining out

A HeNe laser (Oriel) is placed on its permanent support (fig. III.3). This support is adjusted once for all. The prism P is placed so that the laser beam is directed through the middle of the entrance and exit observation

tubes. At about 1 m above the exit observation tube, where the beam hits a screen, a mark is made.

Then the HeNe laser is turned 180° , so that the beam is directed to the ruby laser. The front mirror of the ruby laser and lens L_1 are centered on the HeNe laser beam; next the HeNe laser is taken out of its support, and the beam of the ruby laser is adjusted so as to cover the mark. The HeNe laser is put back into its support, and lens L_2 is so placed in the beam that the latter remains centered around the mark. This can also be checked with the reflections on the HeNe laser head.

These optical components have now been aligned. The detection parts are treated in the same way as lens L_2 ; first, a mirror is placed on the axis of the pyrex discharge tube, and the HeNe beam is reflected horizontally through the middle of the observation tube for the scattered light. Then the optical components are placed on the rails one by one, keeping the direction of the laser beam unaltered, which is controlled with the help of two marks.

The whole alignment process is checked by substituting an electric lamp for the photomultiplier and so tracing back to the ruby laser. In this way diaphragms that may have been overlooked can be found, and this visual test is also a useful check on centring.

d. Spectrum scanning

The scanning of the Thomson spectrum is by tilting the interference filters, and the pass-band shifts from $6944 \text{ \AA}$ to lower wavelengths.

At first sight one filter should be sufficient for the purpose. It turned out, however, that then two He I lines, at $6678 \text{ \AA}$ and $7056 \text{ \AA}$, exceeded the Thomson signal by far because of the finite transmission in the wings of the filter (0.02%). With two filters these lines were suppressed sufficiently. The two filters (Baird Atomic) are identical; in particular, they show the same wavelength-shift if tilted. Therefore, they were mounted in one support and tilted together.

The relation between the wavelength shift $\Delta\lambda_s(\text{\AA})$ and the tilting angle $\theta(\text{deg})$ is

$$\Delta\lambda_s = 0.47 \theta^2 \quad (\text{III.4})$$

This relation was established when tuning the filters so that several spec-

tral lines of Ar were allowed to pass successively. The maximum wavelength shift we used was $110 \text{ \AA}$, which corresponds to $\theta = 15.3^\circ$.

The transmission characteristics of the filters have also been measured. This point will be discussed in Appendix I.1. At $\theta = 0^\circ$ the peak transmission is 19%, and the effective width of the transmission profile (defined as the width of a block-shaped filter having the same peak transmission as the original filter) is $6.3 \text{ \AA}$. The peak lies at $\lambda_{\text{max}} = 6944 \text{ \AA}$.

e. Light measuring

The scattered light is detected by means of an RCA photomultiplier, type C 31034, having a gallium-arsenide photocathode with a quantum efficiency of 12% at $6943 \text{ \AA}$. This photomultiplier is mounted in a Peltier-cooled housing. On the housing an SSR 1120 photon counter is mounted that amplifies, discriminates and shapes the pulses coming from the anode of the photomultiplier. The SSR photon counter has a dead time of 25 ns; the output pulses have a width of 10 ns and are counted. The counter is gated by a retardation unit triggered by the synchronization pulse of the control unit of the ruby laser in such a way that only during the flash of the ruby laser is the counter open. The gating time, 1.5 ms, is as short as practicable in order to reduce the number of plasma-light counts to a minimum.

The complete light measuring unit was checked for linearity in the range of 10 to 1500 counts per ms. Furthermore, we ascertained that the ruby laser triggering unit and the magnetic field of the HCD equipment caused no disturbance.

f. Laser monitor

A small portion of the laser light that passes through the plasma is detected by a PIN diode (see b. Arrangement of components). This diode is biased at -18 V, so that it acts almost as an insulator. Laser light falling onto the diode liberates electrons in the P-I-N-junction. The current is integrated by a Keithly electrometer used as a Coulomb meter. A small additional current is applied to the electrometer to compensate for the leakage current of the PIN diode (fig. III.6). The total electric charge passed through the PIN diode is a relative measure of the laser energy. The amount of electric charge is about 10^{-7} C . The laser monitor was checked for linearity by placing at F_1 (fig. III.3) several calibrated grey filters in the laser beam

fig. III.6
Laser monitor.

and observing the laser monitor and the Rayleigh scattering light. Both were found to be directly proportional to the transmission of the filters.

8. Calibration

Calibration was with the help of Rayleigh scattering as follows: The discharge vessel was filled with 10-15 torr Ar gas, and the scattered signal was measured. Comparison of this signal to that from Thomson scattering yields the electron density, because the cross-sections of Thomson scattering and Rayleigh scattering are known. The latter quantity has been subjected to some doubt in the literature; therefore, we measured this cross-section separately, which experiment is discussed in Appendix I.3. The electron density is calculated with the formula

$$n_e = \frac{S_T(0)}{S_R(0)} \cdot \frac{\sigma_R}{\sigma_T} \cdot n_R \cdot 1.06 \frac{\Delta\lambda_T}{\Delta\lambda_W} \quad (\text{III.5})$$

where: $S_T(0)$ Thomson signal at 6943 Å
 $S_R(0)$ Rayleigh signal at 6943 Å
 σ_R Rayleigh cross-section
 σ_T Thomson cross-section
 n_R density of the atoms (Rayleigh scattering)
 $\Delta\lambda_W$ effective width of the interference filters at 6943 Å
 $\Delta\lambda_T$ half width of the Thomson spectrum.

The electron temperature is determined by $\Delta\lambda_T$ alone

$$T_e = 11.1 \Delta\lambda_T^2 \quad (\text{III.6})$$

T_e in K and $\Delta\lambda_T$ in Å. The derivation of these formulae is to be found in Appendix I.2.

h. Stray light

The observed scattered light is a very small fraction, about 10^{-15} , of the incident laser light. The laser light has to pass through the entrance window and is absorbed in a beam dump. The entrance window scatters some of the laser light and the beam dump is not absorbent to the full 100%. The stray light can enter the detection apparatus via one or more secondary scatterings, and may be many orders of magnitude greater than the light scattered by the plasma. To reduce the stray light itself and avoid its entering the detection apparatus, we placed the diaphragms D_1 , D_2 , D_3 and D_4 in the observation ports (fig. III.3 and III.5). The positions and dimensions of these diaphragms have been chosen so that the stray light was reduced to an acceptable level, while the diameters of the diaphragms were large enough to permit adjustment of the light beams.

The remaining stray light gave a counting of at most 200 photo-electrons per laser shot, the interference filters being tuned to $6943 \text{ \AA}$. We investigated the origin of this light and found it to emanate from the Rayleigh horn; the horn was "illuminated" by light coming back from the laser beam dump.

This stray light did not interfere with the Thomson scattering measurements, because the minimum wavelength shift we adjusted was $13 \text{ \AA}$, in which case the stray light did not pass through the interference filters.

i. Plasma radiation

Plasma radiation can be continuum radiation or line radiation. The reception of the latter could be avoided by tuning the interference filters out of the lines. This is, of course, not possible for continuum radiation. Some suppression of this radiation can be obtained by the use of a polaroid filter because the plasma radiation is unpolarized, whereas the Thomson scattered radiation is plane-polarized. The remaining continuum radiation gave a counting of about 100 per laser shot. This is somewhat more than we expected:

At $6943 \text{ \AA}$, free-free radiation exceeds the free-bound radiation by a factor 10; the free-free radiation is [29]

$$\frac{dP_{ff}}{d\nu} = \sum_Z C_n n_e \cdot n(Z) \cdot Z^2 \left(\frac{X_H}{kT_e} \right)^{\frac{1}{2}} g_{ff} e^{-\frac{h\nu}{kT_e}} \quad (\text{III.7})$$

where: P_{ff} is the radiant power per unit volume of the continuum radiation, in $\text{Wm}^{-3} \text{ s}$.

$n(Z)$ is the particle density of the ions with charge Z

χ_H is the ionization energy of hydrogen

$\bar{g}_{ff}$ is a correction factor, equal to about 1

$C = 1.7 \cdot 10^{-53} \text{ Wm}^3 \text{ s}$.

On the axis of the discharge, $n_e = 3.3 \cdot 10^{19} \text{ m}^{-3}$ and $kT_e = 20 \text{ eV}$; for He^+ , $Z = 1$, and $n(\text{He}^+) = 2.6 \cdot 10^{19} \text{ m}^{-3}$; for He^{++} , $Z = 2$, and $n(\text{He}^{++}) = 0.3 \cdot 10^{19} \text{ m}^{-3}$. Further, $h\nu = 1.8 \text{ eV}$. We obtain

$$\frac{dP_{ff}}{d\nu} = 1.5 \cdot 10^{-14} \text{ Wm}^{-3} \text{ s}. \quad (\text{III.8})$$

With the help of the counting of the continuum radiation we found, after calibrating the detection apparatus with a tungsten ribbon lamp

$$\frac{dP_{ff}}{d\nu} = 3 \cdot 10^{-14} \text{ Wm}^{-3} \text{ s}. \quad (\text{III.9})$$

It is not excluded that the difference between the observed and the measured continuum radiation is due to experimental errors. The excess of continuum radiation cannot be explained by free-bound radiation from highly ionized impurities that have an estimated concentration of less than 1 to 1000; this concentration is too low to create the additional continuum radiation.

Another possible source might be the process: $\text{He} + \text{He}^+ \rightarrow \text{He}_2^+ + h\nu$.

Using the formulae for the case of H [30], we find that this radiation is less than the f-f radiation.

j. Perturbation of the plasma by the laser pulse

In the first place, heating of the plasma by inverse bremsstrahlung may occur. A rough estimate shows that under our conditions (a 30 J laser pulse, $n_e \approx 3 \cdot 10^{19} \text{ m}^{-3}$, $T_e \approx 2 \cdot 10^5 \text{ K}$, focussed laser beam area 3 mm^2) the heating is completely negligible [31].

Secondly, additional ionization of the plasma due to photo-ionization of excited atoms could happen. Such a phenomenon was reported by Lidsky, Rose and Thomson [32].

Our experimental conditions resemble theirs as regards both the plasma and the laser. They observed an additional signal, somewhat delayed in time with respect to the Thomson signal, and of comparable magnitude, but unpolarized. They ascribed this signal to an increase in the recombination radiation. We looked for such a signal by turning the polaroid filter so that it blocked the Thomson scattered signal, and by changing the gating time of the counter. But we did not find any additional light signal. The amount of photo-ionization is calculated as follows: the cross-section for photo-ionization of a hydrogen-like ion in the state i by a photon of frequency ν is given by [33]

$$\sigma_{pi} = 2.815 \cdot 10^{25} \frac{Z^4}{\nu^3 i^5} \text{ (m}^2\text{)}. \quad (\text{III.10})$$

For He^+ , $Z = 2$ and $i \geq 6$, and for He, $Z = 1$ and $i \geq 3$ in the case of the ruby laser (photon energy 1.8 eV). Because of the high power of i , only He atoms need to be considered. We assumed that the cross-section of H is a good estimate for that of He.

The number of photo-ionization processes per unit volume, and per unit time is

$$\left(\frac{dn_e}{dt} \right)_p = J_p \cdot \sum_{i=3}^{\infty} n_i \sigma_{pi} \quad (\text{III.11})$$

where J_p is the photon flux during the laser pulse and n_i is the population density of the i^{th} excited He level. The n_i 's are determined experimentally (Appendix II). The summation in (III.11) is so rapidly convergent that the first term is a good estimate of the total sum. With $n_3 = 10^{15} \text{ m}^{-3}$, and $\sigma_{p3} = 1.4 \cdot 10^{-21} \text{ m}^2$ (formula (III.10)) we find $\sum_{i=3}^{\infty} n_i \sigma_{pi} = 1.5 \cdot 10^{-6} \text{ m}^{-1}$ on the axis of the discharge.

The photon flux for a laser pulse of 30 J during 1.5 ms through an area of 3 mm^2 is $J_p = 2 \cdot 10^{28} \text{ m}^{-2} \text{ s}^{-1}$. The resulting photo-ionization rate is

$$\left(\frac{dn_e}{dt} \right)_p = 3 \cdot 10^{22} \text{ m}^{-3} \text{ s}^{-1}. \quad (\text{III.12})$$

We compare this with the production of electrons due to electronic ionization of He. At an electron density of $3 \cdot 10^{19} \text{ m}^{-3}$ and an electron temperature of $2 \cdot 10^5 \text{ K}$ the collisional-radiative ionization coefficient S is about

$10^{-14} \text{ m}^3 \text{ s}^{-1}$ [34]. The density of the He atoms is $2 \cdot 10^{19} \text{ m}^{-3}$ (Appendix II). Thus, the electronic ionization rate due to electronic collisions is

$$\left(\frac{dn_e}{dt} \right)_e = 2 \cdot 10^{19} \cdot 3 \cdot 10^{19} \cdot 10^{-14} = 6 \cdot 10^{24} \text{ m}^{-3} \text{ s}^{-1} \quad (\text{III.13})$$

So the ionization due to inelastic collisions is about 200 times as large as the photo-ionization.

The same way of reasoning applied to the parameters of the experiment of Lidsky et al. leads to a ratio of electronic to photo-ionization of 40. Therefore, it is not clear to us why Lidsky et al. ascribed the observed additional light to photo-ionization, unless their Thomson scattered signal was 40 times smaller than their continuum signal. In any case, we may neglect photo-ionization in our conditions.

k. Experimental procedure

Before taking any scattering data, the exit window for the scattered light (W_2 , fig. III.3) was cleaned. In the course of time everything inside the discharge vessel was "coated" by, probably, tungsten that evaporated from the electrodes. The coating absorbed light, especially in the UV region (the light intensity measurements were also carried out through window W_2 !). After one day running of the discharge there was no measurable absorption, so we felt every-day cleansing to be sufficient. The entrance window of the laser beam (W_1 , fig. III.3) was cleaned by the laser beam itself; it evaporated the layer from the glass. The efficiency of the beam dumps was not affected.

Next, the discharge tube was filled with 10-15 torr Ar; the Rayleigh signal was measured, and the intersection of the laser beam and the detection beam was controlled by sweeping the laser beam across the detection beam with the help of the adjustable holder of lens L_2 .

Concerning the Thomson scattering measurements, mostly 5 different positions (wavelengths) of the interference filters were chosen. The maximum wavelength shift was $110 \text{ \AA}$. The choice of these positions depended on the expected electron temperature, and on the spectral distribution of the plasma radiation. At every wavelength 20 laser shots were taken, with a repetition frequency of 1 per 2 minutes or 2 per minute, depending on the kind of measuring. Be-

fore and after each series of laser shots, the plasma light was measured. In total, one Thomson spectrum required 100 shots. If desirable, we scanned the scattering volume by adjusting lens L_3 . The scanning range was 15 mm.

After the Thomson scattering measurements were finished, we filled the discharge tube again with 10-15 torr Ar, and the calibration measurements were carried out.

Afterwards, the Thomson scattering signals were normalized with respect to the laser monitor and we applied corrections for the transmission of the interference filters, which is dependent on the wavelength (Appendix I.1.), and for the plasma light. By the method of the least squares the Gaussian profile was calculated, which gave the electron temperature, and, together with the results of calibration, the electron density, including their errors.

1. Some data on the Thomson scattering experiment

The counting of Thomson scattering is about 300 per shot, or less.

The counting of the plasma light is about 100 per shot.

The counting of stray light may be 200 per shot.

The counting of Rayleigh scattering on Ar
at 0.01 torr is 200 per shot.

4. Line broadening

Line profile measurements were performed with the aid of a Fabry-Pérot interferometer. The wavelength was scanned by varying the pressure of the gas atmosphere (Ar) in which the instrument was placed. Before entering the interferometer, the light was dispersed by a prism monochromator.

The width of the instrument's profile ($\approx 25 \text{ m}\text{\AA}$) was determined with the help of line profiles of spectral lines emitted by a gas discharge lamp. The line profiles themselves (width $\approx 10 \text{ m}\text{\AA}$) were measured with the same interferometer, having a plate distance $7\frac{1}{2}$ times as large as "normal". The width of the instrument's profile in this condition was supposed to be negligible.

The profile of a He II line can in practice be broadened by:

- (1) multiplet splitting and fine structure, about $200 \text{ m}\text{\AA}$, or 10 cm^{-1} , for the $4686 \text{ \AA}$ line;

- (2) Zeeman splitting, about 20 mÅ;
- (3) Doppler broadening, about 200 mÅ.

Other broadening mechanisms, such as Stark effect and pressure broadening may be neglected. The number of Zeeman components can be reduced experimentally by the use of a polarizer transmitting only the components that are polarized parallel to the magnetic field (π components). The resulting Zeeman splitting is of the order of 10 mÅ and is in practice negligible.

We separated the remaining phenomena (1) and (3) by assuming that all multiplet components have the same Doppler broadening, and by adjusting the value of this Doppler broadening and the distribution of the intensities of the components in such a way that the calculated line profile was alike the measured profile.

From the value of the Doppler broadening the ion temperature was calculated with $\Delta\lambda_{\frac{1}{2}} = 2\ln 2 (2kT_1/Mc^2)^{\frac{1}{2}}$, where $\Delta\lambda_{\frac{1}{2}}$ is the Doppler half width and the other symbols have their usual meanings.

The temperature of the neutral He atoms is obtained from He I lines. In this case multiplet splitting does not interfere, because it is much larger than the Doppler broadening.

The lines that were investigated are 4686 and 5412 Å (He II), and 4713 Å (He I).

As already mentioned in Chapter I.3, several assumptions were made to obtain solutions of the collisional-radiative (C.R.) equations that are applicable to a not too limited set of plasmas. In the present chapter we shall discuss these conditions in connection with our plasma.

1. Other atomic processes

We shall discuss some examples of other possible atomic processes that might be of importance.

The competing process is $\text{He}_i^+ + e \rightleftharpoons \text{He}_j^+ + e$. The cross-sections for the electronic processes are roughly a factor m_H/m_e greater than those for the atom-atom processes. Furthermore the electrons are over a factor $(m_H/m_e)^{1/2}$ faster than the atomic particles. Therefore the rate coefficients for the electronic processes are over a factor 10^4 greater than those for the atom-atom processes. In our case, the electron density and the density of the He atoms or of the He^+ ions are of comparable magnitude. Thus the atom-atom processes can be neglected.

Anderson et al. [35] reported experimental and theoretical cross-sections for $i = 5, 6$ and 7 . A calculation of the cross-sections for $i = 2, 3$ and 4 was reported by Gillespie [36]. These cross-sections are about 100 times less than those for the processes $\text{He}^+ + e \rightleftharpoons \text{He}_i^+ + e$. The densities of the He^+ ions and the He atoms are of comparable magnitude. Therefore, this process is of no importance.

This process may be compared with the process $\text{He}_i^+ + e \rightleftharpoons \text{He}^{++} + 2e$. The density of the He_j^+ particles is about 10^{-6} times the electron density. Moreover, the electrons have the benefit of their high velocity. Therefore, this process is also of no importance.

(d) Simular arguments can be applied to the following processes:

(e) Radiative absorption

The cross-section for absorption of radiation is dependent on the spectrum of the emission coefficient and on the absorption coefficient. In our experiment, Doppler broadening is the dominant broadening mechanism. Assuming equal absorption and emission profiles, the absorption cross-section for an emission line $i \rightarrow j$ is [37,38]

$$\langle \sigma_{ji} \rangle = \frac{\pi^{\frac{1}{2}} e^2}{2m_e c} \left(\frac{M}{2RT_M} \right)^{\frac{1}{2}} \lambda_{ij} f_{ij} \quad (\text{IV.1})$$

where M is the mass number ($\text{He}^+ : M = 4$)

T_M is the temperature of the atoms (or ions)

f_{ij} is the absorption oscillator strength.

The absorption of the emission line is

$$1 - e^{-n_j \langle \sigma_{ji} \rangle l} \quad (\text{IV.2})$$

where n_j is the particle density of the atoms in state j and l the effective thickness of the absorbing layer.

In our case, only absorption of resonance lines may be considerable.

Therefore $j = 1$; putting $l = 10^{-2}$ m, $T_M = 10^4$ K (obtained with the Fabry-Pérot interferometer) and $n_1 = 3 \cdot 10^{19} \text{ m}^{-3}$ one has:

$\text{He}^+, \quad 2 \rightarrow 1$ shows 32% absorption ($\text{Ly}\alpha$, 304 Å)

$3 \rightarrow 1$ shows 6% absorption ($\text{Ly}\beta$, 256 Å)

$4 \rightarrow 1$ shows 2% absorption ($\text{Ly}\gamma$, 243 Å).

A C.R. calculation, carried out by Drawin and Emard, in which these absorption rates were included, showed that the population densities of the levels 4 and higher are not altered by more than one per cent. as compared with those for an optically thin plasma.

We conclude that the C.R. calculations for optically thin plasmas are applicable to our case.

2. Divergence of the particle fluxes

This contribution appears in the C.R. equation (I.4) as the term $\vec{\nabla} \cdot (n_i \vec{v}_i)$, where $\vec{v}_i$ is the mean macroscopic velocity of the He_i^+ particles. In order to trace the importance of this contribution we consider the distance that a He_i^+ ion can move away during its life-time τ_i , in comparison with the dimensions of the plasma. This distance is $\tau_i \cdot v_i^{\text{th}}$, where v_i^{th} is the mean velocity of the He_i^+ ion.

This velocity has been calculated from the Doppler broadening of the He II line at $4686 \text{ \AA}$ ($4 \rightarrow 3$). One can safely assume that the temperatures of all He_i^+ ions are the same. We found $v_i^{\text{th}} \approx 10^4 \text{ ms}^{-1}$.

The next thing to do is to estimate τ_i . One cannot use the calculations of McWhirter [6] or Drawin [38], concerning relaxation times for establishing (quasi) steady-state populations. These calculations apply to the case that only the population of a level i is perturbed by a small amount, while the other levels are not influenced. The relaxation times thus obtained do not apply to the case that two or more levels are perturbed at the same time. We calculated the τ_i 's as follows:

The levels $i = 4$ and $i = 5$ (lower levels have no interest for us) are, to a great extent, depopulated to lower states by spontaneous emission, and populated from the He_1^+ ground state. The life-times of these states can be put equal to $(\sum_{j \neq i} A_{ji})^{-1}$. The levels $i \geq 6$ show an increasing coupling with neighbouring levels by inelastic electronic collisions. Their electronic coupling to the levels $j \leq 5$ is, however weak. Therefore a reasonable estimate of their life-times is given by $(\sum_{j \leq 5} A_{ji})^{-1}$.

The life-times thus calculated are given in table IV.1.

Table IV.1 Life-times of excited states

i	τ_i (s)	Δx_i (mm)
4	$2.07 \cdot 10^{-9}$	0.0207
5	$5.4 \cdot 10^{-9}$	0.054
6	$1.20 \cdot 10^{-8}$	0.120
7	$2.90 \cdot 10^{-8}$	0.290
8	$5.9 \cdot 10^{-8}$	0.590
9	$1.10 \cdot 10^{-7}$	—
10	$1.90 \cdot 10^{-7}$	—
11	$3.1 \cdot 10^{-7}$	—

The values for τ_i given in the table for the levels $i > 8$ are too high; since because of the strong coupling between these levels their effective life-time is the smallest among the values given.

Table IV.1 also shows the distance that an excited He^+ ion can run away during its life-time, Δx_i . An estimate of the transverse dimensions of the plasma is the halfwidth of the intensity profile of a He II emission line, i.e. about 10 mm.

Starting from the quantities given above, we conclude that the term $\vec{V}(n_i, \vec{v}_i)$ is indeed negligible.

3. Velocity distribution of the electrons

The rate coefficients needed for the calculation of the C.R. coefficients are obtained by integration of the appropriate cross-sections over the velocity distribution of the electrons, viz.

$$C_{ij} = \langle q_{ij} v_e \rangle = \int_0^{\infty} q_{ij}(v_e) f(v_e) v_e dv_e \quad (\text{IV.3})$$

where: $q_{ij}(v_e)$ is the cross-section and $f(v_e)$ is the velocity distribution function of the electrons.

It was assumed that $f(v_e)$ is the Maxwell-Boltzmann (M.B.) velocity distribution function. We shall discuss this assumption.

The energy relaxation time for the electrons τ_E can be calculated from formulae of Spitzer [39]. At an electron temperature of $2 \cdot 10^5$ K and an electron density of $3 \cdot 10^{19} \text{ m}^{-3}$, $\tau_E = 5 \cdot 10^{-8} \text{ s}$. The thermal velocity of the electrons is $3 \cdot 10^6 \text{ ms}^{-1}$. It may be expected that close to and inside the hollow cathode the velocity distribution is not Maxwellian. From the figures given above one concludes that at 15 cm distance from the cathode Maxwellian velocity distribution occurs. The plasma is investigated at 60 cm from the cathode.

We can measure the transverse velocity distribution experimentally with the help of Thomson scattering. On the axis of the discharge, where the electron "temperature" is about 20 eV, the results of these measurements do not show deviations from the M.B. distribution that exceed the experimental inaccuracy (10%) up to an electron energy of 60 eV (fig.V.1). If at energies greater than 60 eV the tail of the distribution function is 10% overpopulated, the rate coefficients do not change more than 10% either. Errors like these do not affect our final conclusions.

4. Population distribution of the sublevels

The values of the transition probabilities A_{ij} used in the C.R. calculations are based on the assumption that the sublevels of each excited main level are populated according to their statistical weights (Chapter I). The energy differences between the sublevels are small-of the order of 10^{-4} eV. These differences lead to a splitting of the spectral lines of the order of 200 mÅ. This is less than the resolving power of the Jarrell-Ash monochromator equipped with 50 μ m slits, but greater than the resolving power of the Fabry-Pérot interferometer.

We measured the so-called "branching ratio" of pairs of spectral lines coming from the same upper level with the help of the Jarrell-Ash monochromator, and we investigated the line profile of the emission line 4686 Å with the Fabry-Pérot interferometer.

a. Branching ratio

The ratio of the intensities of two spectral lines starting from the same upper level, but ending on different lower levels, is called a branching ratio, and can be written as

$$\frac{L_{ij}}{L_{ik}} = \frac{A_{ij}}{A_{ik}} \cdot \frac{\nu_{ij}}{\nu_{ik}} \quad (\text{IV.4})$$

where i denotes the upper level, j and k are lower levels. The transition probabilities A_{ij} and A_{ik} are averaged over the population distribution of the sublevels. Transition probabilities of various sublevel-sublevel transitions of He II were calculated by Green et al. [40]. Assuming a statistical population distribution, branching ratios for several pairs of He II lines were calculated; we compared these with the measured branching ratios. The measurements were done for upper levels $i = 5$ to 11, and for three values of the arc current (50 - 70 - 100 A). The line intensities were measured on the axis of the discharge and at 20 mm from the axis; afterwards, the off-axis line intensities, in the foot of the intensity profile, were subtracted from the on-axis intensities (fig. IV.1). In Chapter V we will show that the excited levels that are responsible for the foot of the intensity profile are populated by He^{++} , while the superposed peak is populated by He^+ (direct or indirect excitation). We are primarily interested in the latter contribution, as will be clear from Chapter V. Consequently we start with considering the population distribution of

fig. IV.1

Radiance profiles $L(u)$ of several He II lines. u indicates the position of height of observation.

sublevels that are populated directly or step-wise from the He^+ ground state. For each discharge current at least two measurements were carried out. No systematical differences between branching ratios for different discharge currents were detected, and all the data on each branching ratio were averaged. The results are shown in table IV.2. The observations for upper levels $i = 10$ and $i = 11$ are less reliable; the line 4339 was somewhat masked by two other lines, and the line 7178 was very weak; this weakness cannot, however, explain the great discrepancy between the experimental and the calculated value of the branching ratio for $i = 11$. We could not obtain data for upper levels lower than 5, because suitable pairs of He II lines were lacking (UV region!).

Table IV.2 Branching ratios $\frac{L_{i,j}}{L_{i,j+1}}$. Population from the He^+ ground state

line (Å)	upper level i	lower level j	branching ratio	
			experimental	calculated
3203	5	3		
10124		4	1.8 ± 0.2	2.57
2733	6	3		
6560		4	1.6 ± 0.2	2.42
2511	7	3		
5412		4	1.7 ± 0.3	2.40
11626		5	2.0 ± 0.3	2.01
2385	8	3		
4859		4	1.5 ± 0.3	2.37
9345		5	2.1 ± 0.3	1.97
2306	9	3		
4542		4	1.7 ± 0.3	2.35
8237		5	1.6 ± 0.2	1.96
4339	10	4		
7593		5	1.5 ± 0.4	1.95
4200	11	4		
7178		5	4 ± 1	1.94

From table IV.2 it is seen that the experimental branching ratios of the type L_{13}/L_{14} are systematically lower than the calculated ones. Two experimental branching ratios, both of the type L_{14}/L_{15} , agree with the calculated ones.

In a second experiment, we scanned the line intensities over the plasma column width. The measurements of all the lines mentioned in table IV.2, except lines 11626 and 9345 Å were carried out at an arc current of 70 A. The results for the upper levels $i = 5, 7, 9$ and 11 are shown in fig. IV.1. The intensity profiles were normalized to the same peak intensity; consequently if two intensity profiles with the same upper level are not coincident in the foot, it means that the population distribution of the sublevels in the foot is different from the one in the peak. The two intensity profiles of the upper levels 5 and 7 coincide. The same holds for upper levels 6 and 8. The intensity profile of the transition $9 \rightarrow 4$ coincides with that of $9 \rightarrow 5$, but the intensity profile of the transition $9 \rightarrow 3$ has a smaller foot

intensity than the preceding ones.

The picture is the same for upper level 10. For upper level 11 the intensity profile of the transition 11→4 is no longer coincident with that of 11→5. We shall discuss these results at the end of this section.

b. Line profile

Only multiplet components of the He II 4686 Å line could be resolved. Table IV.3 contains the thirteen components together with their relative intensities (assuming statistical population distribution) [41] .

Table IV.3 Multiplet components He II 4686 Å line (4→3)

line number	transition 3 — 4	(theory)	
		relative position (Å)	relative intensity (stat. population)
1	$2P_{1/2} - 2D_{3/2}$	-0.427	21.27
2	$2S_{1/2} - 2P_{3/2}$	-0.397	11.12
3	$2P_{1/2} - 2S_{1/2}$	-0.280	1.11
4	$2S_{1/2} - 2P_{1/2}$	-0.236	5.56
5	$2D_{3/2} - 2F_{5/2}$	-0.100	70.00
6	$2P_{3/2} - 2D_{5/2}$	-0.100	38.28
7	$2D_{3/2} - 2P_{3/2}$	-0.047	0.13
8	$2P_{3/2} - 2D_{3/2}$	-0.046	4.25
9	$2D_{5/2} - 2F_{7/2}$	0	100.00
10	$2D_{5/2} - 2F_{5/2}$	+0.027	5.00
11	$2D_{5/2} - 2P_{3/2}$	+0.080	0.63
12	$2P_{3/2} - 2S_{1/2}$	+0.102	2.22
13	$2D_{3/2} - 2P_{1/2}$	+0.114	1.13

The observed line profile for 10 mm out of the axis of the discharge is shown in fig. IV.2. (When we observe the centre of the discharge, the structure of the line profile disappears because of an increase in the Doppler broadening). The calculated line profile is shown, too. It was assumed that every multiplet components had its statistical intensity, except for component No. 9 ($F_{7/2} \rightarrow D_{5/2}$), which was diminished by 5% in order to fit the calculated curve to the observed one.

At 30 mm out of the discharge axis, component No. 9 was diminished by 17%. Going further off-axis, component 9 disappeared.

fig. IV.2

Line profile of He II, 4686 Å.

Solid line, calculated line profile.

Dotted line, measured line profile.

Observation 10mm out of the discharge axis; component No. 9 is diminished by 5% with respect to its statistical intensity.

The lower figure shows the positions and the statistical intensities of the components considered.

c. Discussion

The experimental branching ratios of table IV.2 are, leaving upper level 11 out of discussion, in seven of the nine cases smaller than the calculated ones; in two cases there is agreement. We shall deal with two possible reasons for the deviations, viz. a non-statistical population distribution, and an error in the measurement of the line intensities. We emphasize that at the moment we are only considering that part of the population densities that is contributed by the He^+ ground state. First, a non-statistical population distribution of the sublevels of the upper levels may be the cause of the deviating branching ratios. A considerable part of the radiative emission in a spectral line is due to the highest contributing sublevel. In case of statistical population distribution of the sublevels, 45% of the line emission $6 \rightarrow 4$ is due to $6g \rightarrow 4f$, and 54% of the emission $6 \rightarrow 3$ is due to $6f \rightarrow 3d$. If sublevel $6g$ is 60% overpopulated with respect to its statistical population, and sublevel $6f$ and lower are statistically populated, the branching ratio $L_{6,3}/L_{6,4}$ will have the observed value. All the experimental branching ratios of the $L_{i,3}/L_{i,4}$ type can be explained in this way. The two deviating branching ratios of the $L_{i,4}/L_{i,5}$ type can be explained similarly: in the latter cases the sublevels h should be over-populated. In order to link the two types of deviating branching ratios, one must assume that the higher the angular quantum number of the sublevel the more over-populated it is. Two observations are in contradiction with the foregoing assumption. In the first place, two experimental branching ratios of the $L_{i,4}/L_{i,5}$ type are in agreement with the calculated ones. Secondly, the resolved components of the $4686 \text{ \AA}$ line ($4 \rightarrow 3$) show an intensity distribution that is in disagreement with the supposition that the highest sublevels are over-populated; on the contrary, looking at the centre of the arc, the observed intensity distribution is almost in agreement with statistical population distribution. The situation that the sublevels of main level 4 do have statistical population distribution, while those of the higher levels do not, is very improbable. Supposing that there is some process or other -for instance inelastic electronic collisions- coupling the sublevels, and that it is intense enough to cause statistical population distribution of the sublevels of the main level 4, it will couple the sublevels of the higher main levels even more intensely, because the energy distances between sublevels decrease with increasing main quantum number.

A quantitative treatment of the processes between the sublevels of main levels 2 and 3 may be found in [42]. It is shown there that at electron densities of over 10^{19} m^{-3} the sublevels of main levels 3 and higher should be statistically populated.

Another cause of the deviations of 30 to 40% for 7 of the experimental branching ratios, might be that we made an error in the calibration procedure, so that the He II lines having wavelengths of 4686 Å and lower are in reality 40% more intense than we observed. (From Chapter V it follows that we may include the 4686 line ($4 \rightarrow 3$)). This assumption would explain the five deviating branching ratios of the $L_{i,3}/L_{i,4}$ type, the two of the $L_{i,4}/L_{i,5}$ type (4542 Å and 4339 Å), and would not affect the two correct branching ratios.

The calibration of the optical equipment was carried out using a calibrated tungsten ribbon lamp and a carbon arc lamp. No significant differences between the two lamps were found.

Furthermore, we investigated the coefficient of reflection of the fused silica reflector that we used for the calibration, and the transmission of the fused silica window (W_2). It can be shown that neither of these components can have caused an error as large as 40% in the calibration procedure.

Finally, we investigated experimentally whether the plasma absorbs the He II lines considered, or not. We found no absorption at all, neither a discrete absorption of the He II lines 4686 and 3203 Å, nor a continuous absorption spectrum in the neighbourhood of any of the He II lines mentioned in table IV.2. Besides, self-absorption of the He II lines considered is excluded on theoretical grounds.

In the foregoing we discussed two causes of the deviation of the branching ratios. The first is not very plausible, the second is excluded. Consequently, we are at a loss how to account for the deviations. This will have its repercussion on the determination of the population densities; an experimentally determined population density will depend on the lower level of the transition from which the density is derived.

In conclusion, we want to make a remark about the intensity profiles $L(u)$ (see fig. IV.1) and suggest a further investigation. Lines coming from upper levels 9, 10 and 11 and ending on different

lower levels, are not coincident. The deviations are such that the branching ratios derived from the feet of these intensity profiles are even smaller than those derived from the peaks, which branching ratios are already smaller than the "theoretical" ones.

This may be explained as follows: as regards the feet of the profiles, the levels are populated by He^{++} ; this is not occasioned by direct recombination into the excited levels considered, but via other, higher excited levels. If the cross-sections for the transitions between the sublevels of two neighbouring main levels obey the optical selection rule $\Delta l = \pm 1$, each of the two highest sublevels of a main level is populated by two sublevels of the above-lying main level, and is depopulated to one sublevel of the under-lying main level, whereas each of the other sublevels is populated and depopulated by an equal number of processes. Consequently, one may expect an overpopulation of the highest sublevels, which eventually leads to the deviating branching ratios. However, it is not clear why this process should be manifest only in levels 9 and higher, but not in lower ones.

Besides, line profile measurements show that the highest sublevel of main level 4 is under-populated instead of over-populated at some distance from the discharge-axis.

In any case, we do not have to worry about these branching ratios, because the population densities derived from the feet of the profiles are only of indirect importance to us (Chapter V.1).

For further investigation, we would suggest to perform similar measurements on a hydrogen plasma under comparable conditions. The transition probabilities for He^+ have been derived from transition probabilities of H, so in the H case one expects branching ratios corresponding to those for He^+ . With H, the energy distances between the sublevels are 16 times less than with He^+ , and therefore statistical population distribution of the sublevels is even more probable than it is already with He^+ .

1. Determination of the relevant quantities

a. Introduction

In the first part of this chapter we shall discuss the determination of the electron temperature and density, the density of the He^+ ions, and the population densities of the excited states 4 to 10. The measurements were performed at four different values for the arc current, viz. 35, 50, 70 and 100 A, and were directed to obtain values for the above quantities on the axis of the discharge. For this purpose, we had to do off-axis as well as on-axis measurements. The measurements were spread over a number of periods of time; within each period the measurements were reproducible. Comparable data from different periods showed differences.

It will be found that the contribution of the He^+ ground state to the population densities can be ascertained with much greater accuracy than that of the He^{++} ions. Therefore, we shall concentrate on the former contribution.

The second part of this chapter deals with the comparison of the observed population densities with the calculated ones. From this some aspects of relevant cross-sections will be discussed.

b. Electron temperature and density

These two quantities were determined with the help of Thomson scattering. The amount of scattered light was measured at five different wavelengths between 6832 Å and 6929 Å. At each wavelength 20 laser shots were taken with a repetition frequency of one per two minutes. The countings were corrected for the relative laser energy, the relative transmission of the interference filters and for the plasma light. Fig. V.1 shows an example of an observed Thomson spectrum. The straight line in this figure was fitted to the experimental points with the help of the method of least squares. The slope of this line is a measure of the electron temperature, and its intersection with the ordinate, together with the electron temperature, is a measure of the electron density (formulae III.5 and

III.6). In addition, the method of least squares yields the standard deviation of the slope and of the intersection with the ordinate. The error in the electron temperature is 5%, and that in the electron density is also about 5%. The latter error includes errors originating from the calibration procedure.

fig. V.1

An example of a Thomson spectrum measured.

The error bars indicate the standard errors.

$n_e = 2.6 \cdot 10^{19} \text{ m}^{-3}$, $T_e = 2.25 \cdot 10^5 \text{ K}$. Arc current 70 A.

Table V.1 shows the electron temperatures and densities on the axis of the discharge for four different arc currents.

Table V.1 Electron temperature (T_e) and density (n_e)

arc current (A)	T_e (10^5 K)	n_e (10^{19} m^{-3})
35	1.68	2.9
50	1.79	3.5
70	2.46	3.7
100	2.32	4.3

c. Population densities of the levels 4 to 10

The spectral lines used for the determination of the population densities of the levels 4 to 10 are given in table V.2. The profile of the radiance of each spectral line was scanned over the width of the plasma column. Each profile $L(u)$ shows a "peak" and a "foot" (fig. IV.1). First of all, we shall now show that the "foot" is due to population starting from the He^{++} ion, and the "peak" to population starting from the He^+ ground state. Furthermore, we shall show that it is permissible to extrapolate the "foot" from both sides through the axis of the discharge, so that the "peak" is superposed on a flat "foot".

It should be kept in mind that we are primarily interested in the contribution originating from the He^+ ground state.

Experimentally, we found that each foot has within $\pm 10\%$ a constant level from $u = 15$ mm to $u = 30$ mm, the end of the scanning trajectory. From $u = 15$ mm the radiant intensity slightly increases; at $u = 7$ mm the peak sets in. This observation can be understood when one assumes that the appropriate

Table V.2 Population densities n_i^{exp} of He^+ levels on the discharge axis

level i	wavelength ($\text{\AA}$)	transition i $\rightarrow$ j	error (%)	n_i^{exp} (10^{13} m^{-3})			
				35 A	50 A	70 A	100 A
4	4686	4 $\rightarrow$ 3	6	0.75	1.83	2.65	3.3
5	3203	5 $\rightarrow$ 3	6	0.66	1.69	2.3	2.7
6	6560	6 $\rightarrow$ 4	6	0.65	1.57	1.95	2.3
7	5412	7 $\rightarrow$ 4	6	0.40	0.94	1.14	1.48
8	4859	8 $\rightarrow$ 4	9	0.30	0.65	0.79	0.96
9	4542	9 $\rightarrow$ 4	7	0.17	0.33	0.50	0.56
10	4339	10 $\rightarrow$ 4	25	0.12	0.37	0.34	0.34

local population density does not decrease when the distance to the discharge axis (r) increases, from 15 mm up to at least 30 mm. Most probably the foot will extend beyond $r = 30$ mm. We consider $r = 70$ mm as the upper limit of the spatial extension of the foot, being the distance of diaphragm tube D_3 to the discharge axis (fig. III.1). We estimate the real spatial extension of the foot to be 50 mm.

We shall now turn our attention to the contributions of the He^+ and the He^{++} ions to the local population density of level 7 from $r = 0$ to 13 mm. We make use of data on the electron temperature and density obtained from Thomson scattering; with the help of these two quantities and tables of Drawin and Emard ([7] and private communication) the C.R. coefficients $r_7^{(0)}$ and $r_7^{(1)}$ were calculated, and these values finally yielded the population density due to He^{++} , and that due to He^+ , according to formula (I.8):

$$\frac{n_i}{n_s} = r_i^{(0)} + r_i^{(1)} \frac{n_1}{n_s} \quad (\text{V.1})$$

Furthermore, it was assumed that $\frac{n(\text{He}^{++})}{n} = 0.1$ everywhere; this assumption is rather arbitrary; we shall revert^e to this point later on. The result for an arc current of 70 A is shown in table V.3.

Table V.3 Contribution of He^{++} and He^+ to the population density of level 7 (n_7^{++} and n_7^{+} respectively)

r (mm)	$n_e (10^{19} \text{ m}^{-3})$	$T_e (10^5 \text{ K})$	$n_7^{++} (\text{m}^{-3})$	$n_7^{+} (\text{m}^{-3})$
0	2.57	2.25	$1.0 \cdot 10^{10}$	$7.0 \cdot 10^{12}$
2.7	2.56	1.98	$1.5 \cdot 10^{10}$	$5.0 \cdot 10^{12}$
5.3	2.16	1.40	$1.8 \cdot 10^{10}$	$1.1 \cdot 10^{12}$
8.0	1.84	0.93	$2.4 \cdot 10^{10}$	$1.1 \cdot 10^{11}$
10.7	1.66	0.68	$3.3 \cdot 10^{10}$	$9.3 \cdot 10^9$
13.3	1.54	0.54	$4.1 \cdot 10^{10}$	$8.7 \cdot 10^8$

From this table we conclude that the foot is produced by n_7^{++} . The conclusion is confirmed in the following way:

The "foot" population densities, integrated along the line of sight through the plasma, turn out to be in a Saha-Boltzmann equilibrium for

level $i = 6$ and higher. This observation is in accordance with the magnitude of the C.R. coefficients $r_i^{(0)}$: for $i > 6$ these coefficients are almost equal to 1 in the relevant range of electron density and temperature (they are only weakly dependent on n_e and T_e).

Table V.3 shows that beyond $r = 10$ mm only $n_7(++)$ contributes to the intensity profile $L(u)$ and, as expected, the contribution $n_7(++)$ increases with increasing r . For distances greater than 13 mm out of the discharge axis accurate data on the electron temperature are lacking.

In the next section it will be shown that the assumption $\frac{n(\text{He}^{++})}{n_e} = 0.1$ (independent of r) can account for the value of the radiant intensity of a foot. Having established that the foot of each intensity profile is due to He^{++} , we shall now show that it is permissible to extrapolate the foot through the axis of the discharge; in other words, if the contribution of the He^+ ground state to the population densities were zero, one would observe a rather flat foot, also for $r < 7$ mm where the peak exceeds the foot. We discuss two limiting cases:

- (a) the density of the He^{++} ions is zero for $r < 7$ mm. In this case, looking right through the centre of the plasma, the population density $n_i(++)$ integrated along the line of sight is less than if one looks at 7 mm out of the axis through the plasma; the lacking fraction is $\frac{14}{100}$, because, as mentioned above, the spatial extension of the foot is about 100 mm. Therefore, the foot will show a dip of 14% at the centre of the plasma ($u = 0$). If one extrapolates the foot at a constant level through the axis of the discharge, the remaining "peak" is too small. The error depends on the foot-to-peak ratio; the latter quantity is at most 30%, so the error in the peak height is less than 4%.
- (b) The density of the He^{++} ions is equal to $\frac{1}{2}n_e$ for $r < 7$ mm. Using data on the electron temperature and density (Thomson scattering), and on the C.R. coefficients $r_i^{(0)}$, the density $n_i(++)$, integrated from $r = 0$ to 7 mm can be calculated. This contribution, expressed as a fraction of the observed integrated foot population density, is 35%, 20% and 7% at arc currents of 35 A, 70 A and 100 A respectively. In the worst case (35 A) the error made in the determination of the peak height is 10% (level 10). We do not worry about an error of this magnitude.

We now pick up the line of this section again.

Each "peak" was submitted to an Abel transformation, and from the local emitted radiation power so obtained the population density on the axis of

the discharge was calculated (formula III.1) using values for the transition probabilities that are based on a statistical population distribution of the sublevels. The population densities, together with their experimental errors, are given in table V.2. The line 4339 Å was partly masked by two other lines.

d. Density of the He^+ ions

The density of the He^+ ions is not directly measurable. In many cases the quasi-neutrality relation $n(\text{He}^+) = n_e$ may be applied, but in our case the density of the He^{++} ions is not a priori negligible. On the axis of the discharge the production of He^{++} ions (by ionization of He^+) is about 2% of the production of the He^+ ions (by ionization of He), as calculated with the help of collisional-radiative ionization coefficients of Drawin and Emard [8].

Recombination is negligible for both He^{++} and He^+ ; the ions produced are transported out of the discharge axis by diffusion processes. There are a few publications dealing with this matter. For instance, Taylor [43] found that highly ionized impurities tend to concentrate on the discharge axis if the impurities have the same temperature as the electrons. His way of reasoning, applied to the case that the temperature of the impurities is much less than the electron temperature, leads to the conclusion that the concentration of the impurities is directly proportional to the electron density.

In our plasma, the He^{++} particles may be regarded as impurities in the sense that their production is much less than the production of the He^+ ions. Furthermore, the temperature of the He^{++} ions is expected to be lower than the electron temperature (the temperature of the He^+ ions is about 10^4 K, the electron temperature is about $2 \cdot 10^5$ K). On the basis of Taylor's work we expect that the density of the He^{++} ions is directly proportional to the electron density. However, the model that Taylor used did not consider interaction between unlike ions (He^+ , He^{++}), so we are not sure that it is permissible to apply the results to our plasma. It is our intention to investigate this matter; the best thing we can do at the present time is to suppose that $n(\text{He}^{++})/n_e$ is independent of r .

The density of the He^{++} ions on the axis of the discharge was estimated as follows:

The observed radiance of the foot of the line 7→4 yields a population density integrated along the line of sight. We estimate the length of the contributing column to be 10 cm (foregoing section), and so a mean population density, $\overline{n_7(++)}$, is obtained. As the local population density $n_7(++)$ increases by a factor 4 from $r = 0$ to 13.3 mm (table V.3), and will not decrease for $r > 13.3$ mm the population density $n_7(++)$ at $r = 0$ will be less than $\frac{1}{2}\overline{n_7(++)}$. This upper limit of $n_7(++)$, together with the electron temperature and density at $r = 0$, yields an upper limit for $n(\text{He}^{++})$.

Table V.4 shows the upper limits for the quantities $n(\text{He}^{++})$ and $n(\text{He}^{++})/n_e$ at four different arc currents.

Table V.4 Upper limit of the density of the He^{++} ions

I_{arc} (A)	n_e (10^{19} m^{-3})	T_e (10^5 K)	$r_7^{(0)}$	$\overline{n_7(++)}$ (10^{10} m^{-3})	$n(\text{He}^{++})$ (10^{18} m^{-3})	$\frac{n(\text{He}^{++})}{n_e}$
35	2.9	1.68	0.93	1.5	1.7	0.06
50	3.5	1.79	0.94	2.2	2.5	0.07
70	3.7	2.46	0.95	2.2	3.7	0.1
100	4.3	2.32	0.95	4.8	6.0	0.14

Finally, the density of the He^+ ions is estimated to be

$$n(\text{He}^+) = n_e - n(\text{He}^{++}) \quad (\text{V.2})$$

As the values for $n(\text{He}^{++})$ given above are upper limits, $n_e - 2n(\text{He}^{++}) < n(\text{He}^+) < n_e$. Therefore, the relative error in the estimated He^+ ion density is $n(\text{He}^{++})/n_e$.

2. Comparison of the observed population densities with the calculated ones

a. Choice of collisional-radiative calculation

At the present time, two C.R. calculations for the $e - He^+ - He^{++}$ system and applying to our plasma conditions are available: one by McWhirter and Hearn [6] and one by Drawin and Emard [7].

As shown by table I.3, corresponding C.R. coefficients $r_i^{(1)}$ representing the contribution of the He^+ ground state to the population densities, are considerably different with the two authors. For level $i = 4$, the difference is a factor 4, and it is a factor 6 for $i = 7$; for $i > 7$ the difference decreases.

In the foregoing section we showed that on the axis of the discharge the excited levels are for the greater part populated by He^+ (1s).

In the following it will appear that the experimental population densities are in reasonable agreement with the ones calculated by Drawin and Emard. Therefore, the calculations of McWhirter and Hearn will be discussed only at the end of this chapter (section 2.c.).

b. Comparison of the experimental population densities with the ones calculated with the C.R. model of Drawin and Emard

Let us summarize section 1 of this chapter:

- the population densities were determined on the axis of the discharge with the help of line intensity measurements; we consider that part of the population densities that is contributed by the He^+ ground level.
- the electron temperature and density were measured on the axis of the discharge with the help of Thomson scattering;
- the density of the He^+ ions was derived from the electron density combined with an estimate of the density of the He^{++} ions.

The C.R. coefficients $g_i^{(1)}$ appearing in the expression for the C.R. population densities (see formula (I.7))

$$n_i^{CR} = g_i^{(1)} n(He^+) \quad (V.3)$$

were calculated for our specific sets of (n_e, T_e) by Drawin and Emard. In table V.5 the elements of the C.R. calculation and the results are given. For purposes of comparison, the experimental population densities are also shown.

Table V.5 Calculated population densities n_i^{CR} of excited levels $i = 4$ to 10

level	$I_{arc} = 35 \text{ A}$			$I_{arc} = 50 \text{ A}$		
i	$n_e = 2.90 \cdot 10^{19} \text{ m}^{-3}$			$n_e = 3.50 \cdot 10^{19} \text{ m}^{-3}$		
	$n(\text{He}^+) = 2.73 \cdot 10^{19} \text{ m}^{-3}$			$n(\text{He}^+) = 3.25 \cdot 10^{19} \text{ m}^{-3}$		
	$T_e = 1.68 \cdot 10^5 \text{ K}$			$T_e = 1.79 \cdot 10^5 \text{ K}$		
	$g_i^{(1)}$	$n_i^{CR} (\text{m}^{-3})$	$n_i^{exp} (\text{m}^{-3})$	$g_i^{(1)}$	$n_i^{CR} (\text{m}^{-3})$	$n_i^{exp} (\text{m}^{-3})$
4	4.58(-7)	1.25(13)	0.75(13)	6.61(-7)	2.15(13)	1.83(13)
5	3.64	0.99	0.66	5.06	1.64	1.69
6	2.41	0.66	0.65	3.72	1.06	1.57
7	1.52	0.41	0.40	2.04	0.66	0.94
8	9.81(-8)	0.27	0.30	1.30	0.42	0.65
9	6.53	0.178	0.17	8.65(-8)	0.28	0.33
10	4.53	0.124	0.12	5.98	0.194	0.37

level	$I_{arc} = 70 \text{ A}$			$I_{arc} = 100 \text{ A}$		
i	$n_e = 3.70 \cdot 10^{19} \text{ m}^{-3}$			$n_e = 4.30 \cdot 10^{19} \text{ m}^{-3}$		
	$n(\text{He}^+) = 3.33 \cdot 10^{19} \text{ m}^{-3}$			$n(\text{He}^+) = 3.70 \cdot 10^{19} \text{ m}^{-3}$		
	$T_e = 2.46 \cdot 10^5 \text{ K}$			$T_e = 2.32 \cdot 10^5 \text{ K}$		
	$g_i^{(1)}$	$n_i^{CR} (\text{m}^{-3})$	$n_i^{exp} (\text{m}^{-3})$	$g_i^{(1)}$	$n_i^{CR} (\text{m}^{-3})$	$n_i^{exp} (\text{m}^{-3})$
4	1.59(-6)	5.29(13)	2.65(13)	1.58(-6)	5.85(13)	3.3(13)
5	1.24	4.13	2.3	1.18	4.37	2.7
6	8.09(-7)	2.69	1.95	7.54(-7)	2.79	2.3
7	5.10	1.70	1.14	4.69	1.74	1.48
8	3.29	1.10	0.79	3.00	1.11	0.96
9	2.20	0.73	0.50	1.99	0.74	0.56
10	1.53	0.51	0.34	1.38	0.51	0.34

4.58(-7) denotes $4.58 \cdot 10^{-7}$ etc.

In fig.V.2 the quantities n_i^{CR}/g_i , where $g_i = 2i^2$ is the statistical weight of level i , together with the experimental data are shown. It is seen that

fig. V.2 The population densities of excited levels 4 to 10 as a function of their ionization energy E_i . $\circ$: calculated population densities; $\times$: observed pop. densities; Δ : corrected population densities (levels 4,5 only)

(a) the experimental population densities of levels 4 and 5 are systematically deviating with respect to the other experimental population densities. The difference is removed by adding 40% to the experimental population densities of levels 4 and 5. This adjustment is connected with the branching ratios of table IV.2; the population densities determined from radiative transitions of the type $i \rightarrow 3$ are 40% lower than as determined from transitions of the type $i \rightarrow 4$. The latter type was used for the levels $i = 6$ to 10, the former for the levels 4 and 5.

(b) having applied a correction of +40% to the experimental population densities of the levels 4 and 5, the agreement between the experimental and the C.R. population densities is, in a relative sense, very good; it means that the percentual differences between the experimental and the calculated values are, at any arc current, almost the same ($\pm 10\%$) for all the levels considered.

(c) the deviations of the experimental from the C.R. population densities are: 0%, +45%, -30%, -15% for arc currents 35, 50, 70 and 100A, respectively. Therefore, on the whole the experimental and C.R. population densities are reasonably in agreement. The differences are not completely explained by random experimental errors:

- the error in the electron temperature, 5%, leads to an error of 10% in the C.R. population densities;
- the error in the electron density, 5%, also leads to an error of 10% in n_i^{CR} (by $g_i^{(1)} \cdot n(\text{He}^+)$!);
- the error in $n(\text{He}^{++})$ leads to an error of 10% in $n(\text{He}^+)$ (see foregoing section) and hence in n_i^{CR} ; however, our experiments do not exclude the possibility of the density of the He^+ ions being only half the electron density;
- the error in the experimental population densities may be put at 10% on the whole, apart from systematical errors connected with the branching ratios.

These errors can be considered to be independent of each other. The total error is obtained by adding individual errors quadratically together. The result is 20%. At $I_{\text{arc}} = 50\text{A}$, the experimental population densities are 45% greater than the C.R. population densities; this deviation of 2.2 times the standard error shows that the experimental error may be greater than 20%.

c. Discussion

As mentioned above, the experimentally determined population densities are in reasonably good agreement with the ones calculated by Drawin and Emard, and therefore disagree with the ones calculated by McWhirter and Hearn. The latter are about 4 times as great as the experimental population densities, and the two sets of population densities are also in a relative sense in disagreement. As far as we were able to trace, the differences between the results of the two calculations are caused by (Chapter I.5):

(1) Drawin's rate coefficients for excitation are $1/2$ to $1/3$ times the McWhirter's; this difference is for the greater part consequent upon Drawin's cross-sections for direct excitation, which, near the threshold, are half those of McWhirter's. The remaining factor 1.5 requires further investigation.

(2) Drawin's rate coefficients for (de)excitation between excited levels and for ionization of the excited levels are 2 to 3 times and 4 to 8 times as great as McWhirter's respectively.

We shall first discuss the cross-sections for direct excitation. If a level j is populated by direct excitation from the $\text{He}^+(1s)$ state, the sublevel jp is preferred over the other sublevels; for instance, the excitation cross-section $\sigma(1s-2p)$ is roughly 10 times greater than the cross-section $\sigma(1s-2s)$, and, according to Berry and Roesler [44],

$$\sigma(1s-4s) : \sigma(1s-4p) : \sigma(1s-4d) : \sigma(1s-4f) \approx 1 : 10 : 0.61 : 0.0025.$$

Theoretical calculations have been reported for $\sigma(1s-j_s)$ ($j = 2, 3, 4, 5$), and for $\sigma(1s-2p)$. Only the cross-section $\sigma(1s-2s)$ was investigated experimentally.

As regards $\sigma(1s-2s)$ we mention the calculations of Ormonde, Whitaker and Lipsky [45], 1967, and of McDowell, Morgan and Myerscought [46], 1973. These calculations show that the Coulomb-Born approximation applies to electron energies greater than $10\times$ the threshold energy. Near the threshold, more complex methods have to be used; the discussion of these is beyond the scope of this thesis. The threshold value calculated by McDowell et al. is about half the one calculated by Ormonde et al. Only two measurements of the excitation cross-section $\text{He}^+(1s-2s)$ have been reported; one by Dance, Harrison and Smith [47], 1966, the other by Dolder and Peart [48], 1973. These measurements give near-threshold values for $\sigma(1s-2s)$ that are in better accordance with the calculations of McDowell et al. than with those of Ormonde et al.

With respect to $\sigma(1s-2p)$ the calculations of Burgess [10], 1961, of Burgess, Hummer and Tully [49], 1970, and of Burke, McVickar and Smith [50], 1964, may be mentioned. Their results are in agreement with each other.

The cross-section $\sigma(1s-2p)$ is almost constant for electron energies between the threshold energy and three times the threshold energy (fig. I.2, curve 3). This part of the cross-section contributes more than 90% to the rate coefficient C_{1s-2p} at an electron temperature of $2.5 \cdot 10^5$ K (21.5 eV). Neither theoretical nor experimental data were published on cross-sections of the type $He^+ (1s-jp)$, $j > 2$. Both, McWhirther and Hearn, and Drawin and Emard based their rate coefficients C_{1s-jp} on the theoretical cross-section $\sigma(1s-2p)$ of Burgess. We discussed this in Chapter I.4.

The results of our measurements are in agreement with the calculations of Drawin and Emard; therefore, they indicate very strongly that the near-threshold values of the cross-sections $\sigma(1s-jp)$ are indeed half the theoretical ones. The latter statement is directly demonstrable for level 4; as mentioned earlier in the present work, this level is mainly populated by direct excitation from the ground state and depopulated by spontaneous emission; its population density is therefore directly proportional to the rate coefficient involved, which is directly proportional to the near-threshold value of the corresponding cross-section. In this way the near-threshold value of the cross-section $\sigma(1s-4p)$ is found to be half the one expected from the calculations of Burgess. We estimate the error of the experimentally determined cross-section at 50%. It includes the deviations at different arc currents (fig. V.2), as well as the uncertainty originating from the branching ratios. It was assumed that the error in the density of the He^+ ions does not exceed 10%.

When for a given couple of parameters (n_e, T_e), i.e. for a given value of the arc current, the observed population densities are divided by the factors 1, 1.45, 0.70 and 0.85 for arc currents 35, 50, 70 and 100 A respectively, the calculated and the observed population densities agree within $\pm 10\%$.

The behaviour of the population densities as a function of the main quantum number i depends on the (de)excitation processes between excited levels and on the ionization rates of the excited levels.

Drawin and Emard based the corresponding rate coefficients on theoretical cross-sections of Burgess (Chapter I.4). Because the mean electron energy is much higher than the threshold energies (20 eV compared to a few eV or less)

only the "tails" of the cross-sections are of importance. Here, the cross-sections can be calculated with the help of the Born-Bethe expression [11]. McWhirter and Hearn followed in principle the same procedure; however, they used approximation formulae for the rate coefficients instead of calculating them exactly like Drawin and Emard did. As a result, their rate coefficients are several times as small as the ones of Drawin and Emard (at electron temperatures relevant to our experiment), and finally the dependence of their C.R. population densities on the main quantum number is different from that of Drawin and Emard.

On the other hand, the fact that our experimental population densities are relatively within 10% in agreement with the C.R. population densities of Drawin and Emard, means an experimental justification of the cross-sections they used for (de)excitation between excited levels and for ionization of excited levels. An alteration of these cross-sections by more than 50% would be observable. This may be illustrated by table V.6. This table shows, for $n_e = 3.70 \cdot 10^{19} \text{ m}^{-3}$, $n(\text{He}^+) = 3.33 \cdot 10^{19} \text{ m}^{-3}$, and $T_e = 2.46 \cdot 10^5 \text{ K}$:

- (a) the calculated population densities, if the rate coefficients C_{ij} ($i, j \neq 1$) and S_i are 2 times greater as "normal".
- (b) the calculated population densities, if the rate coefficients are unaltered (corresponding to the n_i^{CR} tabulated in table V.5, for $I_{\text{arc}} = 70 \text{ A}$).
- (c) the calculated population densities, if the rate coefficients C_{ij} ($i, j \neq 1$) and S_i are half as great as "normal".
- (d) the observed population densities at $I_{\text{arc}} = 70 \text{ A}$; all densities have been multiplied by a factor 1.43, and superposed on this, the observed population densities of the levels 4 and 5 have been multiplied by a factor 1.4.

Table V.6 The effect of variation of rate coefficients

level	case (a)	case (b)	case (c)	case (d)
i	$n_i^{\text{CR}} (\text{m}^{-3})$	$n_i^{\text{CR}} (\text{m}^{-3})$	$n_i^{\text{CR}} (\text{m}^{-3})$	$n_i^{\text{exp}} (\text{m}^{-3})$
4	4.73(13)	5.29(13)	5.67(13)	5.30(13)
5	3.12	4.13	5.11	4.60
6	1.85	2.69	3.76	2.80
7	1.12	1.70	2.54	1.60
8	0.70	1.10	1.70	1.13
9	0.46	0.73	1.16	0.72
10	0.31	0.51	0.82	0.49

This table clearly demonstrates that the observed population densities are not in agreement with cases (a) and (c).

In the foregoing we showed that the observed population densities are in reasonably good agreement with the ones calculated with a collisional-radiative model and observed electron temperatures and densities. In reverse, one can deduce the electron temperature and density from the observed population densities with the help of the results of collisional-radiative calculations. This procedure is of interest to many astrophysical and laboratory (fusion) plasmas, as mentioned at the end of Chapter I. For example, we consider our measurements for an arc current of 70 A (table V.5) and put the question: what electron temperature and density would one deduce from the observed population densities? The answer is (see table V.7) : applying the collisional-radiative calculations of Drawin and Emard,

$$3.3 \cdot 10^{19} < n_e < 4.4 \cdot 10^{19} \text{ (m}^{-3}\text{)}, \quad 2.39 \cdot 10^5 > T_e > 1.87 \cdot 10^5 \text{ (K)},$$

applying the collisional-radiative calculations of McWhirter and Hearn,

$$n_e \sim 10 \cdot 10^{19} \text{ (m}^{-3}\text{)}, \quad T_e \sim 1 \cdot 10^5 \text{ (K)},$$

whereas Thomson scattering yielded

$$n_e = 3.7 \cdot 10^{19} \text{ (m}^{-3}\text{)}, \quad T_e = 2.46 \cdot 10^5 \text{ (K)}.$$

This clearly demonstrates that the resulting n_e and T_e strongly depend on the type of collisional-radiative calculation one applies.

Table V.7 Determination of n_e and T_e from observed population densities

$n_e \text{ (m}^{-3}\text{)}$	$T_e \text{ (K)}$	$n_{10}^{\text{exp}} \text{ (m}^{-3}\text{)}$	$n_{10}^{\text{CR(1)}}$	$n_7^{\text{exp}} \text{ (m}^{-3}\text{)}$	$n_7^{\text{CR}} \text{ (m}^{-3}\text{)}$
$3.3 \cdot 10^{19}$	$2.39 \cdot 10^5$	$3.4 \cdot 10^{12}$	$4.0 \cdot 10^{12}$		
$3.7 \cdot 10^{19}$	$2.12 \cdot 10^5$	$3.4 \cdot 10^{12}$	$3.4 \cdot 10^{12}$		
$4.4 \cdot 10^{19}$	$1.87 \cdot 10^5$	$3.4 \cdot 10^{12}$	$2.9 \cdot 10^{12}$		
$\approx 10 \cdot 10^{19}$	$\sim 1 \cdot 10^5$	$3.4 \cdot 10^{12}$		$1.1 \cdot 10^{13}$	$\approx 1.1 \cdot 10^{17}$

In all cases $n_4^{\text{exp}} = n_4^{\text{CR}}$ and $n(\text{He}^+) = 0.9.n_e$. CR(1) denotes:
calculated with C.R. coefficients of Drawin and Emard; CR(2) denotes:
from McWhirter and Hearn.

APPENDIX I. Additional information on the Thomson scattering experiment

1. Transmission characteristics of the two interference filters

The transmission of the interference filters at a tilting angle of 0° , i.e. at normal incidence, has been measured in relation to the wavelength of the incident light. The optical arrangement is shown in fig. AI.1. It is built

fig. AI.1

Optical arrangement for the determination of the transmission properties of the interference filters.

The focal lengths of the lenses, the openings of the diaphragms and the distances between the components are given in mm. L denotes a projection lamp, D a diffusing screen, IF the interference filters. RG 665 is a red transmitting blocking filter. The optical arrangement starting from mirror M_1 is the same as shown in fig. III.1.

up so that the interference filters were illuminated in the same way as they were illuminated in the scattering experiment. The transmitted light was focussed on a diffusing screen using two lenses and the line intensity optical components. The illuminated part of the screen had an area of $2 \text{ cm} \times 2 \text{ cm}$. The screen was placed in front of the entrance slit of the Jarrel Ash monochromator. The slit was illuminated in this way in order to exclude any beam selection by the monochromator, as will occur by direct focussing on the slit. Both the entrance and the exit slits were $10 \mu\text{m}$ wide, the entrance slit was 2 mm high. The grating had 1180 lines/mm and was blazed for $3000 \text{ \AA}$. The resolution of the monochromator in the appropriate wavelength region was $0.25 \text{ \AA}$, as derived from the measured

half-width of a narrow spectral line.

The transmission curve of the interference filters (fig. AI.2) was constructed by scanning the wavelength with the monochromator. The effective width of the filters was determined from the transmission curve. A correction of $-0.25 \text{ \AA}$ was applied to allow for the finite resolution of the monochromator. The position of the peak wavelength (the wavelength at maximum transmission) was located with the help of an emission line of Ar at $6965.4 \text{ \AA}$.

The results are:

position of the peak	$6944.2 \text{ \AA}$
transmission at peak wavelength	19%
effective width at peak wavelength	$6.3 \text{ \AA}$.

fig. AI.2

Transmission characteristic of the interference filters illuminated at normal incidence.

If in the scattering experiment the filters are tilted, the effective width increases owing to the divergence of the light: at $\theta = 0^\circ$, $\Delta\lambda_w$ is $6.3 \text{ \AA}$, at 12.5° ($6871 \text{ \AA}$): $12 \text{ \AA}$ and at 20° ($6753 \text{ \AA}$): $20 \text{ \AA}$. As the effective width increases the peak transmission decreases, and as a result the transmission integrated over the complete profile decreases. This integrated transmission has been determined in an arbitrary unit in relation to θ by placing a tungsten ribbon lamp at the position of the plasma (PL, fig. III.1) and measuring the photomultiplier signal.

If the integrated transmission is put at 100 for $\theta = 0^\circ$, it is 91 for $\theta = 10^\circ$, 74 for $\theta = 15^\circ$, 49 for $\theta = 20^\circ$. At greater tilting angles the

transmission rapidly decreases. The Thomson signals have been corrected for this behaviour of the integrated transmission. The Thomson spectrum itself is not significantly influenced by the finite width of the interference filters. This width is about 10 times less than that of the Thomson spectrum; the latter has a Gaussian profile, the transmission curve of the interference filters is approximatively Gaussian. The correction factor for the width of the interference filters has therefore the negligible value of $(0.1)^2$.

2. Calculation of T_e and n_e from the Thomson scattering data

The Thomson spectrum $S(\Delta\omega)$ can be represented by

$$S(\Delta\omega)d(\Delta\omega) = n_e \left(\frac{m_e}{2\pi k^2 \kappa T_e} \right)^{\frac{1}{2}} e^{-\frac{m_e (\Delta\omega)^2}{2 k^2 \kappa T_e}} d(\Delta\omega) \quad (\text{AI.1})$$

where:

$\Delta\omega$ is the frequency shift (in rad s^{-1}) between the scattered light (ω_s) and the incident light (ω_o)

$\vec{k}$ is the scattering vector, defined by $\vec{k} = \vec{k}_s - \vec{k}_o$, where $\vec{k}_s$ and $\vec{k}_o$

are the wave vectors ($|\vec{k}| = 2\pi/\lambda$) of the scattered and the incident light respectively

m_e is the electron mass

κ is Boltzmann's constant

n_e is the electron density

T_e is the electron temperature.

The area of the Thomson spectrum, $\int S(\Delta\omega)d(\Delta\omega)$, is equal to n_e . Instead of $\Delta\omega$ one can also use $\Delta\lambda = |\lambda_s - \lambda_o|$, the absolute difference between the wavelengths of the scattered and incident light. Using the relation $k = k_o/2$ (scattering at 90°) the spectral function becomes

$$S(\Delta\lambda)d(\Delta\lambda) = n_e \frac{2\pi c}{\lambda^2} \left(\frac{m_e \lambda^2}{16\pi^3 \kappa T_e} \right)^{\frac{1}{2}} e^{-\frac{m_e c^2}{4\kappa T_e \lambda^2} \Delta\lambda^2} d(\Delta\lambda) \quad (\text{AI.2})$$

The wavelength shift $\frac{1}{2}\Delta\lambda_T$, for which $S(\frac{1}{2}\Delta\lambda_T) = \frac{1}{2} S(0)$, is given by

$$\frac{m_e c^2}{4\kappa T_e \lambda^2} \left(\frac{1}{2}\Delta\lambda_T\right)^2 = \ln 2 \quad (\text{AI.3})$$

Inserting the numerical values for the constants, one obtains

$$T_e = 11.1 \Delta\lambda_T^2 \quad (\text{AI.4})$$

T_e in K, $\Delta\lambda_T$ in Å.

The fraction of the total Thomson scattered light present in a wavelength interval $\Delta\lambda_W$, centered on λ_0 is

$$\frac{S_T(0)}{S_T} = \frac{S(0)\Delta\lambda_W}{n_e} = \frac{c}{\lambda} \left(\frac{m_e}{4\pi\kappa T_e}\right)^{\frac{1}{2}} \Delta\lambda_W \quad (\text{AI.5})$$

The ratio of the total Thomson scattered light from an electron gas with density n_e , to the Rayleigh scattered light $S_R(0)$ from an atomic gas with density n_R is

$$\frac{S_T}{S_R(0)} = \frac{n_e \sigma_T}{n_R \sigma_R} \quad (\text{AI.6})$$

where σ_T is the Thomson cross-section and σ_R is the Rayleigh cross-section. (AI.6) with (AI.5) yields

$$\frac{S_T(0)}{S_R(0)} = \frac{n_e \sigma_T}{n_R \sigma_R} \frac{c}{\lambda} \left(\frac{m_e}{4\pi\kappa T_e}\right)^{\frac{1}{2}} \Delta\lambda_W \quad (\text{AI.7})$$

After insertion of the numerical values for c , λ , m_e and κ , and using relation (AI.4), relation (AI.7) becomes

$$n_e = 1.06 \frac{S_T(0)}{S_R(0)} \frac{\sigma_R}{\sigma_T} \frac{\Delta\lambda_T}{\Delta\lambda_W} n_R \quad (\text{AI.8})$$

We measured Rayleigh scattering of Ar; $\sigma_R/\sigma_T = 1/427$ (see next section).

Supposing the gas pressure to be 10 torr at 20°C, $n_R = 3.30 \cdot 10^{23} \text{ m}^{-3}$;

$\Delta\lambda_W$ is 6.3 Å (Appendix I.1). (AI.7) now reads

$$n_e = 1.30 \cdot 10^{20} \frac{S_T(0)}{S_R(0)} \Delta\lambda_T \quad n_e \text{ in } \text{m}^{-3}, \Delta\lambda_T \text{ in } \text{Å} \quad (\text{AI.9})$$

3. Rayleigh scattering cross-section

As outlined in Chapter III.3 and in the foregoing section, the cross-section for Rayleigh scattering of Ar is of direct importance to the determination of the electron density. This cross-section is related to the index of refraction by

$$\sigma_R = \frac{8\pi^3}{3n_R^2} \frac{(\mu^2 - 1)^2}{\lambda^4} \quad (\text{AI.10})$$

where: σ_R is the total scattering cross-section (for a solid angle of 4π sr)

n_R is the particle density of the gas

μ is the index of refraction of the gas having density n_R

λ is the wavelength of the incident radiation.

The quantity $(\mu^2 - 1)^2 / n_R^2$ is almost wavelength independent. Values for $\mu - 1$ can be found in [51]; at a temperature of 0° C, and a pressure of 76 cm Hg, $\mu - 1 = 281 \cdot 10^{-6}$ in the case of Ar. For $\lambda = 6943 \text{ \AA}$ one finds

$$\sigma_R = 1.56 \cdot 10^{-31} \text{ m}^2 \quad (\text{AI.11})$$

Some investigators have tried to determine this cross-section experimentally, using a ruby laser. Their results were not always in agreement with the theoretical cross-section. George et al. [52] reported experimental cross-sections for various gases, among which Ar, that are about twice as large as the theoretical values. They used a normal mode ruby laser having an output energy of 0.1 J per shot. Decroisette and Cabannes [53], also using a normal mode ruby laser, found an experimental value for Ar that was in agreement with the theoretical one; however, their theoretical cross-section is a factor 4 too small, and therefore their experimental cross-section is 4 times too small.

Rudder and Bach [54] obtained cross-sections for various gases (among which Ar) that are in excellent agreement with the theoretical ones. They used a Q-switched laser of 8 MW power.

These results induced us to check the Rayleigh cross-section of Ar experimentally.

The experimental set-up is completely identical with that of the Thomson scattering experiment. Now we had to perform a calibration of the optical equipment. We chose the following procedure:

- (1) Measurement of the absolute laser light energy passing through the scattering volume.

(2) Determining the efficiency of the detection device.

The absolute laser energy was measured at the position of diaphragm D_2 (fig. III.3) with the help of a laser calorimeter of Instrument Technology. The inaccuracy of the calorimeter is 5% (manufacturer's rating). We had to use an extra lens to focus the laser light in the calorimeter. Therefore, a correction of +8% was applied to the reading of the laser calorimeter. The light energy passing through the scattering volume was 29 J per laser shot on the average.

The efficiency of the detection device was determined by "replacing" the scattering volume by a calibrated tungsten ribbon lamp. In this way we established the relation between the spectral radiance at the position of the scattering volume and the counting rate on the counter. Finally, we needed the value of the effective width of the interference filters (section 1 of this appendix) and the transmission of one additional grey filter (transmission $2.73 \cdot 10^{-2}$). Using the results of this calibration and the data of a Rayleigh scattering experiment with Ar at a pressure of 15 torr, we calculated the differential Rayleigh cross-section for a scattering angle of 90°

$$\left(\frac{d\sigma_R}{d\Omega}\right)_{90^\circ}^{\text{exp}} = (1.72 \pm 0.17) \cdot 10^{-32} \quad (\text{m}^2) \quad (\text{AI.12})$$

The theoretical differential cross-section for scattering at 90° is:

$$\left(\frac{d\sigma_R}{d\Omega}\right)_{90^\circ}^{\text{th}} = \frac{3}{8\pi} \sigma_R = 1.84 \cdot 10^{-32} \quad (\text{m}^2) \quad (\text{AI.13})$$

The agreement is satisfactory. Therefore, we have no reason to doubt the theoretical cross-section.

In addition we carried out some supplementary measurements on Rayleigh scattering:

- The pressure of the Ar gas was varied between 2 torr and 160 torr. This did not affect the Rayleigh cross-section.
- We attenuated the laser energy in steps, ultimately $3.6 \cdot 10^{-6}$ times the unattenuated energy, using several calibrated grey filters. The Rayleigh cross-section was independent of the laser energy.
- The Rayleigh cross-section of other gases than Ar, viz. He, Xe, N_2 was measured. The experimental cross-sections were in agreement with the theoretical ones.

APPENDIX II. An investigation on the population densities of He levels

The purpose of this investigation was:

(a) to determine the He^+ density on the axis of the arc,

(b) to determine the density of the neutral He atoms on the axis of the arc.

As will be shown below the first goal was not attained, the second was.

These investigations have only been carried out for one value of the discharge current, viz. 70 A.

(a) Determination of the density of He^+ ions

Line intensities of spectral lines of the series $2^3\text{P}-n^3\text{D}$ were determined in order to see whether the population densities of the $n^3\text{D}$ levels show a Saha-Boltzmann distribution for high quantum numbers n . If this were the case, the density of the He^+ ions could be calculated. Upper levels up to $n = 18$ were investigated.

As is the case with the intensity profiles of the He II lines, the intensity profiles $I(u)$ of the He I lines show a "foot" and a "peak". Similar to the He II case the "foot" is due to population from the He^+ ground state, the "peak" is due to (in)direct excitation from the $\text{He } 1s_0$ ground state. Calculations made by Drawin and Emard ([55] and private communication) show that on the axis of the discharge ($T_e = 2.5 \cdot 10^5 \text{ K}$, $n_e = 3.7 \cdot 10^{19} \text{ m}^{-3}$) the contribution of He exceeds that of He^+ by a factor 3 for level 18. Outside the axis the former contribution rapidly decreases owing to the decrease in the electron temperature.

The foot of the intensity profile is not suitable for an Abel transformation, because its behaviour close to the discharge-axis and the boundary conditions are unknown. Therefore, it is not possible to determine the density of the He^+ ions on the axis of the discharge from spectral lines of He I.

(b) Determination of the density of the He atoms

We measured the line intensities of three He I lines in absolute value: $5016 \text{ \AA}$ ($3^1\text{P} - 2^1\text{S}$), $3889 \text{ \AA}$ ($3^3\text{P} - 2^3\text{S}$) and $3554 \text{ \AA}$ ($10^3\text{D} - 2^3\text{P}$). Experimentally, the contribution of the He^+ ions was separated from that of the He atoms. After an Abel transformation (line 3554 with poor precision), the population densities of the excited levels involved on the axis of the discharge were obtained.

On the other hand, these population densities were calculated with collisional-radiative coefficients of Drawin and Emard (private communication) for the optically thin case.

The C.R. contribution of the He ground state to the population density of level i is

$$n_i = r_i^{(1)} n(\text{He}) g_i e^{\frac{E_{o,i}}{kT_e}} \quad (\text{AII.1})$$

where: $r_i^{(1)}$ is a collisional-radiative coefficient

$n(\text{He})$ is the population density of the ground state 1^1S_o

g_i is the statistical weight of sublevel i ; $g(^3P) = 9$,
 $g(^3D) = 15$, $g(^3P) = 3$

$E_{o,i} = E_i - E_o$, the difference between the ionization energies of the levels i and o (ground level).

The other symbols have their usual meanings. The C.R. coefficient $r_i^{(1)}$ is a function of the electron density and temperature. The latter quantities were determined with Thomson scattering. The only unknown quantity is $n(\text{He})$.

Comparison between the calculated and the measured densities yields $n(\text{He})$.

The result is:

level 3^1P gives $n(\text{He}) = 1.2 \cdot 10^{19} \text{ m}^{-3}$

level 3^3P gives $n(\text{He}) = 2.3 \cdot 10^{19} \text{ "}$

level 10^3D gives $n(\text{He}) = 1.7 \cdot 10^{19} \text{ "}$

The accuracy of these results is in the first place determined by the C.R. coefficients $r_i^{(1)}$, which should be accurate within a factor 2.

Absorption of radiation may have some influence upon the calculated population density of level 3^1P . We conclude

$$n(\text{He}) = (1 - 3) \cdot 10^{19} \text{ (m}^{-3}\text{)} \quad (\text{AII.2})$$

This density may be compared with that far from the axis of the arc, where the background pressure is about $2 \cdot 10^{-3}$ torr, giving $n(\text{He}) = 6.6 \cdot 10^{19} \text{ m}^{-3}$.

Therefore, approaching the centre of the discharge the density is lowered, owing to ionization to He^+ .

APPENDIX III. Measurements with a Langmuir probe and a microwave interferometer

1. Langmuir probe measurements

With the help of a Langmuir probe the electron density was measured from 10 to 100 mm out of the axis of the discharge. At shorter distances the probe became heated too much.

Theories on the interpretation of the probe curve are very complex, and in the case of a magnetic field, in which the electron Larmor radius is less than the dimensions of the probe, they are not complete (see contribution of Chen in [29]). In spite of the theoretical difficulties it seems to be fairly safe to calculate the ion density from the ion saturation current with

$$I_+^s = \frac{1}{2} n_+ q_+ A \left(\frac{k T_e}{m_+} \right)^{\frac{1}{2}} \quad (\text{AIII.1})$$

where: I_+^s is the ion saturation current

n_+ is the ion density

q_+ is the ion charge

A is the probe area

m_+ is the ion mass

T_e is the electron temperature.

The electron temperature may be determined from the slope of the electron retardation curve with

$$\frac{d \ln I_e}{dV} = \frac{q_e}{k T_e} \quad (\text{AIII.2})$$

The electron current I_e is obtained by subtracting the extrapolated ion saturation current from the measured probe current. V is the probe potential.

We used two types of probe: a spherical one, 3 mm in diameter, and a cylindrical one, 0.4 mm in diameter and 5.5 mm long.

The results as regards the electron density and temperature are almost the same for the two probes, see fig. A III.1 and fig. A III.2. In these figures the results obtained from Thomson scattering are also indicated.

The latter measurements gave electron densities that are about 50% greater than those derived from the probe measurements, and electron temperatures

fig. AIII.1

Electron density as a function of the radial distance to the arc-axis.

- ▲ measured with a spherical probe
- measured with a cylindrical probe
- x obtained from Thomson scattering.

that roughly agree with those measured by means of the probes. The region of overlap between the two diagnostic methods is too small to permit a more decisive comparison. It should be emphasized that the Langmuir probe detects only the electrons in the tail of the velocity distribution, while Thomson scattering is concentrated on the bulk of the electrons.

fig. AIII.2 *Electron temperature as a function of the radial distance to the arc-axis.*
For symbols see fig. AIII.1.

Finally, we give some characteristic distances:

Debye length in our plasma: $5 \cdot 10^{-6}$ m

electron Larmor radius: $5 \cdot 10^{-5}$ m

ion Larmor radius: $2 \cdot 10^{-3}$ m.

2. Microwave interferometer

The experimental set-up is like that of a Mach-Zehnder interferometer (fig. A III.3). A klystron produces a 4 mm electromagnetic wave, which is split up into two parts by a directional coupler; one part is guided through the plasma, the other is led through an attenuator and two phase shifters. The two waves interfere near a detector. The electromagnetic wave is an ordinary wave in the plasma in the applied magnetic field. The index of

fig. AIII.3

Microwave equipment.

P, isolator; Q, attenuator; R, frequency meter; S, directional coupler; T, phase shifter; V, tuner; W, horn; X, hybrid tee; Y, calibrated attenuator.

refraction of the plasma is in this case

$$\mu = \left(1 - \frac{n_e}{n_c}\right)^{\frac{1}{2}} \quad (\text{AIII.3})$$

Where n_e is the actual electron density, and n_c the critical electron density, defined by

$$n_c = \frac{\epsilon_0 m_e}{q_e} \omega_0^2 \quad (\text{AIII.4})$$

where ω_0 is the frequency of the klystron. In our case, $\omega_0 = 4.2 \cdot 10^{11}$ rad s⁻¹, and $n_c = 5.6 \cdot 10^{19}$ m⁻³. The phase shift of the transmitted wave with respect to the situation without plasma is

$$\Delta\phi = 2\pi \int_0^d \frac{\mu(y) - 1}{\lambda} dy \quad (\text{AIII.5})$$

in case the plasma can be considered as an infinite slab in the $x - z$ direction, having a thickness d in the y direction.

Furthermore, it is assumed that the index of refraction does not change much over a distance of one wavelength. The phase shift was measured

for several plasma conditions (pressure, arc current, magnetic field), and is 10 - 15 radians. For one case we calculated the phase shift with the help of formula (AIII.5) and combined Thomson scattering and probe measurements (fig. A III.1). The result, 10 radians, is somewhat lower than the actual phase shift. The plasma region that extends to 13 mm from the discharge-axis contributed 60% to the calculated phase shift.

It is not possible to obtain a value for the electron density on the discharge-axis from microwave measurements alone. One needs additional information on the radial distribution of the electron density. In our case, this information could only be supplied by Thomson scattering. Therefore, the microwave measurements gave us no new information. These measurements only confirmed the results already obtained by Thomson scattering.

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SUMMARY

The population densities of excited atomic levels in a gas discharge are, generally, functions of many variables like the electron temperature, the electron density, the temperatures and the densities of the atoms and ions present in the plasma, the radiation field, etc. Under certain circumstances only the electron temperature and density, the density of the atoms in the ground state and of those in the next ionized state are of importance to the population densities. This is the domain of the collisional-radiative model. Using it, the population densities can be calculated as functions of the four quantities mentioned above by taking into account all the relevant populating and depopulating processes.

Two collisional-radiative calculations on hydrogen-like ions are checked by determining experimentally the population densities of excited He^+ ions, the electron density and temperature and the densities of the He^+ and He^{++} ions. The observations are made on a hollow cathode arc discharge. Line intensities have been used to measure the population densities. Thomson scattering has been applied to determine the electron density and temperature, and the densities of He^+ and the He^{++} ions have been derived from the electron density, and line intensities.

The assumptions incorporated in the collisional-radiative calculations, have been investigated. There is evidence to doubt the statistical population distribution of the sublevels of the excited levels.

The observed population densities are in agreement with the ones calculated by Drawin and Emard, and disagree with those of McWhirter and Hearn. The results of the two calculations are different partly because of differences between the values of cross-sections for electronic excitation of the He^+ levels from the He^+ ground level.

The cross-section for the excitation process $1s + 4p$ near the threshold energy has been calculated from the experimental results. It is half the one expected from calculations of Burgess.

No cross-section of this type has been determined experimentally before.

De bezettingen van de aangeslagen atoomniveaus in een gasontlading zijn in het algemeen functies van een groot aantal variabelen zoals de electronentemperatuur, de electronendichtheid, de temperaturen en de dichtheden van de in het plasma voorkomende atomen en ionen, het stralingsveld, enz. Onder bepaalde omstandigheden zijn slechts de electronentemperatuur en -dichtheid, de dichtheid van de grondtoestand en van de eerst volgende geïoniseerde toestand van de atomen van belang voor de bezettingen van de aangeslagen toestanden. Dit is het terrein van het zg. collisional-radiative model. Met behulp van dit model worden de bezettingen berekend als functie van de bovengenoemde vier grootheden door de relevante bevolkingsprocessen en ontvolkingsprocessen in rekening te brengen.

Twee van dergelijke berekeningen betreffende waterstofachtige ionen zijn aan de praktijk getoetst door middel van meting van de bezettingen van de aangeslagen toestanden van He^+ ionen, de electronendichtheid en -temperatuur en de dichtheden van de He^+ en He^{++} ionen.

De waarnemingen zijn verricht aan een holle kathode boogontlading. Bezettingen zijn gemeten met behulp van lijnintensiteiten, de electronentemperatuur en -dichtheid met Thomson-verstrooiing en de dichtheden van de He^+ en de He^{++} ionen zijn afgeleid van de electronendichtheid en lijnintensiteiten.

De veronderstellingen, ingebouwd in de collisional-radiative berekeningen, zijn onderzocht. Er zijn aanwijzingen gevonden dat de subniveaus van de aangeslagen toestanden niet bezet zijn volgens hun statistische gewichten.

De waargenomen bezettingen komen duidelijk beter overeen met die berekend door Drawin en Emard dan met die van McWhirter en Hearn. De twee berekeningen leveren verschillende resultaten, onder meer vanwege een verschil in grootte van de werkzame doorsneden voor aanslag van de hogere toestanden vanuit de He^+ grondtoestand door electronen.

Uit de experimentele resultaten is de werkzame doorsnede voor de aanslag $1s \rightarrow 4p$ voor een electronenenergie in de buurt van de drempelenergie bepaald. Deze werkzame doorsnede is ca. de helft van wat op grond van berekeningen van Burgess te verwachten is.

Een experimentele bepaling van deze werkzame doorsnede is nog niet eerder gedaan.

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STELLINGEN

behorend bij het proefschrift van

W. Kohsiek

Eindhoven, 25 juni 1974.

1. Bij toepassing van de formule van Bohm betreffende het verband tussen de ionenverzadigingsstroom door een sonde aan een plasma onttrokken en de ionendichtheid in het plasma, dient de ionenverzadigingsstroom te worden bepaald bij een sondepotentiaal die nagenoeg gelijk is aan de plasmapotential.

Chen, in "Plasma Diagnostic Techniques", e.d. Huddleston and Leonard, New York 1965, blz. 150.

2. In tegenstelling tot wat Bates et al. beweren, kan hun rate coefficient voor ionisatie van waterstofachtige ionen meer dan een factor 2 keer zo groot zijn als die welke resulteert uit de werkzame doorsnede van Burgess.

Bates, Kingston en McWhirter, Proc.Roy.Soc.London, A 267, (1962)297
Burgess, Mem.Soc.Roy.Liège (5) 4 (1961) 299.

3. De relaxatietijden voor instelling van (quasi)stationaire bezettingen, berekend door McWhirter en Hearn, en door Drawin, mogen niet toegepast worden op systemen waar meer dan één aangeslagen toestand een bezetting heeft, verschillend van zijn evenwichtsbezetting.

McWhirter en Hearn, Proc.Phys.Soc. 82 (1963) 641
Drawin, J.Q.S.R.T. 10 (1970) 33.

4. Het is niet toegestaan zonder meer de "branching ratio", welke gebaseerd is op een statistische bezettingsverdeling van de atomaire subniveaux, toe te passen om in geval van een $e\text{-He}^+ \text{-He}^{++}$ plasma met een electronendichtheid van ca. 10^{19} m^{-3} de intensiteit van de éne spectraallijn te bepalen uit de intensiteit van de andere.

Dit proefschrift, hoofdstuk IV.

5. De door Cabannes en Decroisette gemeten werkzame doorsnede voor Rayleigh-verstrooiing door Ar gas is een factor 4 te klein.

Cabannes en Decroisette, J.Q.S.R.T. 8 (1968) 619.

6. De werkzame doorsnede voor Rayleigh-verstrooiing door Ar gas van licht geproduceerd door een z.g. normal mode robijnlaser is in overeenstemming met de waarde volgend uit de brekingsindex van Ar gas.

Dit proefschrift, appendix I.

7. De werkzame doorsnede voor aanslag van een subniveau p van het He^+ ion vanuit de He^+ grondtoestand ($1s$) door electronen is nabij de drempelenergie de helft van die welke door Burgess berekend is.

Dit proefschrift, hoofdstuk V.

8. Het verdient aanbeveling ook in Engelstalige teksten als decimaalscheider uitsluitend de "," te gebruiken in plaats van de ".". Het ware bovendien te wensen dat de aanbevelingen van de S.U.N. commissie (zoals die te vinden zijn in het Jaarboek van de Nederlandse Natuurkundige Vereniging 1970) op dit punt meer duidelijkheid boden.

9. Zowel uit oogpunt van milieuhygiëne, als van energiebesparing verdient het paard als trekker van landbouwwerktuigen de voorkeur boven de tractor.

0. De rol die Mühleisen in het "Handbuch der Physik" toekent aan ioniserende straling, aangewend om met een sonde op snelle wijze het lucht-elektisch veld te meten, is onjuist.

Mühleisen, in "Handbuch der Physik", Band XLVIII, Springer, 1957.

1. De VanderWaals capillairviscosimeter is principieel ongeschikt om een eventuele logarithmische term in de dichtheidsontwikkeling van de viscositeitscoëfficiënt van edelgassen aan te tonen.

Kawasaki en Oppenheim, Phys. Rev. A 139 (1965) 1763

Trappeniers, Botzen, van Oosten, van den Berg, Physica 30 (1964)

12. Op grond van de wet van Pascal is het aanvechtbaar dat het paleontologische reuzereptiel de Brachiosaurus (zie afbeelding) een leefwijze onder water had zoals weergegeven in de afbeelding. Eerder is het waarschijnlijk dat het beest slechts tot aan zijn rug ondergedompeld was, waarbij de lange nek dienst deed om voedsel te zoeken.

Björn Kurtén, Het tijdperk van de Dinosauriërs, Uitg. de Haan/Meulenhoff, 1968.

Edwin H. Colbert, The Dinosaur Book. Ed. American museum of natural history, New York, 1945.

13. Teneinde te voorkomen ook maar de schijn te wekken Klaas Dijkstra's "Pleidooi voor de platte aarde" au serieux te nemen, dient men iedere discussie over en verwijzing naar dit werk te voorkomen.

Klaas Dijkstra, Pleidooi voor de platte aarde, Uitg. Boersma, Enschede 1967.