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On the Applicability of the Line Reversal Method for Measuring the Electron Temperature in a Weakly Ionized Alkali Seeded Rare Gas Plasma

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ABSTRACT

When an electric current is passed through an alkali seeded rare gas plasma the electron temperature rises above the gas temperature. This effect was investigated in a streaming argon-potassium plasma at atmospheric pressure and a gas temperature of 2000°K. To measure the electron temperature the applicability of the line reversal method was investigated. General conditions under which the electron temperature equals the line reversal temperature were derived. In the case studied, the plasma had to be optically thick to the resonance radiation used. The electron temperature was measured with this method as a function of the current density. The results were compared with the values calculated from the theory and those calculated from the measured conductivities. For electron temperatures above 2400°K all three give the same results.

Introduction

The line reversal method has been used successfully for quite some time now for measuring the temperature of flames. The purpose of this investigation is to find whether this method is suitable for measuring the electron temperature in an alkali seeded rare gas, such as is used in magnetohydrodynamic energy converters.

The method is based on the observation that a spectral line seen against the continuous background of a reference radiation source appears to be emitting or absorbing, depending on the brightness of the background. If the line vanishes in the continuum and Kirchhoff's law is valid for the plasma investigated, the temperature of the plasma is equal to the black-body temperature of the reference source. This procedure was first employed by Kurlbaum 1) for the experimental determination of flame temperatures. The agreement between the reversal temperature and the gas temperature of thermal flames was demonstrated by Kohn 2) and Griffiths and Awbery 3) . This method was then extended by Rössler 4) and Moutet 5) to short-time non recurring phenomena. The line reversal method has also been used for determining temperature in shock wave experiments (Clouston, Gaydon and Glass 6)). Finally, Bundy and Strong 7) describe a method that utilizes an interferometer to determine the temperature distribution in flames of complicated structure. Before the temperatures of the various regions of the flame can be calculated from the measured reversal temperature, the optical layer thickness of the individual regions has to be known. These authors also confine their attention to flames with local thermal equilibrium.

What this paper does is investigate the applicability of the line reversal method to a plasma in which the electron temperature may be different from the gas temperature. For this purpose the theoretical principles of the method are set out and a relation established between the measured reversal temperature, the electron temperature and the gas temperature.
In addiation a method is described that allows the reversal

temperature to be determined readily and at the same time gives the optical layer thickness of the plasma in the frequency interval used. The report ends with an account of measurements on a streaming argon - potassium mixture at 2000°K and atmospheric pressure in which the electron temperature rises above the gas temperature under the influence of an electric current. This elevation is measured as a function of the electric current strength by the line reversal method. The behaviour observed is compared with the values obtained from the theory and those calculated from the measured electrical conductivity.

1. Theory

The principle of the line reversal method is as follows: The light from a continuously emitting reference lamp is passed through the plasma whose temperature is to be measured. Let the intensity of the light on entry into the plasma at the point $x = 0$ be J_ν, to which a black-body temperature T_S can be assigned so that

$$J_{\nu,x=0} = B_\nu(T_S) \equiv 2h\nu^3/c^2 \cdot \left(e^{\frac{h\nu}{kT_S}} - 1\right) \qquad (1)$$

where B_ν is the Kirchhoff-Planck function. On passing through a layer of matter this light changes in intensity in accordance with the radiation transport equation:

$$dJ_\nu/dx = \varepsilon_\nu - \kappa'_\nu J_\nu \qquad (2)$$

If the emissivity ε_ν and the absorption coefficient κ'_ν at which the induced emission processes have to be counted as negative absorption are constant along the direction of observation, the intensity emitted from the plasma at the point $x = l$ is

$$J_{\nu,x=l} = B_\nu(T_S)\, e^{-\kappa'_\nu l} + \frac{\varepsilon_\nu}{\kappa'_\nu}\left(1 - e^{-\kappa'_\nu l}\right) \qquad (3)$$

This intensity is compared with that of the reference radiation source, the light from which passes unperturbed through the plasma outside the line of observation. The temperature T_s is chosen such that the line vanishes in the continuum of the reference source, i.e. with the reversal condition

$$J_{\nu, x=l} = J_{\nu, x=0} = B_\nu(T_s) \tag{4}$$

the resulting reversal intensity is

$$B_\nu(T_s) = \frac{\varepsilon_\nu}{\kappa'_\nu} \tag{5}$$

Before anything can be said about the reversal temperature T_s it is thus necessary to investigate the ratio $\varepsilon_\nu/\kappa'_\nu$, which is related to the population ratio of the energy level of the line used. For this purpose we consider the two energy states i and k, whose population densities are n_i and n_k respectively, whose statistical weights are g_i and g_k respectively, and whose energy difference is $h\nu$. With the Einstein-Milne transition probabilities A_{ki} for spontaneous emission, B_{ik} for absorption and B_{ki} for induced emission, and the line profile $\varphi(\nu)$ normalized to 1, which is assumed to be the same for emission and absorption, ε_ν and κ'_ν take the following form:

$$\varepsilon_\nu = \frac{h\nu}{4\pi}\varphi(\nu)\, n_k A_{ki} \quad (6a) \qquad \kappa'_\nu = \frac{h\nu}{4\pi}\varphi(\nu)\left[n_i B_{ik} - n_k B_{ki}\right] \quad (6b)$$

Taking into account the relations between the transition probabilities

$$\frac{A_{ki}}{B_{ik}} = \frac{g_i}{g_k} 2h\nu^3/c^2 \quad (7a) \qquad \frac{B_{ki}}{B_{ik}} = \frac{g_i}{g_k} \tag{7b}$$

and introducing the population temperature T_p, which is defined as follows:

$$\frac{n_k}{n_i} = \frac{g_k}{g_i}\, e^{-\frac{h\nu}{kT_p}} \tag{8}$$

we obtain $\varepsilon_\nu / k_\nu' = B_\nu(T_p)$ With eq. 5 this gives in the reversal case $T_s = T_p$. The line reversal method thus measures the population temperature of the line used. The scope of this method is restricted by the condition that the emission and absorption coefficients of the plasma investigated be uniform in the direction of observation. The decisive advantage of this method is that with suitable calibration of the reference lamp its temperature in the reversal case is equal to the population temperature. Additionalinformation on optical layer thicknesses or transition probabilities is not required. The accuracy of the method is illustrated in graph form (Fig.1). The intensity J_ν observed at an arbitrary point within the line is plotted as a function of the reference intensity $B_\nu(T_s)$ The result is a straight line (cf. eq. 3) whose slope $\operatorname{tg}\alpha = e^{-k_\nu' l}$ is given by the optical layer thickness and whose value at the intensity 0 of the reference source is $\frac{\varepsilon_\nu}{k_\nu'}(1 - e^{-k_\nu' l})$. The intersection of this line with the straight line $J_\nu = B_\nu(T_s)$ yields the reversal intensity. The equality of the two intensities can only be determined, however, to within a value given by the measuring accuracy. There is thus no exact reversal point but a range, the size of which is obtained from a simple geometric relation in Fig. 1: $\Delta B_\nu = \Delta J_\nu \,/\, 1 - e^{-k_\nu' l}$.

The inaccuracy of the reversal temperature in the domain of the Wien approximation of Planck's law is found from this using the relation $\Delta B_\nu / B_\nu = \Delta T / T \cdot h\nu / kT$ to be

$$\frac{\Delta T}{T} = \frac{\Delta J_\nu}{B_\nu} \Big/ \frac{h\nu}{kT}\left(1 - e^{-k_\nu' l}\right) \qquad (9)$$

Thus, the smaller the temperature and the larger the optical layer thickness and the energy difference between the levels of the line used, the less pronounced is the relative error in the determination of the temperature. For alkali resonance lines at a temperature of 2000°K $h\nu/kT$ is of the order of 10. If in this case the error made in measuring the reversal temperature ought not to be larger than the error made in measuring the intensity ($\Delta T/T \approx \Delta J_\nu/B_\nu$), the optical layer thickness in the frequency interval used has to be at least 0.1.

The relationship between the measured population temperature and the electron (or gas) temperature has to be determined in the line reversal method just as in almost all other spectroscopic temperature measurements, and from case to case it has to be investigated.

For this purpose we consider the population ratio between the ground and resonance states of alkali atoms in a weakly ionized alkali seeded rare gas plasma in which the temperature and density of the electrons may deviate appreciably from the corresponding values of the carrier gas. To get an initial idea of the dependence of the population on the electron temperature, we first take a two-level atom model as a basis. A check will then be made later to see whether this model reproduces the present conditions in good approximation.

To calculate the population ratio one now has to put up a balance of the individual processes, which for the steady state is: The number of population processes of a level is equal to that of the depopulation processes. Allowance is made for excitation or depopulation collisions involving electrons and atoms or any impurity molecules present and for spontaneous emissions and absorption processes. The induced emission processes are ignored since the alkali resonance lines are only investigated in a temperature range around 2000°K in which $h\nu/kT \approx 10$. The population balance is then

$$n_0 \left(n_e K^e_{01} + n_A K^A_{01} + \int_{\Omega,\nu} B_{01} J_\nu \frac{\varphi(\nu)}{4\pi} d\nu d\Omega \right) = n_1 \left(n_e K^e_{10} + n_A K^A_{10} + A_{10} \right) \tag{10}$$

The probabilities K^e_{01} and K^A_{01} for electron and atom collisonal excitation respectively are connected with the corresponding coefficients K^e_{10} and K^A_{10} respectively by relations resulting from the principle of detailed equilibrium:

$$\frac{K^e_{01}}{K^e_{10}} = \frac{g_1}{g_0} e^{-\frac{h\nu}{kT_e}} \equiv b(T_e) \tag{11a}$$

$$\frac{K^A_{01}}{K^A_{10}} = \frac{g_1}{g_0} e^{-\frac{h\nu}{kT_A}} \equiv b(T_A) \tag{11b}$$

If we introduce the intensity $\overline{J}_\nu = \int_\Omega J_\nu \frac{1}{4\pi} d\Omega$ averaged over the entire solid angle Ω and characterize the radiation intensity in the line region by a temperature T_r with the definition:

$\int \bar{J}_\nu \varphi(\nu) d\nu = B_\nu(T_r)$ the number of absorption processes can also be written as follows:

$$n_0 B_{01} B_\nu(T_r) = n_0 A_{10} b(T_r) \qquad (12)$$

This results in a population ratio of

$$\frac{n_1}{n_0} \equiv b(T_p) = \frac{n_e K_{10}^e b(T_e) + n_A K_{10}^A b(T_A) + A_{10} b(T_r)}{n_e K_{10}^e + n_A K_{10}^A + A_{10}} \qquad (13)$$

We now introduce the rarefaction ratio of the radiation intensity to be integrated over the line profile to the reversal intensity, i.e. to the intensity present if there were no radiation loss:

$$V = \frac{B_\nu(T_r)}{B_\nu(T_p)} = \frac{b(T_r)}{b(T_p)} \qquad (14)$$

V is best understood in terms of the local radiation loss R in the frequency range of the line in question:

$$R = h\nu \left(n_1 A_{10} - n_0 B_{01} \int_\nu \bar{J}_\nu \varphi(\nu) d\nu \right) = h\nu\, n_1 A_{10} (1 - V) \qquad (15)$$

(1 - V) is thus the fraction of emission processes which is not compensated by absorption processes and causes radiation loss. The population ratio, therefore, can also be written in the form:

$$b(T_p) = \frac{n_e K_{10}^e b(T_e) + n_A K_{10}^A b(T_A)}{n_e K_{10}^e + n_A K_{10}^A + A_{10}(1 - V)} \qquad (16)$$

With $W = \frac{n_e K_{10}^e}{A_{10}}$, i.e. the ratio of electron collision depopulation rate to radiation depopulation rate at the resonance level, and $a = \frac{n_A K_{10}^A}{n_e K_{10}^e}$, the ratio of the atom to electron collision depopulation rates, eq. 16 becomes

$$b(T_p) = \frac{b(T_e) + a\, b(T_A)}{1 + a + \frac{1 - V}{W}} \qquad (17)$$

It can be seen that the population temperature is governed by the electron and atom collision processes only if $\frac{1-V}{W} \ll 1 + a$.

This is only the case if either

a) the collision processes exceed the radiation processes, i.e. $W \gg 1+a$, where V can take an arbitrary value, or
b) the plasma is optically thick to the investigated line so that $V \approx 1$. W may then be smaller than 1.

In order to decide whether collision or radiation processes are more numerous, we calculate the ratio W: The excitation coefficient for electron collisions is governed by the following equation:

$$K_{01}^{e} = \int_{v=\sqrt{\frac{2h\nu}{m_e}}}^{\infty} Q_{01}^{e}(v_e)\, v_e\, f(v_e)\, dv_e \qquad (18)$$

where $f(v_e)$ is the electron velocity distribution. If we put the excitation cross section $Q_{01}^{e}(v_e)$ in front of the integral as a mean constant cross section $\overline{Q}_e$ and integrate over a Maxwell velocity distribution function, we obtain

$$K_{01}^{e} = \left(\frac{h\nu}{kT_e} + 1\right)\sqrt{\frac{8kT}{\pi m_e}}\ \overline{Q_e}\ e^{-\frac{h\nu}{kT_e}} \qquad (19)$$

and with eq. 11a we then have

$$W = \frac{n_e K_{01}^{e}}{A_{10}\, b(T_e)} = \frac{g_0}{g_1}\left(\frac{h\nu}{kT_e} + 1\right)\sqrt{\frac{8kT_e}{\pi m_e}}\ \frac{n_e \overline{Q}_e}{A_{10}} \qquad (20)$$

The data used in the calculation are taken from Landolt-Börnstein 8) and Mitchell and Zemanski 19) and set out in Table 1.

Table 1

Atom	Transition	g_0/g_1	$h\nu$ [eV]	Wavelength [Å]	Mean life-time $\tau = \frac{1}{A_{10}}$ [sec]	$\overline{Q}_e$ [cm²]
Na	$3^2S - 3^2P$	2/6	2,1	5896/5890	$1{,}6 \cdot 10^{-8}$	$3 \cdot 10^{-15}$
K	$4^2S - 4^2P$	2/6	1,6	7699/7665	$2{,}7 \cdot 10^{-8}$	$\sim 10^{-15}$

The behaviour of W is shown in Fig. 2 for Na and K as a function of the electron temperature, where the values calculated with the Saha equation for a local thermal equilibrium are used for the electron density. It can be seen that at $T_e =$ 2000°K the collisional depopulation rate is still much smaller

than the spontaneous emission rate. Not until the electron temperature exceeds 3500°K is the collisional depopulation rate larger. If the population temperature in the temperature range of interest is to be governed only by the collision processes, the plasma must be optically thick to the radiation of the line used ($V \sim 1$), which is in fact the case here. Calculations of Strong and Bundy [7] and Lutz [10] show that the alkali resonance lines often used in the line reversal method are optically thick under the present conditions. With an alkali seed concentration of 10^{16} atoms/cm^3 the absorption coefficient attains the order of 10^4 cm^{-1} in the centres of the lines. Decisive in this case is the deviation of the radiation intensity that has to be integrated over the line profile from the equilibrium intensity $B_\nu(T_s)$. In an optically thick line there are practically no radiation losses in the regions of the line where $\kappa_\nu l > 1$ except for a boundary layer of thickness $1/\kappa_\nu$. In the line wings with $\kappa_\nu l < 1$ the radiation intensity does deviate appreciably from the equilibrium value, but owing to the small absorption coefficient here these regions contribute correspondingly little to the integral $\int \overline{J}_\nu \varphi(\nu) d\nu$. What is therefore of prime importance is not the radiation loss in the frequency interval used for measuring the reversal temperature but the total radiation loss within the line. If the radiation loss R is known, eq. (15) can be used to estimate the order of magnitude of 1 - V. Using the radiation loss, discussed in detail later in Section 4, it is found that 1 - V is of the order of 10^{-3}. It can thus be seen that we have here the conditions necessary for using the line reversal method in an potassium seeded argon plasma to determine a temperature governed only by the temperature of the electrons and possibly by that of the atoms as well. The radiation losses do not adversely affect the population temperature.

This still leaves the question whether exciation collisions involving rare gas atoms and any impurities present also help to determine the population temperature. Since atoms require at least twice the excitation energy for an excitation collision that electrons do, the atom collision excitation or depopulation rates are generally negligible; since, however, the electron density is here several orders of magnitude smaller than that of the atoms, they may have some influence,

and this will be estimated in the following. For this purpose we have to find the ratio a. The coefficient K_{10}^e is obtained from eqs. 11a and 19:

$$K_{10}^e = \frac{g_0}{g_1}\left(\frac{h\nu}{kT_e}+1\right)\sqrt{\frac{8kT_e}{\pi m_e}}\,\bar{Q}_e \qquad (21)$$

For the atom and molecule collisional depopulation rates we use the quenching cross section q_A, which is based on measurements of the quenching of resonance radiation in the presence of extraneous gases and is defined as follows:

$$n_A K_{10}^A = n_A q_A \sqrt{\frac{8kT_A}{\pi}\left(\frac{1}{m_A}+\frac{1}{m_K}\right)} \approx n_A q_A \sqrt{\frac{16kT_A}{\pi m_A}} \qquad (22)$$

where m_K is the mass of the alkali atom and $m_A \approx m_K$. The following expression is then obtained for a:

$$a = \frac{g_1 n_A q_A}{g_0 n_e \bar{Q}_e}\sqrt{\frac{2 m_e T_A}{m_A T_e}} \Big/ \left(\frac{h\nu}{kT_e}+1\right) \qquad (23)$$

Using the cross sections given by Landolt-Börnstein 8 and Tsuchiya 11 ($\bar{Q}_e = 3\cdot10^{-15}\,cm^2$ for electrons, $q_A \approx 10^{-19}\,cm^2$ for argon, and $q_{N_2} = 10^{-15}\,cm^2$ for nitrogen), we can estimate the conditions under which the resonance state of the alkali atoms is governed more by electron than by atom collision processes. The cross sections given apply to the Na-D lines, but the corresponding values for the potassium resonance doublet used in our measurements ought to be of the same order. On this assumption a is plotted in Fig. 3 as a function of the gas temperature for a plasma on which the measurements described later were made. This plasma consists of argon at atmospheric pressure with 0,2 % potassium, the degree of ionization in equilibrium with T_e being given by the Saha equation. Since small quantities of nitrogen impurities may be present in the experiment, a parameter for between 10^{-3} and 1 % nitrogen was included. The calculations were made for the case where the electron and gas temperatures are equal. Under the experimental conditions with a gas temperature of 2000°K and a nitrogen concentration of less than 10^{-2} %, the electron collision processes are more numerous. The population temperature and hence the reversal temperature are

equal to the electron temperature. If the electron temperature rises above the gas temperature the weight of the electron collision processes increases even more.

Let us conclude by considering to what extent the two-level atom model taken as a basis is justified. Hitherto only the transitions between the ground and resonance levels of the alkali atoms have been considered in calculating the population ratio. It will now be investigated whether the transitions to and from the higher levels can be ignored in the present case.

To make allowance for these transitions a population balance has to be made out for each individual level. Disregarding the ionization processes and atom collision transitions, the equation for the ground state is

$$n_0 \left[n_e K^e_{01} + b(T_{r1}) A_{10} + \sum_{m \geq 2} (n_e K^e_{0m} + b(T_{rm}) A_{m0}) \right] = \\ = n_1 \left[n_e K^e_{10} + A_{10} \right] + \sum_{m \geq 2} n_m (n_e K^e_{m0} + A_{m0}) \quad (24)$$

To each of the four expressions already taken into account we now have to add the sum of all the transitions between the ground state and higher excited states whose populations n_m, transition probabilities K^e_{0m}, A_{0m} etc. and whose radiation intensities are characterized by a temperature T_{rm} analogous to the radiation temperature T_r of the resonance doublet. With the present subscript m, m = 0 stands for the ground state. The relations between the transition probabilities are used to rearrange eq. 24, the following abbreviations being made:

$b_m(T_e) = \frac{g_m}{g_0} e^{-\frac{h\nu_m}{kT_e}}$ Boltzmann population ratio of the m-th level to the population of the ground state, in thermal equilibrium with T_e

$b_m(T_{pm}) = \frac{n_m}{n_0}$ Actual population ratio of the m-th level to the ground state

$\rho_m = \frac{n_m}{(n_m)_{equ.}}$ Deviation of the population of the m-th level from its population in thermal equilibrium with T_e

$V_m = \frac{b_m(T_{rm})}{b_m(T_{pm})}$ Deviation of the mean radiation intensities from the black-body radiation intensity with the appropiate population temperature T_{pm}.

Equation 24 is now resolved for $\frac{n_1}{n_0} = b_1 (T_p)$ to give

$$\frac{n_1}{n_0} = \frac{b_1(T_e)}{1 + \frac{1-V_1}{W}} \left(1 + \sum_{m \geq 2} \frac{b_m(T_e)}{b_1(T_e)} \left[\left(1 - \frac{\varrho_m}{\varrho_0}\right) \frac{K^e_{m0}}{K^e_{10}} - (1-V_m) \frac{\varrho_m A_{m0}}{\varrho_0 A_{10} W} \right] \right) \tag{25}$$

From calculations on hydrogen and hydrogen-like atoms made by Mc Whirter et al. 12 we know that it generally holds that

$$\varrho_0 = \frac{n_0}{(n_0)_{equ.}} \geq \frac{n_m}{(n_m)_{equ}} = \varrho_m \qquad \text{or} \qquad 0 < \frac{\varrho_m}{\varrho_0} \leq 1$$

In addition, as the energy level increases both the excitation cross sections and the transition probabilities decrease, and so

$$\left(1 - \frac{\varrho_m}{\varrho_0}\right) \frac{K^e_{m0}}{K^e_{10}} < \frac{K^e_{m0}}{K^e_{10}} < 1 \qquad \text{and} \qquad (1-V_m) \frac{\varrho_m A_{m0}}{\varrho_0 A_{10} W} < \frac{A_{m0}}{A_{10} W} < \frac{A_{20}}{A_{10} W} \approx 1$$

The whole correction sum in eq. 25 can thus be estimated as follows: $\sum_{m \geq 2} \ldots < \sum_{m \geq 2} b_m(T_e) / b_1(T_e)$

Owing to the large energy differences between the higher levels and the resonance level in comparison with the difference between the latter and the ground state of the alkali atoms and owing to the relatively low temperature of our plasma, however, the population of the higher levels is very small compared with that of the resonance level. For the potassium resonance series at $T_e = 2000^{\circ}K$, for example, we obtain

$$\sum_{m \geq 2} b_m(T_e) / b_1(T_e) = (2.5 + 0.12 + 0.02 + \ldots)\, 10^{-4}$$

The sum expression in eq. 25 is thus small relative to 1. Making allowance for the higher levels thus represents in our case only a very slight correction of the population balance of the two-level atom model.

The results can be summarized as follows: The reversal intensity is

$$B_\nu(T_s) = \frac{\varepsilon_\nu}{\kappa_\nu} = B_\nu(T_F) = \frac{B_\nu(T_e) + a\, B_\nu(T_A)}{1 + a + (1-V)/W}$$

Under the conditions investigated we get $a < 10^{-1}$, $1-V < 10^{-3}$ and $W > 10^{-2}$; since $B_\nu(T_A) \leq B_\nu(T_e)$ is also valid, $B_\nu(T_s)$ should be equal to $B_\nu(T_e)$ in good approximation and therefore the reversal temperature should be equal to the electron temperature.

2. Experimental set-up

The experimental set-up is shown schematically in Fig.4. The argon is heated by two arc heaters, facing one another and projecting into a mixing chamber. Here the argon is seeded with potassium. The mixture then flows through a connecting pipe into the actual measuring section and is expelled. The gas is conveyed in tantalum or tungsten pipes which are thermally and electrically insulated by a boron nitride sheath embedded in magnesium oxide. The whole system is assembled in a vacuum-sealed stainless steel casing. In the centre of the mixing chamber the potassium is injected in vapour form from a tungsten tube heated by the plasma. The amount of potassium injected can be regulated exactly by means of a piston with a continuously variable feed rate. In order to neutralize fluctuations in the amount of potassium injected due to irregularities in the feed rate and boiling delays, the tungsten pipe was tipped with a nozzle with a cross section of 0.15 mm^2. This resulted in the formation of a potassium vapour atmosphere under pressure in the hot top part of the pipe, thus providing a buffer volume.

The measuring channel (Fig. 5) has a rectangular cross section of 1 x 1.3 cm^2 and is 10 cm long. The walls are made of boron nitride. An electric current can be passed through this section in the direction of flow. Current is supplied by way of tantalum electrodes projecting into the channel at right angles. The potential distribution along the channel can be measured with tungsten probes. Two apertures opposite one another in the walls of the channel allow optical observations and measurements to be made perpendicular to the direction of flow.

The gas temperature was checked by means of a tungsten-rhenium thermocouple. Another way was to measure the electron temperature by the line reversal method with no current flowing in the plasma, where $T_e \approx T_{gas}$. This reversal temperature was about 100°K higher than that measured with the thermocouple.

Since, however, the temperature of the thermocouple is reduced by thermal conduction and radiation losses, it is easy to see why this systematic error occurs. When the thermocouple was shifted it was found that the temperature was almost constant over the channel cross section. Only in a boundary layer about 1 mm thick could a fairly pronounced temperature decrease be observed. For all measurements the conditions were as follows:

Argon mass flow:	14,5 g/sec	Channel cross section:	1.3 cm^2
Gas temperature:	2000 °K	Mean velocity:	$4 \cdot 10^4$ cm/sec
Gas pressure:	1.1 atm	Potassium seed:	0,026 g/sec (0,175%)

3. Experimental application of the line reversal method

3.1 Measuring set-up

The typical set-up used in the reversal method is shown in Fig. 6. Light from a tungsten strip lamp is focused along the axis of the measuring channel. Together with the light from the plasma it is imaged at the same solid angle on the entrance aperture of a light pipe, this opening being matched to the plasma cross section. It is then focused on the entrance slit of a Zeiss MG 4 monochromator by the exit of the light pipe, which is slit-shaped. The spectrally resolved light is recorded by a RCA 7102 photomultiplier. The wavelength setting of the monochromator can be rotated continuously through the spectral region of interest by means of a variable speed motor. In this way the intensity of the line can be recorded in emission and absorption as a function of the wavelength by means of an oscilloscope (see Fig. 6) or - after appropriate amplification - by means of an X-Y recorder (see Fig. 7). Using the relationships described in Fig. 1 we can determine the reversal temperature quickly and accurately as follows: The intensity J_ν recorded by the photomultiplier is plotted as a function of the tungsten strip lamp intensity $B_\nu (T_s)$, which is measured direct by means of an extra photoelectric cell; this is done twice, first with a fixed wavelength setting of the monochromator outside of the line and

then within the line. The intersection of the two lines obtained in this way yields the reversal intensity of the tungsten strip lamp (see Fig. 8). The B_ν (T_s)-axis can be calibrated direct in temperatures so that the reversal temperature can be read off immediately. The advantage of this "intersection reversal method" is primarily that the reversal temperature can be measured relatively fast and without undue effort. The time required is governed by how quickly the intensity of the reference light source is varied: Using a Moseley X-Y recorder we took 1 - 2 sec. This time can, however, be shortened several orders of magnitude by recording with an oscilloscope and using a rotating sector or similar device with which to vary the intensity in a defined manner. Another adventage of this method is that - provided the resolution of the monochromator is sufficient - the optical layer thickness of the plasma in the frequency interval used is also easy to determine (cf. Fig. 1).

3.2 Influence of the optical layer thickness

It was shown in Chapter 2 that the reversal temperature is only equal to the electron temperature if the electron collision depoulation rate in the resonance level is much higher than the difference between the spontaneous emission rate and absorption rate for the resonance radiation. In our potassium seeded argon plasma the number of electron collision processes exceeds the emission rate only at electron temperature above 3000°K. In the region of interest here the plasma must therefore be optically thick to the potassium resonance lines. For this reason the optical layer thickness of the plasma had to be checked for the potassium seed concentration used ($n_K = 7 \times 10^{15} cm^{-3}$). In the centres of the lines of the potassium resonance doublet it was found by the "intersection method" to be larger than 2. Owing to the limited resolution of the monochromator, however, frequency ranges of the line to which the plasma is optically thin were also involved here, and so the optical layer thickness determined is much lower than the actual thickness in the centres of the lines. For small optical layer thickness one would expect the measured

reversal temperature at $T_e < 3000^{\circ}K$ to be lower than the electron temperature. This was checked experimentally by varying the thickness for the lines to be investigated. For this purpose a small amount of Na (1 %, i.e. $n_{Na} = 7 \times 10^{13} cm^{-3}$) was added to the potassium. The reversal temperature was then $100^{\circ}K$ lower than the temperature measured on the potassium resonance lines. If the sodium number density was raised, the Na lines had the same reversal temperature as the K lines.

3.3 Flushing and the influence of boundary layers

To prevent the potassium from diffusing through the openings in the channel wall and being deposited on the windows in the steel casing of the channel, pure argon is blown in through the openings. This flushing also breaches the highly absorbent, relatively cold boundary layer of the plasma. Fig. 9 shows clearly how, with little flushing at first, the lines of the potassium resonance doublet are almost completely self-absorbed. Only the radiation in the line wings makes its way to the outside because it can no longer be absorbed by the colder boundary layers. With intensified flushing the reversal temperature also rises quickly from values which are much too small at first. It then attains a value which is constant within the measuring accuracy and which cannot be raised once the argon mass flow through the two windows (each 2 mm^2 in area) reaches 1.5 g/sec, even if this flow rate is substantially increased. But a certain self-reversal of the lines does persist. There is apparently still a slight temperature profile present. This can be seen in the determination of the reversal temperature by the fact that the reversal condition for the various regions of the line is no longer satisfied for a uniform reference temperature, as is expected in the case of a homogeneous plasma.For this purpose we consider at a resolution of about 8 Å the potassium resonance doublet in emission, absorption and reversal for various intensities of the reference source (Fig. 7). These photometric curves were evaluated graphically. The temperatures obtained are plotted in Fig. 10 as a function

of the wavelength. The line wings thus yield temperatures up to 100°K higher than those in the line centres. As the line centres are practically unaffected by the absorption in the boundary layers, it is to be expected that their reversal temperature corresponds to the temperature in the inner regions of the channel, while the reversal temperature of the line centres is reduced in keeping with the temperature profile as a result of self-absorption.

4. Measurement of the electron temperature elevation

When a current flows through the potassium seeded argon plasma the electron temperature rises above the gas temperature. In order to measure the dependence of the electron temperature on the electric current density, an electric current was passed through the measuring channel in the direction of flow. The current was supplied by way of two pairs of electrodes, each in parallel connection (cf. Fig. 5), so as to achieve as uniform a current density distribution as possible. The measurements were made in that part of the discharge where there were no electrode effects. This was checked by measuring the potential distribution between the two pairs of electrodes with static probes. In calculating the mean current density $j = I/A$ allowance was made for a non-conducting boundary layer 1 mm thick, found from the measurements of the temperature profile. Curves calculated for boundary layers 0.5 and 1.5 mm thick are also given in Fig.11 for the sake of comparison. Since only a relatively rough mean value of the current density can be obtained to the unfavourable ratio between the surface and volume of the channel, only the mean reversal temperature between the line wings and the centre were plotted in Fig.11.

Besides measuring the electron temperature elevation by the line reversal method, we determined the electrical conductivity as a function of the current density. It was calculated with Ohm's law $j = \sigma E$ from the measured electric field and

the mean current density. The results agree with those of Zukoski, Cool and Gibson [13], who determined the conductivity under the same conditions as the authors here. In order to compare these measurements with the behaviour of the electron temperature as a function of the electric current density, the corresponding electron temperature was calculated from the measured conductivity using the Saha equation. This is also given in Fig. 11.

The agreement between the two different methods of determining the electron temperature is very good. Only in the current density range below 0.5 A/cm^2 does the temperature calculated from σ exceed that determined from the reversal temperature. This is probably because the electron density in this region has a higher value than is compatible with the local electron temperature according to the Saha equation. This electron density is due to the charged particles, which are dragged by the flow from the hotter plasma heating system into the measuring channel as a result of finite recombination rates.

The measured electron temperature also agrees well with the theoretical distribution. This is obtained from the energy balance of the electrons, which is

$$\frac{j^2}{\sigma} = n_e \nu_e \frac{m_e}{m_A} 3k(T_e - T_A) + R(T_e) \qquad (26)$$

where the thermal conduction losses are ignored and it is assumed that the electron density n_e in equilibrium with T_e is given by the Saha equation. j^2/σ is the energy input to the electrons per sec and cm^3; the first term on the right-hand side stands for the energy loss due to elastic collisions with atoms (ν_e : electron collision frequency, T_e and T_A: electron and atom temperatures respectively, a Maxwellian velocity distribution being assumed for each, and m_e/m_A: mass ratio); the second term R (T_e) represents the energy lost by the electrons in inelastic collisions, i.e. lost in the form of radiation.

Zukoski, Cool and Gibson [13] made allowance for the radiation loss by calculating for various potassium lines the total radiation flux emitted from the plasma and averaging the total radiation loss over the entire volume. For the probabilities of a photon breaking out without reabsorption these authors use values calculated by Holstein [14] in the case of a uniform, but non-stationary decaying plasma for various geometries and broadening mechanisms.

The results of this investigation now suggest a more simple estimate of the radiation loss using experimentally determined values. The reasons are as follows: The major part of the radiation losses occurs in the region of the potassium resonance doublet since the radiative transitions between the resonance and ground states far outnumber all others. Attention can therefore be confined to these lines. Moreover, the plasma is optically thick in the centre of the lines, where it practically attains the intensity of a black body at the electron temperature. In view of the large optical thickness, isotropy may also be assumed. The radiation output emitted from the plasma surface per unit area is therefore

$$\frac{dE}{dt\, dA} = \pi \int_L J_\lambda \, d\lambda$$

The integral of the intensity over the line region can be replaced by a rectangle of suitable width $\Delta\lambda$ which is equal in height to the maximum intensity:

$$\int_L J_\lambda \, d\lambda = J_{\lambda,max} \, \Delta\lambda \approx B_\lambda(T_e) \, \Delta\lambda$$

The required width $\Delta\lambda$ is estimated from the experimentally determined behaviour of the intensity as a function of the wavelength (Fig. 7) to be about $\Delta\lambda = 40$ Å altogether for $T_e = 2000\ ^\circ K$ at $j = 0$. Like Zukoski [13], we use as local radiation loss the total radiation loss averaged over the emitting volume:

$$R(T_e) \approx \frac{A}{V} \pi \, B_\lambda(T_e) \, \Delta\lambda \qquad (27)$$

where A/V is the ratio between the surface area and volume of the plasma. The value of $\Delta\lambda$ varies only slightly with increasing electron temperature; moreover, as the electron temperature increases the radiation losses become smaller rela-

tive to the elastic losses anyhow. The constant value of 40 Å given for the two lines of the potassium resonance doublet could therefore be used for the calculation. It can be seen in Fig. 11 that the electron temperature elevation calculated in this way agrees well with the results of both the line reversal method and the conductivity measurement.

This method of calculating the electron temperature fails, however, at current densities below 0.1 A/cm^2. Without current in the plasma the electron temperature would have to drop below the gas temperature until the radiation losses are compensated by the energy gained by the electrons by virtue of the temperature difference $T_A - T_e$ resulting from elastic collisions with the heavy particles. No such drop could, however, be detected. The reason may again be that the strong heating of the gas by the arc heaters makes the electron density in the measuring channel higher than is compatible with the local electron temperature. The electron density then cannot drop below a certain value with decreasing electron temperature. If this value is assumed to be about $10^{13} cm^{-3}$, as suggested by the conductivity measurement at low current densities, the decrease in the electron temperature resulting for $j = 0$ from the above energy balance is only $100^{\circ}K$. A strong reduction of the electron temperature in the case $j = 0$ may also be prevented because the above calculation overestimates the radiation losses at low electron temperatures. That is excitation by atom collisions gains in importance with decreasing electron temperature and hence the radiation loss no longer occurs wholly at the expense of the electron energy. This complex of problems is treated in detail by Schwenn [15].

5. Influence of nitrogen impurities

As seeding with molecular gases greatly reduces the degree of the electron temperature elevation, it was also investigated what influence nitrogen impurities have on this elevation mechanism. The electron temperature elevation in a rare gas

depends for its effectiveness on the energy gained by the electrons from the electric field being lost slowly in essentially elastic collisions with rare gas atoms; for the mean thermal energy of the electrons (0.2 eV) is much lower than the lowest excitation energy of argon (11.5 eV). In molecular gases, on the other hand, inelastic collisions have also to be taken into account because of the possibility of oscillations and rotations being excited. For this purpose it is usual to introduce into the energy balance a formal energy loss factor δ which states by how much the mean energy loss of an electron in a collision with a molecule is higher than that in a purely elastic collision.

With nitrogen present the values measured for the reversal temperature were apparently between those of the electron temperature and gas temperature owing to the possibility of the nitrogen molecules exciting the alkali atoms, this being in agreement with what was said in Section 1. We therefore discuss here only the measurements of the electrical conductivity as a function of the mean current density for various nitrogen concentrations (Fig. 12). The expected decrease of the electrical conductivity with increasing nitrogen concentration was obtained. Butthis drop is larger, particularly for small nitrogen concentrations, than that calculated with an energy loss factor of $\delta = 13$, as was given by Anderson and Goldstein [16]. It should be noted, however, that the factors given in the literature differ appreciably. It was not possible in our case anyhow to determine an energy loss factor for approximating the experimental values that was effective for all current densities and nitrogen concentrations. Our measurements show, at any rate, that the mean energy loss of an electron in a collision with a nitrogen molecule is higher than that usually given in the literature, and how important it is to work with a very pure rare gas if the electron temperature elevation is to remain effective.

Conclusion

Our calculations and measurements show that it is possible in a non-equilibrium plasma, such as is used in rare gas MHD

generators to determine the electron temperature by the line reversal method. At low current densities and correspondingly low electron densities below about 10^{14} per cm^3 the plasma has to be optically thick to the alkali resonance lines used before local thermal equilibrium can set in. There thus arises the problem of colder boundary layers making the measurement unreliable, but the perturbing influence of these layers can be greatly reduced by having clean window flushing free of absorbing atoms. Another possibility is to determine the reversal temperature in the line wings, these being scarcely affected by the boundary layers. On the other hand, the very fact that the plasma is optically thick to the potassium resonance lines and the radiation here attains approximately black-body radiation intensity at the electron temperature allows the radiation loss to be estimated readily after the spectral intensity distribution is found experimentally. The measurements showed that the behaviour of the electron temperature determined by the line reversal method as a function of the mean electric current density is in good agreement with the values calculated from the conductivity values and the theory.

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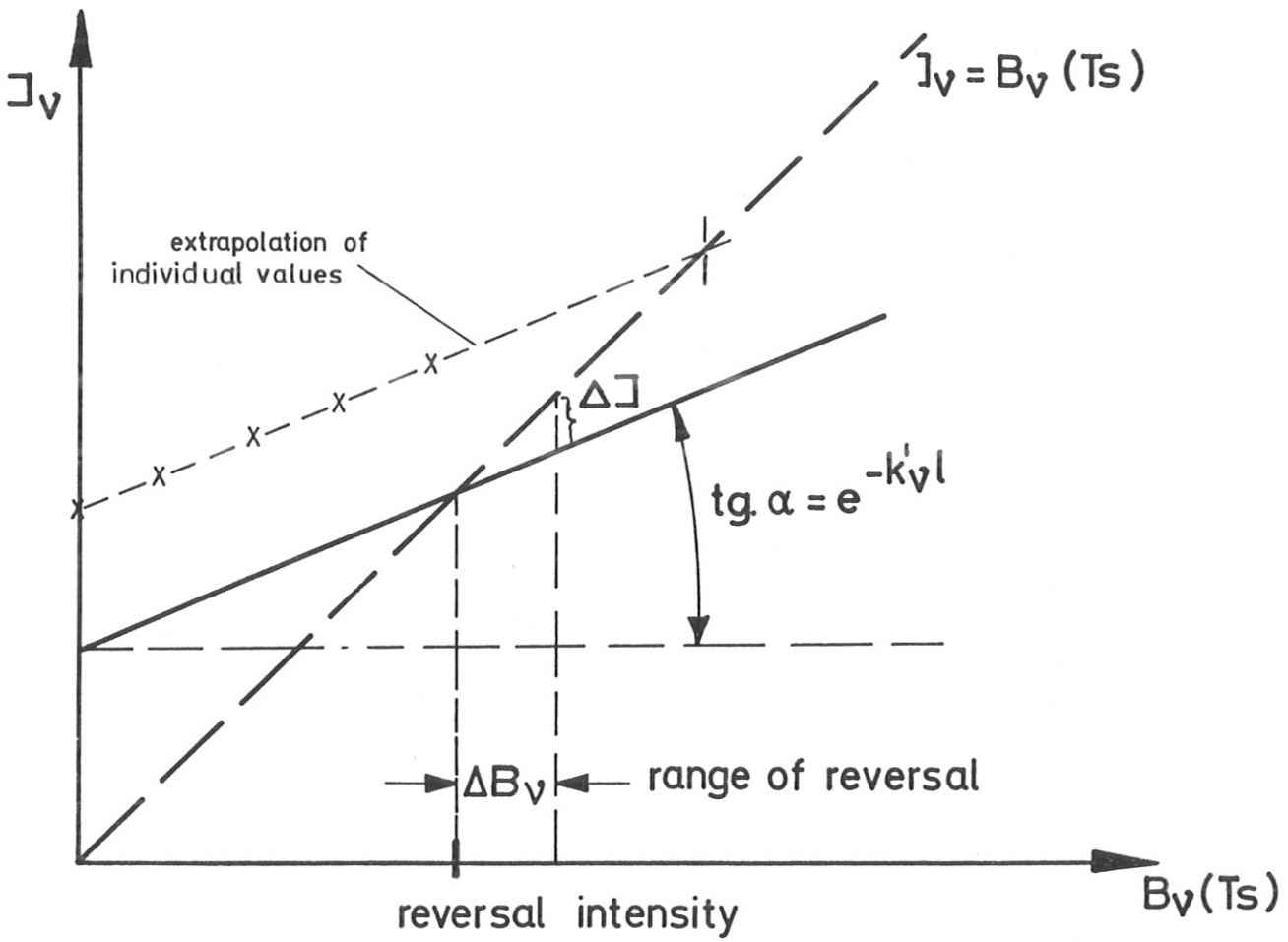

Fig. 1 Principle of the line reversal method.

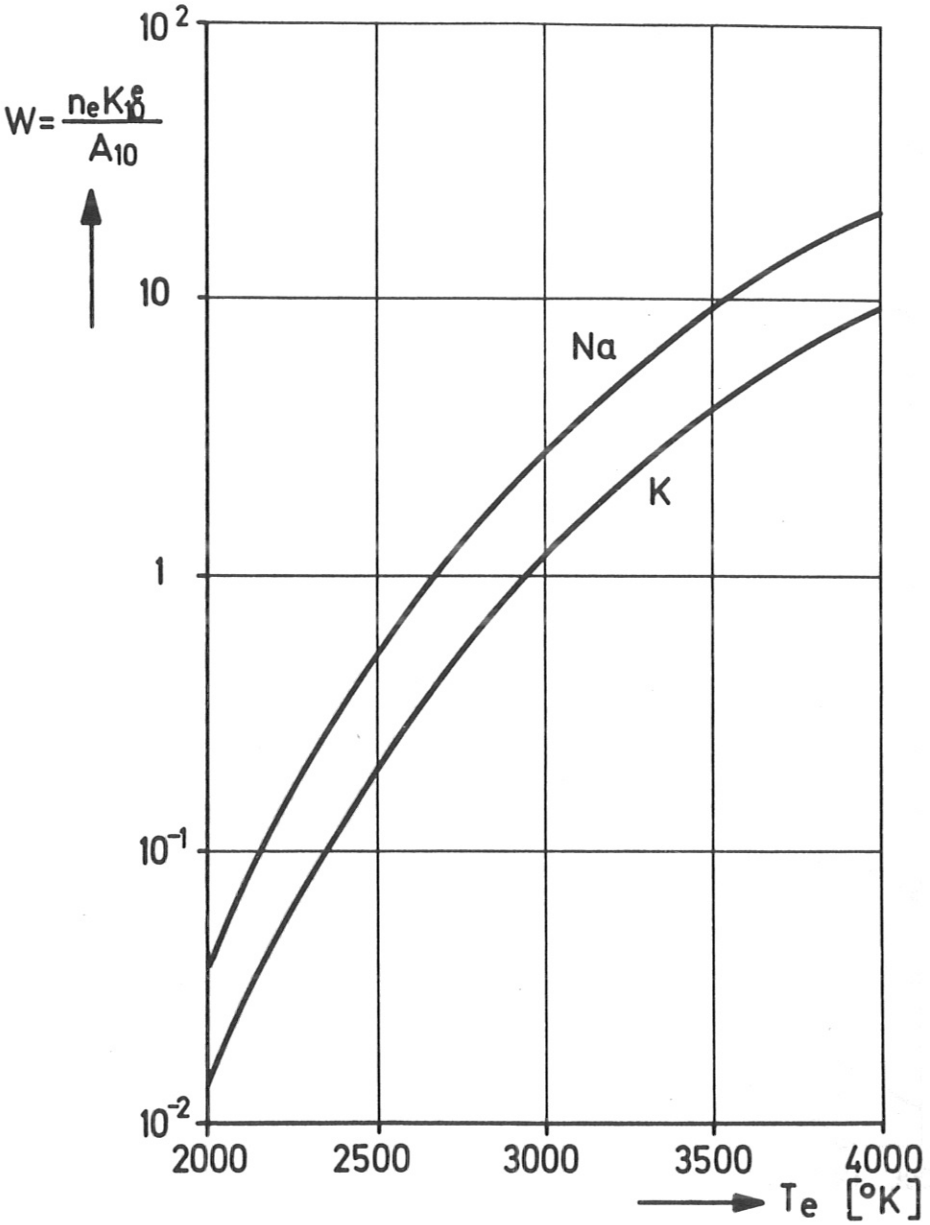

Fig. 2 Ratio of the collisional and radiation depopulation rates at the resonance level as a function of the electron temperature .

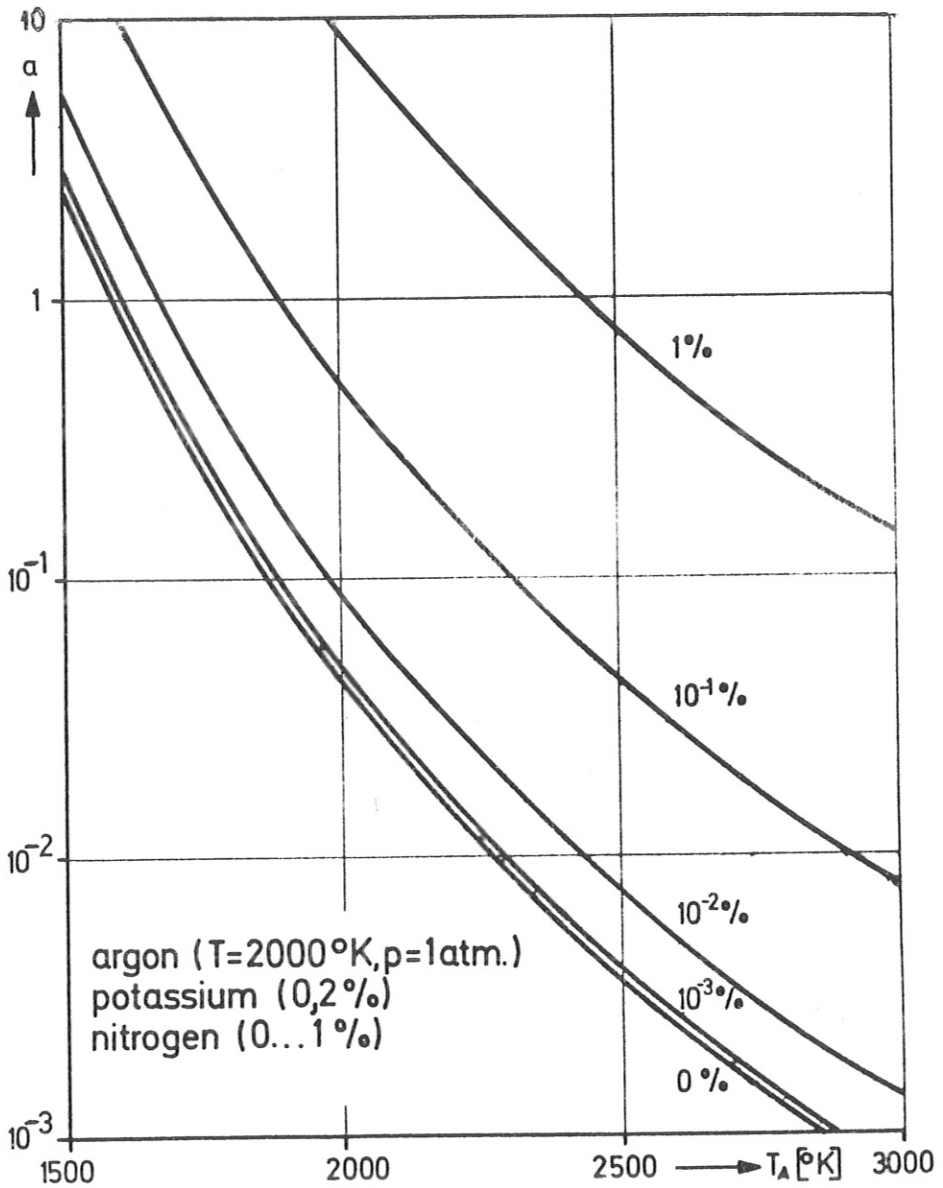

Fig. 3 Ratio of the resonance level depopulation rates induced by neutral and electron collisions as a function of gas temperature for several N_2 concentrations.

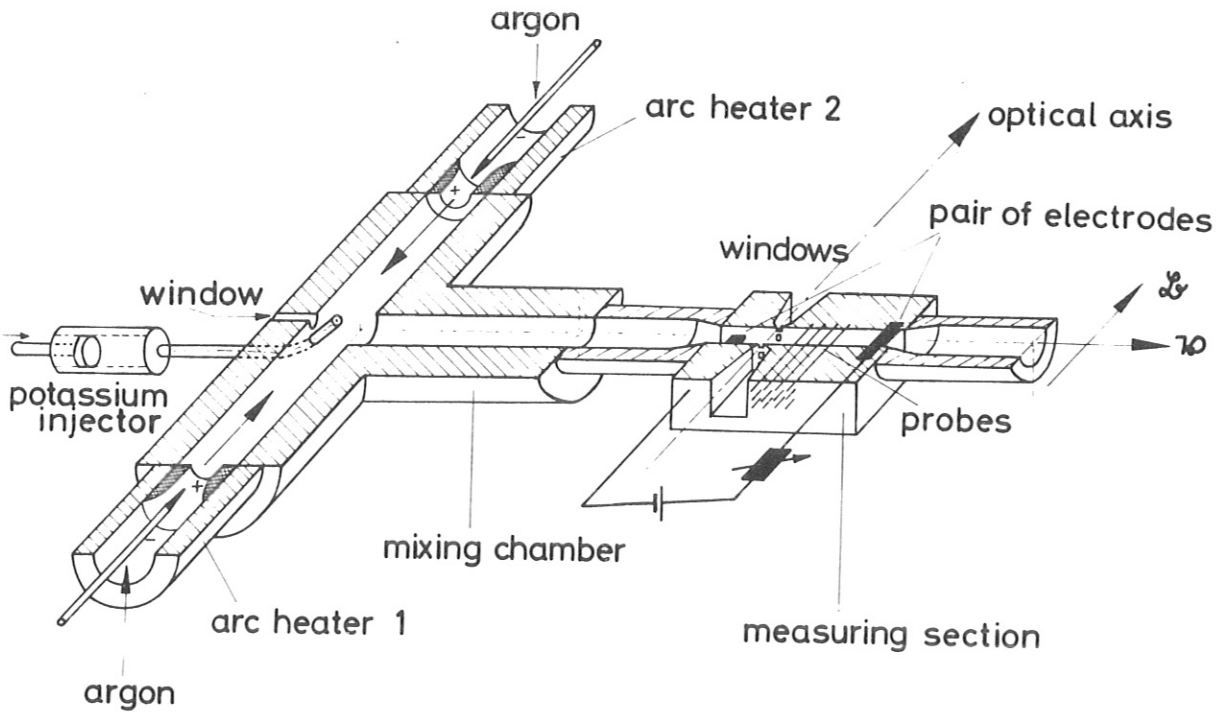

Fig. 4 Experimental equipment.

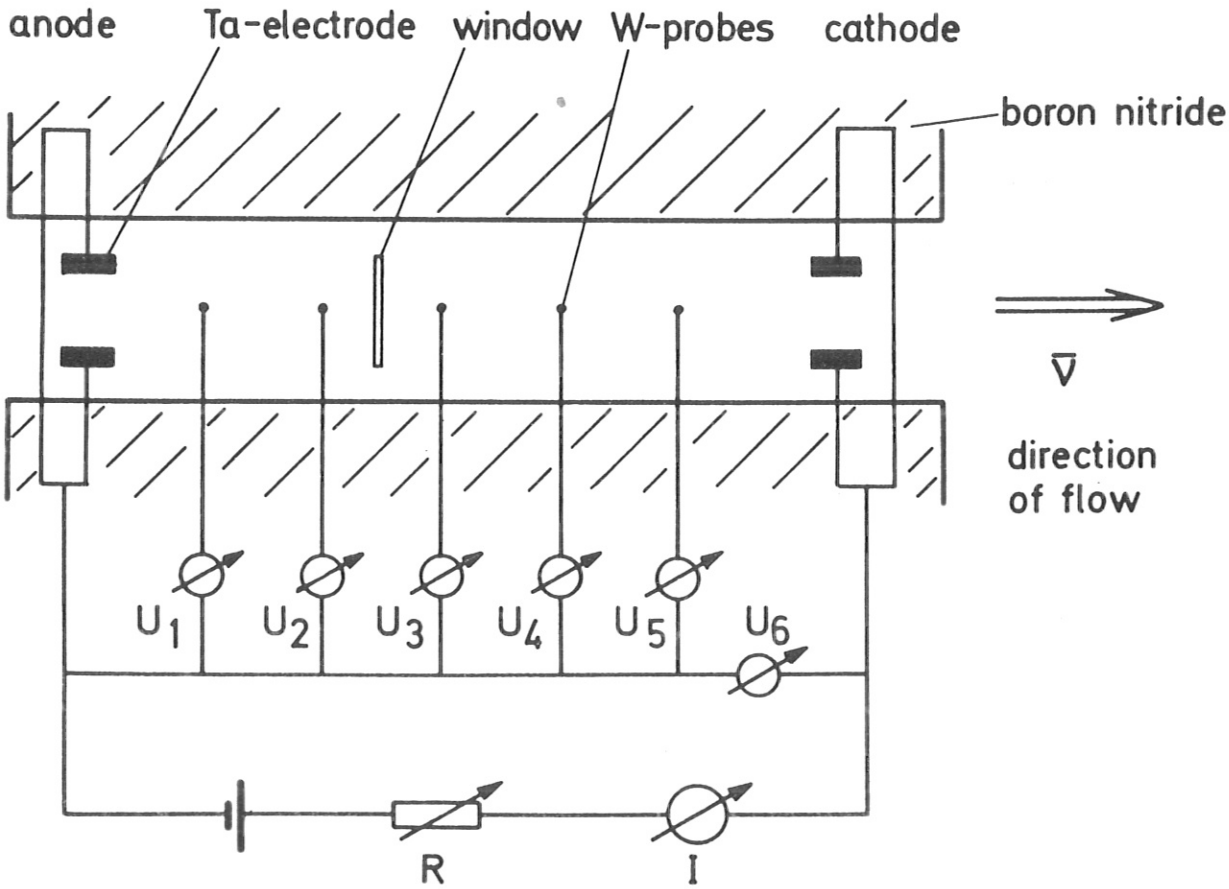

Fig. 5 Test section.

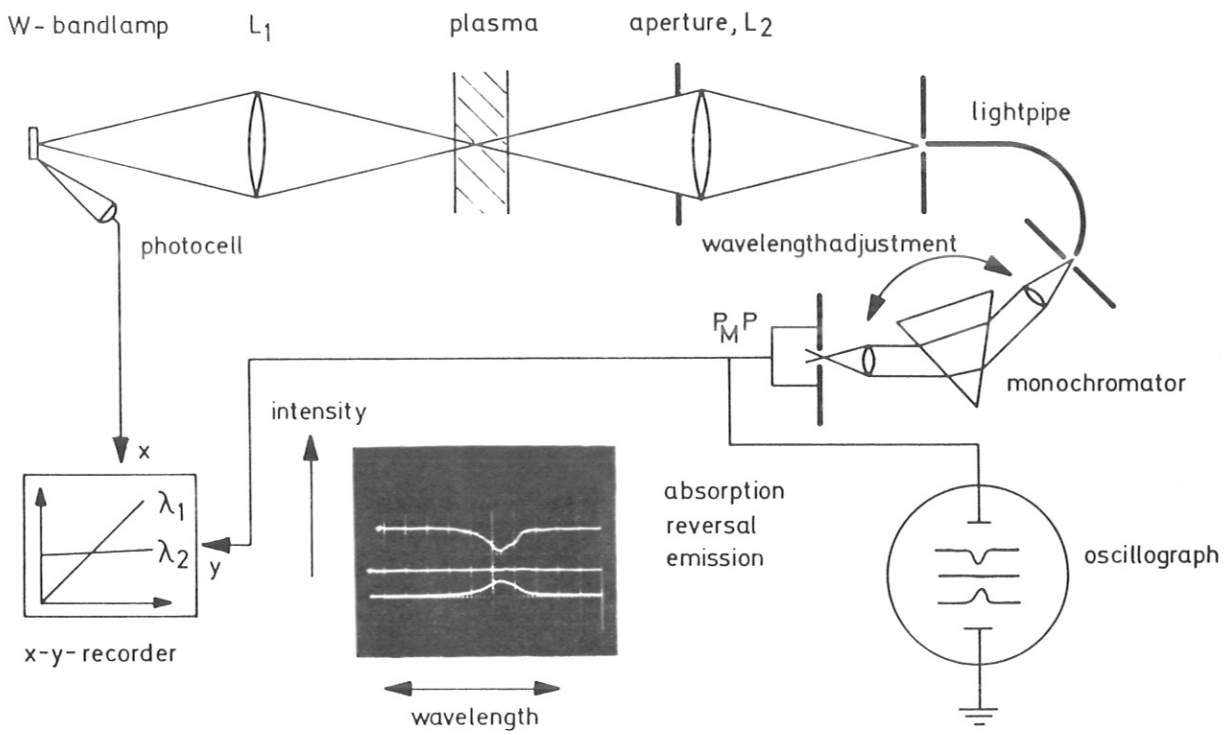

Fig. 6 Line reversal apparatus.

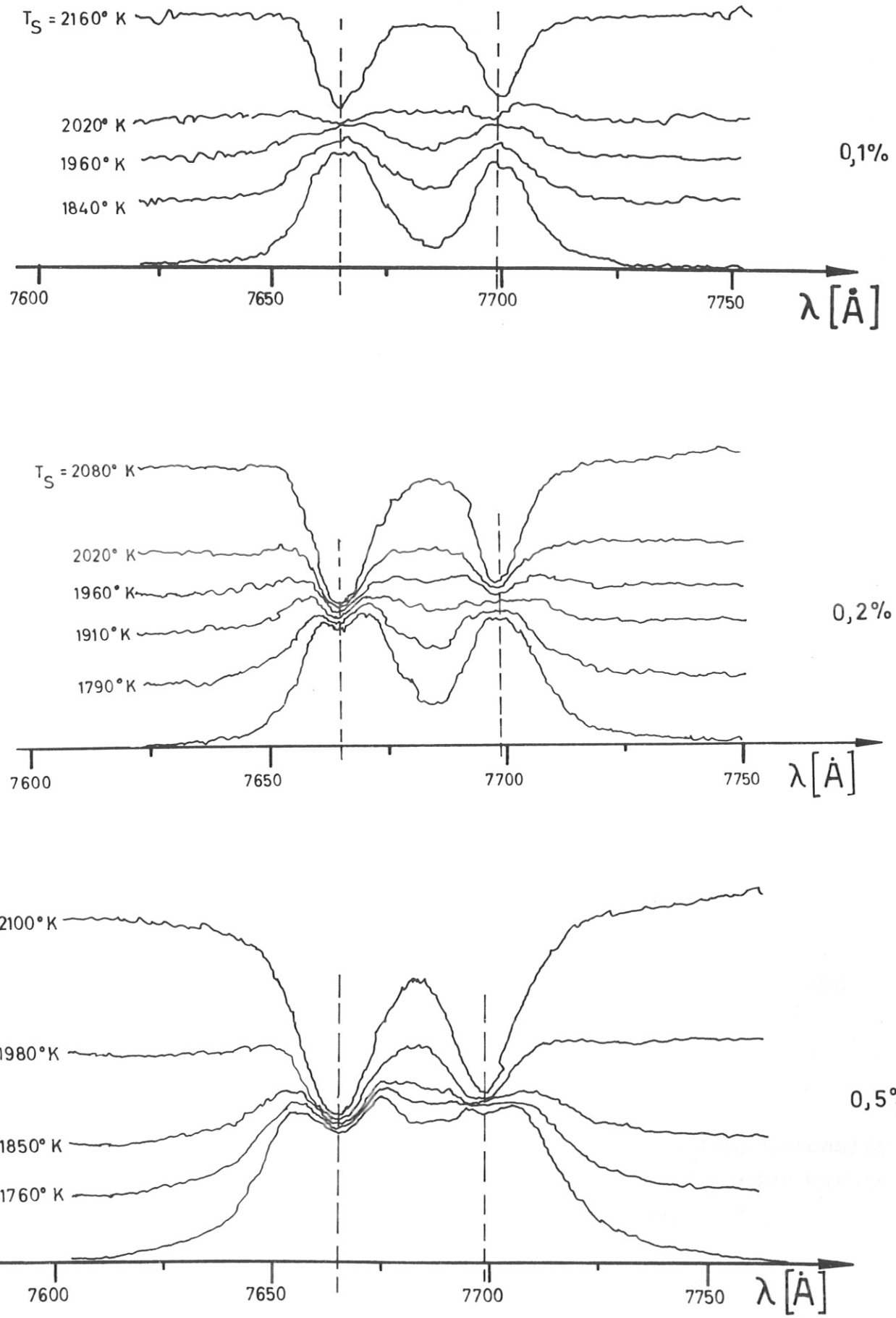

Fig. 7 Potassium resonance lines at different intensities of the reference light source for three seed concentrations.

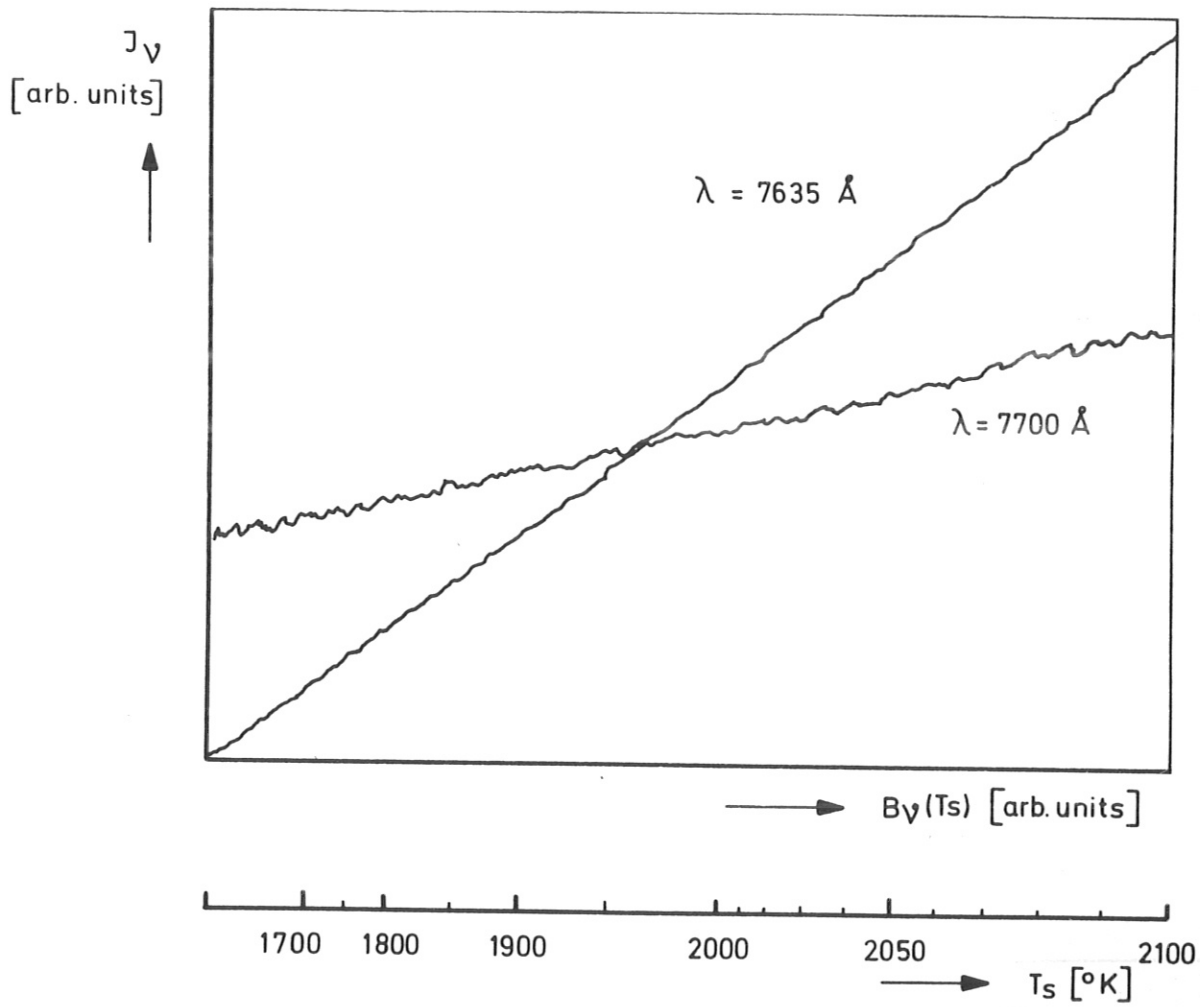

Fig. 8 PMP signal intensity as a function of W-band lamp intensity measured by a photocell and plotted on a X-Y recorder within and outside of the potassium resonance line.

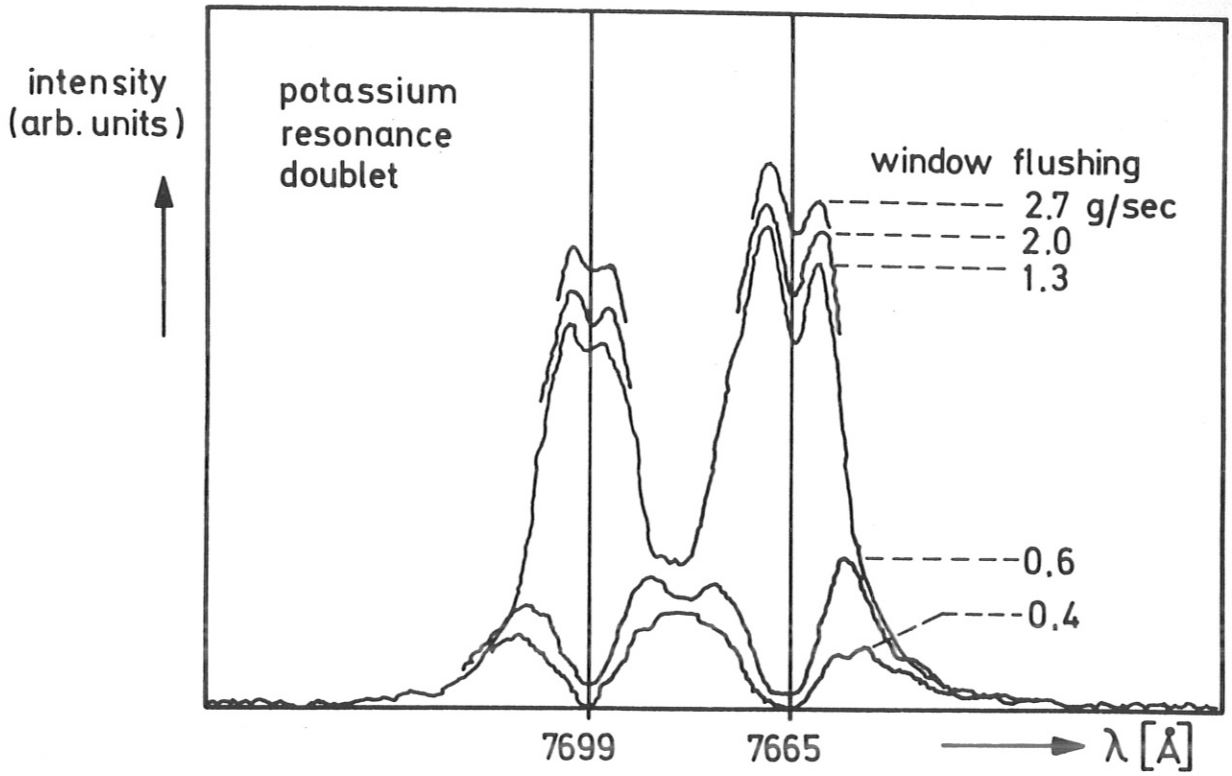

Fig. 9 The potassium resonance lines at various bleed-in rates.

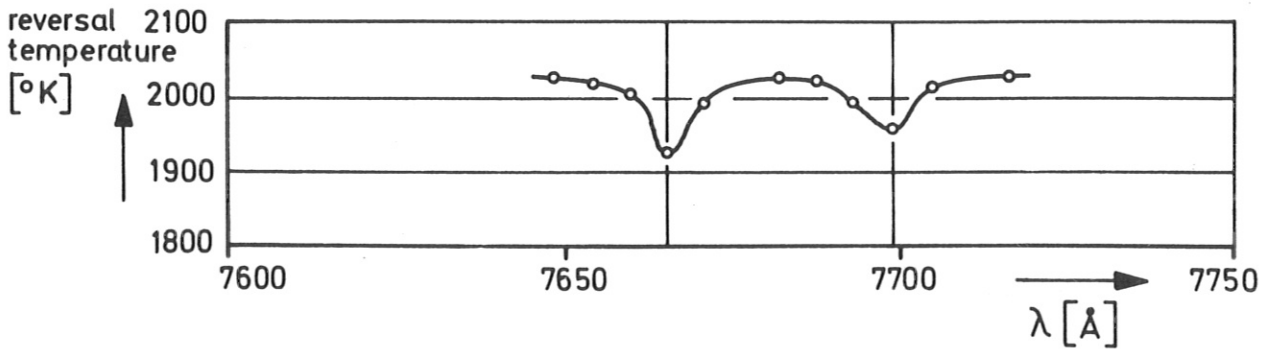

Fig. 10 Line reversal temperature as a function of the wavelength.

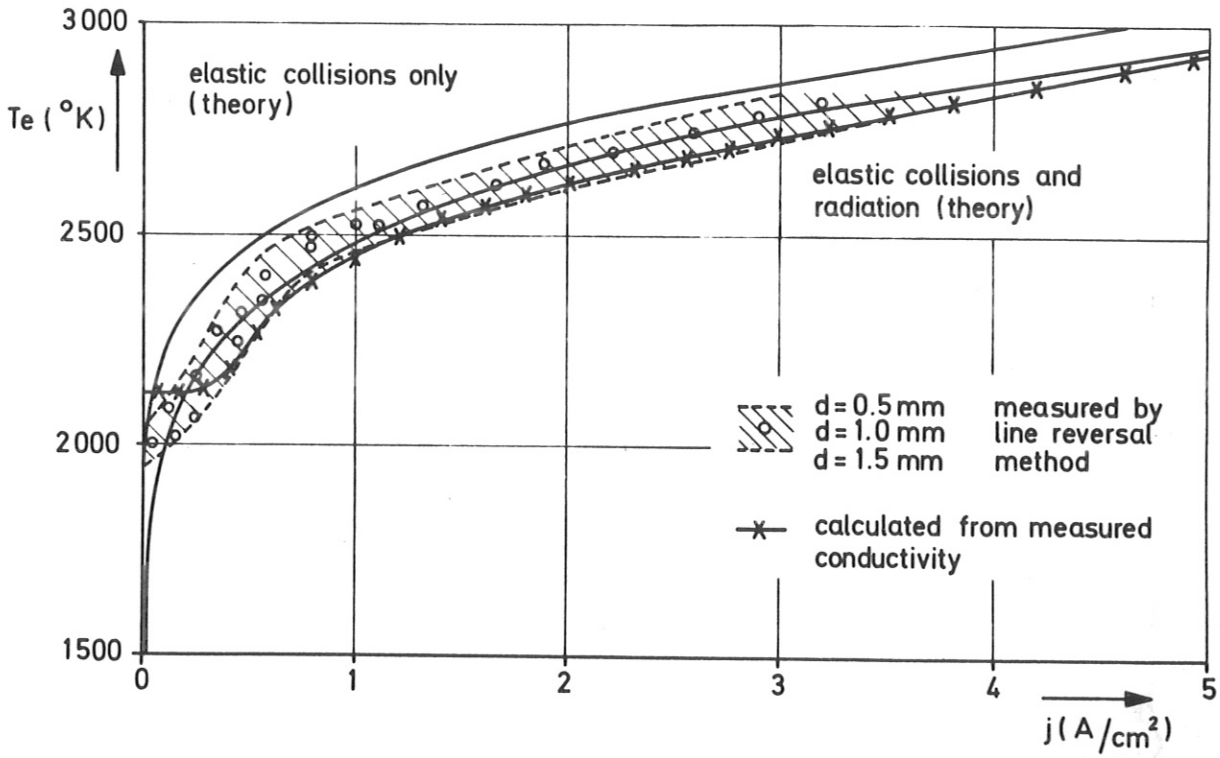

Fig. 11 Electron temperature as a function of the average current density.

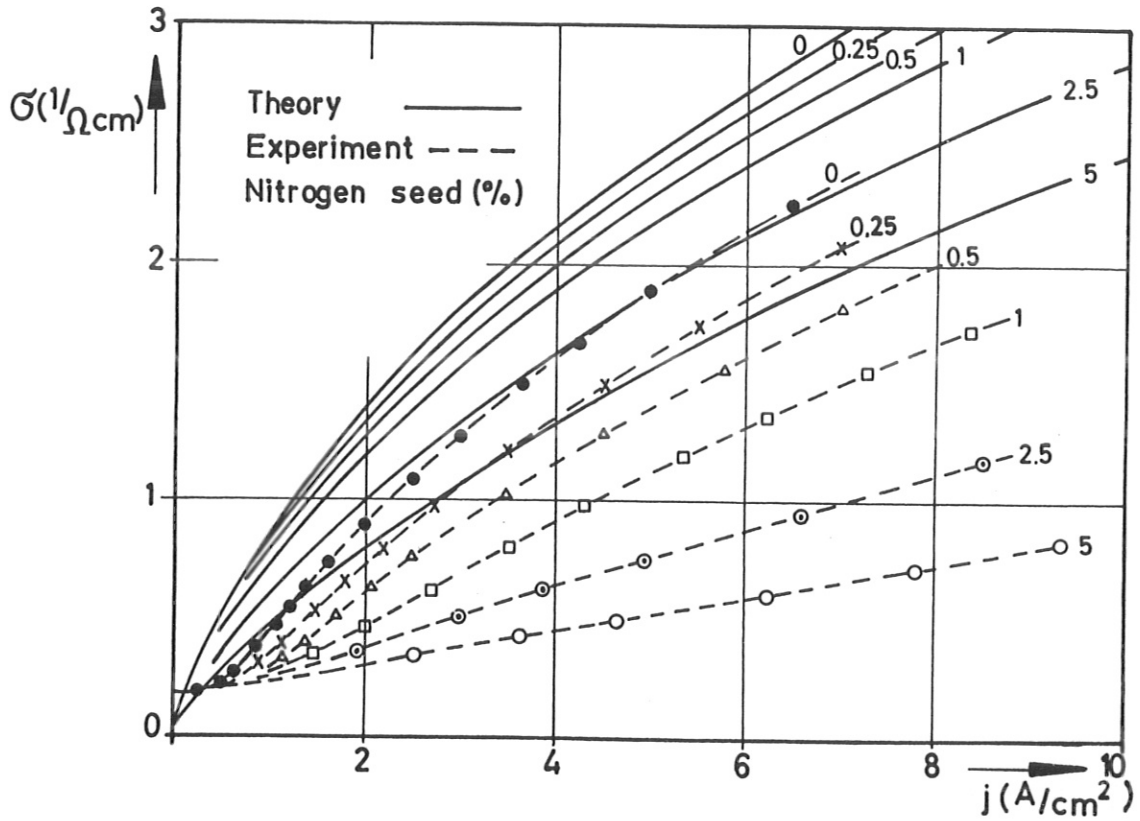

Fig. 12 Electrical conductivity as a function of average current density at various nitrogen concentrations