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Citation for published version (APA):

Kroesen, G. M. W., Schram, D. C., Wilbers, A. T. M., & Meeusen, G. J. (1991). Analysis of the evolution of a supersonically expanding plasma. *Contributions to Plasma Physics*, 31(1), 27-42.

Document status and date:

Published: 01/01/1991

Document Version:

Publisher's PDF, also known as Version of Record (includes final page, issue and volume numbers)

Please check the document version of this publication:

- A submitted manuscript is the version of the article upon submission and before peer-review. There can be important differences between the submitted version and the official published version of record. People interested in the research are advised to contact the author for the final version of the publication, or visit the DOI to the publisher's website.
- The final author version and the galley proof are versions of the publication after peer review.
- The final published version features the final layout of the paper including the volume, issue and page numbers.

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Analysis of the Evolution of a Supersonically Expanding Plasma

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Abstract

A supersonically expanding arc plasma in argon is analyzed both experimentally and theoretically. The plasma is created in a cascaded arc and extracted through a hole in the anode. It emanates in a large vacuum system, where it expands supersonically. This expansion is limited by a shock wave. After the shock wave a subsonic plasma beam is created. A quasi one-dimensional model, based on the conservation of mass, momentum and energy is presented. The shock wave is treated as a discontinuity. The electron density, the gas velocity and the gas temperature are measured as a function of the position in the expansion by means of Stark broadening and Doppler spectroscopy. The model calculations agree well with the measurements, especially in the first part of the supersonic flow.

1. Introduction

A supersonically expanding arc plasma in argon was recently reported as a source of excitation, dissociation and ionization to be used for the deposition of thin films [1]. The reactor, which is outlined in Fig. 1, consists of a cascaded arc which is connected to a vacuum system. The arc plasma expands supersonically, and if a reactive gas like methane is injected, it will be dissociated completely. The carbon atoms are transported

Fig. 1. Schematic view of the reactor used in the reported experiments. When the setup is applied for plasma deposition a substrate can be mounted on the sample support. The cascaded arc (on the left side of the graph) is depicted in Fig. 3.

to the substrate by a plasma beam. On the substrate an amorphous, hydrogenated carbon film is grown [2]. In this paper we will constrict ourselves to the description of the evolutions of the most important plasma parameters in the flowing plasma. Since it is not our aim to perform an extensive study of the supersonic flow and the shock waves associated with it, an attempt will be made to develop a simple model that can be coupled easily to chemical reaction rate equations. This last aspect will be treated in a separate paper.

The model which has been used is quasi one-dimensional. It is based on the conservation equations for electron density, argon density, plasma momentum, electron energy and thermal gas energy. In the nozzle where the expansion starts, this quasi one-dimensional approach will be valid, and since the streamlines of the plasma are already straight at the end of the nozzle, these calculations can be continued in the vacuum system. The shock wave is treated as a discontinuity with the aid of the Rankine-Hugoniot relations [3]. In the subsonic relaxation zone immediately after the shock wave the same set of equations is used, but there the plasma diameter is assumed to be determined by radial ambipolar diffusion.

The results of the model calculations will be compared with measurements of the electron density, the gas velocity and the gas temperature by means of H_β -Stark broadening and Doppler spectroscopy respectively. The obtained profiles will be compared with those reported in the literature.

II. Model

A. Supersonic expansion

The source terms of the mass and energy balances of electrons and heavy particles are derived elsewhere [4]. If we consider a stationary situation ($\nabla \cdot \mathbf{j} = 0$) then the number balances for electrons and for argon atoms and ions read:

$$\nabla \cdot (n_e \mathbf{u}) = n_e(n_{ar} - n_{ars}) K_{1+} \quad (1)$$

and

$$\nabla \cdot (n_h \mathbf{u}) = 0. \quad (2)$$

Here $\mathbf{u}$ denotes the plasma velocity, n_e , and n_{ar} represent the densities of electrons and argon atoms, n_{ars} is the density of argon atoms according to the Saha equation and n_h is defined as the total heavy particle-density (argon atoms and ions). In the expansion the loss of electrons due to radiative recombination is much smaller than the transport term, and it will be neglected. If the momentum balances of all species present in the plasma are added and the Lorentz force is set to zero, the result is:

$$\rho(\mathbf{u} \cdot \nabla) \mathbf{u} + \nabla p = 0. \quad (3)$$

Assuming that the heavy particles all have the same temperature T_h , the total energy balance for all heavy particles reads

$$\nabla \cdot \left(\frac{3}{2} n_h k T_h \mathbf{u} \right) + n_h k T_h (\nabla \cdot \mathbf{u}) = -\nabla \cdot \mathbf{q}_h - \mathbf{\Pi}_h : \nabla \mathbf{u} + Q_{eh}, \quad (4)$$

where $\mathbf{q}_h$ denotes the heavy particle heat flux ($\mathbf{q} = -\kappa \nabla T$), $\mathbf{\Pi}_h$ the viscosity tensor and Q_{eh} the elastic energy transfer between electrons and heavy particles:

$$Q_{eh} = 3 \frac{m_e}{m_{ar}} n_e \left(\frac{1}{\tau_{ei}} + \frac{1}{\tau_{e0}} \right) k(T_e - T_h). \quad (5)$$

Here τ_{ei} and τ_{e0} denote the collision times of an electron with argon ions or atoms respectively.

Finally, the electron energy balance is needed:

$$\begin{aligned} \nabla \cdot \left(\frac{3}{2} n_e k T_e \mathbf{u} \right) + n_e k T_e (\nabla \cdot \mathbf{u}) = -\nabla \cdot \mathbf{q}_e - \mathbf{\Pi}_e : \nabla \mathbf{u} - n_e (n_{ar} - n_{ars}) K_{1+} E_{ion} \\ \pm 4\pi C_1 n_e^2 \langle Z \rangle T_e^{1/2} \frac{k}{hc} \xi_{ff} - Q_{eh} - Q_{line} \pm Q_{rec}. \end{aligned} \quad (6)$$

Here Q_{line} and Q_{rec} denote the energy losses of the electrons caused by escape of line radiation and by escape of radiation resulting from recombination to excited levels. The term containing ξ_{ff} represents the inelastic free-free transitions (Bremsstrahlung). Whereas radiative recombination can be neglected in the continuum equation (1), it has to be taken into account here because the energy associated with this process is much larger than the thermal energy of the electrons. Implicitly we have excluded Ohmic dissipation by convective currents which might be present in the expansion.

If the ionization degree α is defined as $\alpha = n_e/n_h$ the mass density ϱ can be expressed by

$$\varrho = m_{ar}(n_{ar} + n_e). \quad (7)$$

Furthermore

$$p = k(n_e T_e + n_h T_h) = \frac{k\varrho}{m_{ar}} (\alpha T_e + T_h). \quad (8)$$

Now the eqs. (1) to (6) can be rewritten as:

$$\nabla \cdot (\alpha \varrho \mathbf{u}) = \alpha \varrho (n_{ar} - n_{ars}) K_1, \quad (9)$$

$$\nabla \cdot (\varrho \mathbf{u}) = 0, \quad (10)$$

$$m_{ar} \varrho (\mathbf{u} \cdot \nabla) \mathbf{u} + \nabla (k \varrho (\alpha T_e + T_h)) = 0, \quad (11)$$

$$\nabla \cdot \left(\frac{3}{2} \varrho T_h \mathbf{u} \right) + \varrho T_h (\nabla \cdot \mathbf{u}) = Q_h \frac{m_{ar}}{k} \quad (12)$$

and

$$\nabla \cdot \left(\frac{3}{2} \alpha \varrho T_e \mathbf{u} \right) + \alpha \varrho T_e (\nabla \cdot \mathbf{u}) = Q_e \frac{m_{ar}}{k} \quad (13)$$

respectively, where Q_h and Q_e represent the rhs of (4) and of (5).

To study the supersonic expansion of the plasma in the nozzle the equations (9) to (13) will be applied to a volume segment which is depicted schematically in Fig. 2. A quasi one-dimensional approach is followed, as is usual in treating nozzle expansions [3]. This implies that the energy and momentum transfer of the plasma to the nozzle wall are neglected and the radial diffusion and heat conduction terms are neglected because they are much smaller than the axial expansion terms. Also the parameters in the equations are regarded as constant over the nozzle cross section. If furthermore viscosity effects are assumed to be negligible, then the set of equations transforms to:

$$A \frac{\partial}{\partial x} (\varrho u) + \varrho u \frac{\partial A}{\partial z} = 0, \quad (14)$$

$$A \frac{\partial}{\partial x} (\alpha \varrho u) + \alpha \varrho u \frac{\partial A}{\partial z} = \alpha \varrho (n_{ar} - n_{ars}) K_1 A, \quad (15)$$

$$\rho m_{ar} u \frac{\partial u}{\partial x} + \frac{\partial}{\partial x} (k \rho (T_h + \alpha T_e)) = 0, \quad (16)$$

$$A \frac{\partial}{\partial x} \left(\frac{3}{2} \alpha \rho T_e u \right) + \frac{5}{2} \alpha \rho u T_e \frac{\partial A}{\partial x} + \alpha \rho T_e A \frac{\partial u}{\partial x} = Q_e \frac{m_{ar}}{k} A, \quad (17)$$

$$A \frac{\partial}{\partial x} \left(\frac{3}{2} \rho T_h u \right) + \frac{5}{2} \rho u T_h \frac{\partial A}{\partial x} + \rho T_h A \frac{\partial u}{\partial x} = Q_h \frac{m_{ar}}{k} A \quad (18)$$

Fig. 2. The reference volume over which the transport equations are integrated to make them quasi one-dimensional. At axial position x all relevant quantities (e. g. the electron temperature) have the value $A(T_e)$, at position $x + dx$ they have the value $A + dA(T_e + dT_e)$.

where A is the nozzle cross section area. This set of equations can be rewritten as

$$\frac{\partial \alpha}{\partial x} = \frac{\alpha}{u} \frac{S_e}{n_e}, \quad (19)$$

$$\frac{\partial u}{\partial x} = 5u \frac{kT}{3u^2 m_{ar} - 5kT_a} \frac{1}{A} \frac{\partial A}{\partial x} - \frac{Q_e + Q_h}{3u^2 m_{ar} - 5kT_a} \frac{2m_{ar}}{\rho}, \quad (20)$$

$$\frac{\partial \rho}{\partial x} = -3\rho \frac{kT}{3u^2 m_{ar} - 5kT_a} \frac{1}{A} \frac{\partial A}{\partial x} - \frac{Q_e + Q_h}{3u^2 m_{ar} - 5kT_a} \frac{2m_{ar}}{u}, \quad (21)$$

$$\begin{aligned} \frac{\partial T_e}{\partial x} = & -2T_e \frac{u^2 m_{ar}}{3u^2 m_{ar} - 5kT_a} \frac{1}{A} \frac{\partial A}{\partial x} + \frac{Q_e + Q_h}{3u^2 m_{ar} - 5kT_a} \frac{4m_{ar} T_e}{3\rho u} \\ & + \frac{2}{3} Q_e \frac{m_{ar}}{\rho \alpha u k} - \frac{\alpha T_e}{u} (n_{ar} - n_{ars}) K_1, \end{aligned} \quad (22)$$

$$\begin{aligned} \frac{\partial T_h}{\partial x} = & -2T_h \frac{u^2 m_{ar}}{3u^2 m_{ar} - 5kT_a} \frac{1}{A} \frac{\partial A}{\partial x} + \frac{Q_e + Q_h}{3u^2 m_{ar} - 5kT_a} \frac{4m_{ar} T_h}{3\rho u} + \frac{2}{3} Q_h \frac{m_{ar}}{\rho u k} \end{aligned} \quad (23)$$

where the effective temperature T_a is defined by

$$T_a = T_h + \alpha T_e, \quad (24)$$

and S_e represents the source term for electron production. The speed of sound c_{pi} in the two-temperature plasma can be deduced from (19) to (23):

$$c_{pi} = \left[\frac{5kT_a}{3m_{ar}} \right]^{1/2}. \quad (25)$$

The eq. (19) to (23) describe the evolution of the ionization degree, the plasma velocity, the densities of the various species and the electron and heavy particle temperatures in plasma that flows supersonically through a nozzle with no wall friction or wall heat interaction. They will be solved numerically.

B. Shock wave

For a normal shock wave the conservation equations for mass, momentum and energy when crossing the shock (Rankine-Hugoniot relations [3]) read

$$\rho u = \hat{\rho} \hat{u}, \quad (26)$$

$$p + \rho u^2 = \hat{p} + \hat{\rho} \hat{u}^2, \quad (27)$$

and

$$h + \frac{1}{2} u^2 = \hat{h} + \frac{1}{2} \hat{u}^2. \quad (28)$$

Here $\hat{p}$, $\hat{\rho}$, $\hat{u}$ and $\hat{h}$ denote the pressure, mass density, velocity and enthalpy after the shock. HSU has derived an expression for the enthalpy h [5]. If the temperature dependence of the partition functions is neglected the result reads:

$$h = \frac{1}{\rho} \left(\frac{5}{2} k n_h T_h + \frac{5}{2} k n_e T_e + n_e E^{1+} \right) \quad (29)$$

This can be transformed to

$$h = \frac{5}{2} \frac{p}{\rho} + \frac{\alpha E^{1+}}{m_{ar}}. \quad (30)$$

As the residence time of the plasma in the shock is small compared to the time needed for substantial recombination the ionization degree α can be kept constant over the shock.

Using the equation of state ($p = \rho R T_a$) and introducing the Mach number $M = u/c_{p1}$ the equations (26) to (30) can be transformed into

$$\hat{p} = p \left(1 + \frac{5}{4} (M^2 - 1) \right), \quad (31)$$

$$\hat{u} = u \left(1 + \frac{3}{4} \left(\frac{1 - M^2}{M^2} \right) \right), \quad (32)$$

$$\hat{\rho} = \rho \frac{1}{1 + \frac{3}{4} \left(\frac{1 - M^2}{M^2} \right)}, \quad (33)$$

$$\hat{T} = T \left(1 + \frac{3}{16} \left(\frac{(5M^2/3 + 1)(M^2 - 1)}{M^2} \right) \right). \quad (34)$$

Each of the relevant densities will be transformed like the mass density ρ . The plasma equilibrium is assumed frozen in the shock wave. At all considered axial positions in the nozzle these Rankine-Hugoniot relations are used to calculate the pressure if a shock was to occur. If this pressure equals the experimentally determined background pressure the location of the shock is found.

C. Subsonic relaxation

In this section the eqs. (19) to (23) will be used again. Since the axial gradients are now no longer larger than the radial gradients, the divergence of the plasma beam will be determined by diffusion. As an approximation the angle φ which the plasma boundary makes with the plasma axis at each axial position can be expressed by

$$\tan \varphi = F \frac{D}{r_p u}, \quad (35)$$

where D represents the diffusion coefficient and r_p the plasma radius. As the plasma mainly consists of argon atoms and ions, D can be approximated by the ambipolar diffusion coefficient for an argon plasma [4]

$$D = \frac{5.52 \cdot 10^{18}}{n_{ar} + n_e} T_e T_i^{-0.36}. \quad (36)$$

The coefficient F depends on the actual radial profile of the particle densities in the plasma beam. However, a number of common profiles (Gauss, Lorentz, Bessel, parabola) yield approximately the same result ($F = 2 \pm 20\%$). The radial heat conduction is estimated in a similar way. The approach followed in this section is similar to the one-dimensional treatment of jet and duct flows by THOMPSON [6].

III. Experiment

The basic concept of the reactor is described above. Fig. 1 outlines the system. In Fig. 3 the cascaded arc used in these experiments is depicted. The precise construction of it is described elsewhere [7]. Basically the arc consists of an anode, a stack of water cooled, electrically isolated copper plates and three cathodes. The complete arc assembly is fixed to the (grounded) anode supporting flange, which is mounted onto a cylinder that can slide in another cylinder over a length of 0.2 m while the vacuum is maintained. This allows axial motion of the arc in the vacuum system without the

Fig. 3. Schematic view of the cascaded arc. The exact construction is discussed extensively elsewhere [7].

need of bringing the apparatus up to atmospheric pressure and axial scans of the optical measurements to be performed without moving the optical system.

The pumping system incorporates a mechanical rotary piston pump and two roots blowers; it is connected to the vacuum vessel by a stepper motor controlled valve. This system makes it possible to set the pumping speed from 10 l/sec up to 700 l/sec. In this way the working pressure in the system and the gas flow can be controlled independently. The gases are fed to the arc entrance through mass flow controllers.

a. The optical System

The optical system which has been implemented to allow for the execution of various spectroscopic techniques is depicted in Fig. 4. Lens L_4 and mirror M_4 are mounted together on a support, which is vertically translatable and can be rotated around a

Fig. 4. The optical system used for the spectroscopic measurements. The various components are introduced in the text. All indicated distances are in mm.

vertical axis. Since the part of the lightbeam between the lenses L_4 and L_2 is parallel, lateral scanning of the plasma without significantly changing the image quality is possible. It is also very simple to change the angle which the observation direction makes with the plasma axis. When mirror M_5 is inserted, a small hollow cathode arc will be imaged on the monochromator instead of the expanding plasma. This hollow cathode arc has a spectrum which almost exclusively consists of Ar II and Ar I lines. The continuum radiation is of negligible strength. Furthermore, the lines do not show any Stark- or Doppler-shift, so they can be used as wavelength reference. A Fabry-Pérot interferometer can be inserted in the parallel part of the path of light. Lens L_2 images the detection volume in the plasma on the diafragma D_1 . Finally, lens L_1 images the diafragma D_1 on the entrance slit of the monochromator. On the exit slit an RCA 5624 photomultiplier is installed, which is connected with the data acquisition system. To

eliminate most of the signal noise, the lightbeam is chopped by the chopper CH, and the signal is fed to a lock-in amplifier. A Jarrel-Ash 0.5 m monochromator selects the wavelength, whereas the Fabry-Pérot interferometer (Halle, pressure scanned) is added in the case of high resolution spectroscopy. For the monochromator usually slit widths of 100 μm were used. With a dispersion of 1.6 nm per mm this results in a spectral width of 0.16 nm.

To get an impression of the size of the detection volume the relative transmitted intensity was calculated for a number of points in the plasma with a matrix method. Fig. 5 illustrates the shape of this detection volume. One can clearly see that the diameter of the imaged part of the plasma is small, but the length is considerable. With an angle of 67 degrees between the observation direction and the plasma axis (as is used mostly for the Doppler-shift measurements) the axial integration distance values around 1 cm.

Fig. 5. Impression of the shape of the detection volume. On the vertical axis the normalized transmitted intensity is plotted against the position in the detection volume. In a cylinder with a length of 30 mm and a diameter of 4 mm around the focal point the relative transmittance is 100%.

The apparatus profiles of the monochromator and of the Fabry-Pérot interferometer were measured with a low pressure argon lamp. The profiles appear to be Voigt profiles. The transmission of a Fabry-Pérot interferometer is given by the Airy-function, which is periodic with a period called the free spectral range (*FSR*). The *FSR* depends on the wavelength λ , the distance d of the two parallel interferometer plates and the refractive index n of the medium between the two plates according to

$$FSR = \lambda^2 / 2nd. \quad (31)$$

When the reflection coefficient of the mirrors is high enough the maxima of the Airy profile can be approximated by Voigt profiles. In Tab. 1 a short list is given of the physical phenomena that cause broadening of the apparatus profile and their associated finesses (ratio of spectral width and *FSR*). The overall finesse of the interferometer is about 40.

The signal from the lock-in amplifier is processed by a M68 000 microcomputer, which is connected to a LSI 11/23 hostcomputer with a background memory of 0.5 Gbyte (see Fig. 6). The operating language is PEP, a Pascal-like, interpreting language [8], which controls the interface system.

Table 1. Finesses of the Fabry-Pérot interferometer

name	formular	value
reflection	$\pi \sqrt{R/(1 - R)}$	62
surface roughness	$S/2$	100
pinhole (diam. d , focal length f)	$4\lambda f^2/(r^2 d)$	215
non-parallelism (mirror dist. d , angle β)	$\lambda/(2d\beta)$	350

The signal of the pressure in the Fabry-Pérot interferometer and the intensity signal of the lock-in amplifier are sampled by a multiplexed 12-bit ADC. The M68000 stores the data on the background memory, and also performs the Voigt-analysis. The line profiles can be plotted on a pen plotter.

Fig. 6. Data acquisition system used for Doppler spectroscopy. Components: monochromator (mon), photo multiplier (PM), chopper (CH), Baratron pressure transducer connected to the Fabry-Pérot interferometer (FP), high-Z amplifier (HZA), look-in amplifier (LIA), amplifier (ampl), analog to digital converter (ADC), recorder (rec), M68000 laboratory micro computer (comp), text terminal (term 1), graphics terminal (term 2), plotter (plot) and printer (print).

The spectral lines emitted by the plasma show Doppler broadening. The thermal movements of the radiating atoms cause a Doppler shift in all observed frequencies. If the velocity distribution is taken Maxwellian with a temperature T , then the line profile has a Gaussian shape with a full width at half maximum (FWHM) γ (in m/s) according to

$$\gamma = \lambda_0 \sqrt{8 \ln^2 kT/mc^2} \tag{37}$$

where m denotes the atomic mass. If all the atoms have an average velocity component v in the direction of the observer the wavelength of the spectral line appears to be

shifted over a range $\Delta\lambda$, with

$$\Delta\lambda = -\lambda_0 v/c \quad (38)$$

where c denotes the velocity of light.

The Stark effect is caused by interaction of charged particles and the radiating atom, and usually results in a Lorentzian line broadening profile (with the prominent exception of hydrogen atoms). The width is proportional to the electron density. GRIEM has accumulated the Stark parameters for a large number of atomic transitions [9]. For the 420.0 nm argon line that is used in the experiments the Stark width is about 60 fm.

In the case of hydrogen also ion broadening has to be taken into account. GRIEM calculated the total width of the hydrogen lines as a function of electron density [10]. The result is given by:

$$n_e = C\Delta\lambda^{3/2} \quad (39)$$

where t is a constant that only slightly depends on the electron temperature, and $\Delta\lambda$ is the full width at half maximum of the line. For the H_β line ($\lambda = 486.1$ nm), for $T_e = 1$ eV, C values $1.2 \cdot 10^{22} \text{ m}^{-3} \text{ nm}^3$ [2]. The other broadening mechanisms, like resonance broadening, van der Waals broadening, etc, all yield negligible contributions to the line broadening.

IV. Results

In the expansion of the plasma into the vacuum system the temperature and velocity of the argon atoms have been determined by Doppler spectroscopy. Fig. 7 presents typical examples of normalized line profiles measured by the Fabry-Pérot interferometer. These profiles are subjected to Voigt-analysis on the laboratory computer to extract the Doppler shift and the Doppler width.

Fig. 7. Examples of normalized line profiles as measured with the Fabry-Pérot interferometer at several positions in the expansion. Also shown is the reference profile (ref. prof) as obtained from the low pressure argon lamp.

Fig. 8. Three-dimensional impression of the spatial profile of the plasma velocity in the expansion.

Fig. 8 presents a three-dimensional impression of the velocity profile. The beam-behavior can be seen clearly in this graph: the width of the lateral velocity profile increases and the centerline velocity decreases while the plasma moves downstream. The same trend is observed with respect to the gas temperature (see Fig. 9). The electron density seems to behave somewhat differently (Fig. 10). The centerline value decreases with increasing distance to the nozzle, but the width of the profile does not increase. The electron densities at the plasma boundary have values significantly smaller than $1 \cdot 10^{20} \text{ m}^{-3}$. As GRIEM's data [9] for the Stark parameters are only listed for larger values of the electron density an extrapolation has been used. It is not expected however that the error made in this extrapolation causes dramatic errors in the determined electron density. Therefore we might conclude that the lateral profile of the electron density is narrower than the temperature profile.

Fig. 9. The measured lateral profiles of the gas temperature at three axial positions.

Fig. 10. The measured lateral profiles of the electron density at three axial positions.

The eqs. (19) to (23) and (31) to (34) are solved by Runge-Kutta integration using Pascal on a PC/AT. Fig. 11 compares the calculated axial velocity profiles with the values measured by Doppler-shift. Measurements at the standard condition (100 Pa) and at a background pressure of 20 Pa are shown. There is a good qualitative agreement between model and measurements, especially in the first (supersonic) part. In the supersonic part of the flow field the calculated profiles agree with the results of calculations by CHOU and coworkers [11]. In the subsonic relaxation zone the velocity profile agrees qualitatively with centerline profiles calculated by WITZE [12]. POISSANT and DUDECK [13] found an identical profile in the subsonic part. The model predicts the position of the shock wave (at standard conditions) at $z = 80$ mm, whereas the measure-

Fig. 11. Comparison of the measured and the calculated axial profiles of the plasma velocity in the expansion. Parameter is the background pressure. The full curves with the points represent the measurements. The dashed curve correspond to the calculated profiles.

ments indicate that the shock starts at about 50 mm. In the relaxation zone just after the shock the measured and calculated velocities differ by a factor 3. The integration over the detection volume implies in this case not only a lateral integration over 40 mm, but also an axial integration over ca. 20 mm. The measured axial profile is therefore integrated over 20 mm. This length is much smaller than the length scales over which the velocity changes (at least in the subsonic relaxation zone), so the errors are small. In the relaxation zone the measured gradients are much smaller than the calculated ones.

All the mentioned discrepancies between model calculations and measurements most probably arise from the fact that the models are quasi one-dimensional whereas the expansion and especially the shock wave are essentially three-dimensional.

The calculated axial profile of the electron density is shown in Fig. 12 together with the measurements. Again there is a good agreement between model and experiment.

Fig. 12. The measured (points and full curve) and calculated axial profiles of the electron density. Parameter is the background pressure.

After having passed the shock wave the electron density increases in the relaxation zone because of deceleration and cooling of the plasma. After the plasma has travelled some distance, three-particle recombination starts to eliminate electrons and ions. This is illustrated clearly in Fig. 13. Here the axial profile of the total electron flux is plotted. Again the experimental points (in the supersonic part) approximately fit the model curves. POISSANT and DUDECK [13] used crossed electrostatic probes to determine the plasma parameters in a supersonically expanding argon arc. They found similar axial profiles for the electron density, including the increase in the first part of the relaxation zone. GOLDFARB [14] also reported an increase of the electron density in a paper discussed by IGRA [15]. In the supersonic part the profiles are similar to those obtained by JENKINS [16], BOWEN et al. [17] and CHOU and coworkers [11].

Three-particle recombination supplies energy to the electrons. Therefore the electron temperature remains higher than the gas temperature. Fig. 14 compares the measured values of the gas temperature with the calculated axial profiles of electron- and gas temperature.

The temperatures decrease strongly in the supersonic part of the expansion. The shock front converts axial translational energy into temperature. Then the temperature decreases because of heat interaction with the surroundings.

Fig. 13. The calculated (full curves) and experimentally obtained (points) axial profiles of the total electron flux. The electron flux φ_e is defined here as $\varphi_e = \int_r 2\pi r n_e(r, z) u \, dr$.

Fig. 14. Axial centerline profiles of the various temperatures in the plasma. The full curve with the points represents the measurements of the gas temperature T_h using Doppler broadening. The dashed and dotted curves correspond to the calculated profiles of gas- and electron temperature respectively.

The experimental values appear to be higher than the calculated ones. Again this is probably caused by defects in the model. In the supersonic part additional heating by viscosity, convective currents and adsorption of radiation emitted in the arc and in the subsonic relaxation zone might diminish the temperature drop due to expansion.

The calculated profiles for the electron- and gas temperature agree qualitatively with profiles calculated by BOWEN and PARK [17], CHOU and TALBOT [11], GOLDFINCH [18] and ZAKREWSKII and coworkers [19]. FRASER, ROBEN and TALBOT [20] used an electron beam to excite an argon-ion line in a supersonically expanding low-density RF-plasma torch and analyzed the decay of the transition with a Fabry-Pérot interferometer. They

Fig. 15. The calculated shape of the plasma beam (dashed curves) and the experimentally determined profiles (full curves). The experimental profiles are obtained using a cathetometer to determine the boundaries of the radiating part of the beam. Parameter is the background pressure.

Fig. 16. Calculated axial profile of the gas pressure. Parameter is the background pressure.

determined the electron temperature with probe measurements. Their axial profiles for electron temperature, gas temperature and gas velocity are very similar to the ones presented in this work.

Fig. 15 compares the boundary of the plasma as assumed by the models with the dimension of the radiating part of the plasma as measured with a cathetometer. The agreement is fairly good. The experimentally determined beam shapes are similar to the profiles calculated by Love [21] for jets exhausting into still air, although the pressures are not in the same range.

Finally, Fig. 16 illustrates the calculated axial profile of the total gas pressure. In the supersonic expansion the pressure drops drastically. Across the shock wave the pressure increases strongly (cf. Fig. 5.15 in OWCZAREK's book [3]). Then the background pressure is gradually approached.

V. Summary and Conclusions

A quasi one-dimensional model is presented to describe the evolution of the most important plasma parameters in a supersonically expanding arc plasma in argon. The results of the calculations are compared with spatially resolved measurements of the electrons density, the gas velocity and the gas temperature. The agreement of model and experiments is good in the first part of the supersonic flow. The measurements clearly indicate that the shock wave cannot be treated as a discontinuity; the calculated profiles show much larger gradients than those observed experimentally. The obtained axial profiles are similar to those reported in the literature.

Acknowledgement

These investigations in the program of the Foundation for Fundamental Research on Matter (FOM) have been sponsored by the Netherlands Technology Foundation (STW). The authors wish to thank M. J. F. VAN DE SANDE and J. J. BLEIZE for technical assistance and T. BROEKHUIZEN, M. CREUWELS, P. VAN DEN DONK, H. KUPPENS, R. MIESEN, C. SEVERIJNS, G. SIPS, M. TIMMERS and Dr. C. J. TIMMERMAN for taking part in the measurements.

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