Formation of plasma ionization state in the course of evaporation by a laser

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The plasma produced when a solid target is evaporated by a laser is located in a stationary layer (of the Knudsen type) adjacent to the target. The boundary of this layer in the nonstationary flow region (the jet itself) is the Jouguet point. Under the assumptions indicated, a theory of the stationary layer is presented and the boundary condition for the jet, namely the connection between the electron temperature and the degree of ionization at the Jouguet point, is obtained.

The dependence of these parameters on the gasdynamic flow in the jet is obtained. A comparison with experiment is made.

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According to Ref. 4, the rate of the ionization \( I^{n-1} \rightarrow I^{n} \) (\( m \) is the number of equivalent electrons in the outer electron shell of the ion) is

\[
S_{i}(T_{i}) = 10^{-4} \frac{m}{2m+1} \left( \frac{E_{i}}{kT_{i}} \right)^{\frac{3}{2}} e^{-\frac{A_{i} \beta}{kT_{i}}} \left[ \text{cm}^{3} \text{sec}^{-1} \right] \quad (1.1)
\]

Here \( E_{i} \) is the ionization potential, \( \beta = E_{i}/T_{i} \), \( A \) and \( \beta \) are constants, and \( \beta \) is the angular momentum of the outer-shell electron.

The rate of the photorecombination \( I^{n-1} \rightarrow I^{n} \) is expressed as follows:

\[
a_{i}(T_{i}) = 10^{-4} \frac{4\pi m^{2} \beta^{2}}{4\pi^{2}} \left( \frac{E_{i}}{kT_{i}} \right)^{\frac{1}{2}} \left( e^{-\frac{A_{i} \beta}{kT_{i}}} - 1 \right) \left[ \text{cm}^{3} \text{sec}^{-1} \right] \quad (1.2)
\]

where \( A \) and \( \beta \) are constants.

In the analytic calculations we shall replace \( m \) by its mean value \( \bar{m} = 2l + 1 \) and approximate the dependence of the ionization potential on the degree of ionization by a power-law function:

\[
E_{i} = E_{i}(1 + l)^{\gamma} \quad (1.3)
\]

The constants \( E_{i} \) and \( k \), as incidentally also the constants \( A \), \( a \), \( B \), and \( b \), are different for different shells. The quantity of interest to us

\[
\gamma_{i}(T_{i}) = \frac{S_{i}(T_{i}) - a_{i}(T_{i})}{S_{i}(T_{i})} \quad (1.4)
\]

can then be written in the form

\[
\gamma_{i}(T_{i}) = \frac{\gamma_{i}(T_{i})}{\frac{E_{i}}{kT_{i}}} = \left[ \frac{e^{-\frac{A_{i} \beta}{kT_{i}}} - 1}{e^{-\frac{A_{i} \beta}{kT_{i}}} - 1} \right] \quad (1.5)
\]

Here

\[
S_{i} = 10^{-4} \frac{m}{2m+1} A \left( \frac{E_{i}}{kT_{i}} \right)^{\frac{1}{2}} \quad (1.6)
\]

\[
r = 10^{-4} \left( \frac{E_{i}}{kT_{i}} \right)^{\frac{1}{2}} \quad (1.7)
\]

The constant quantities \( E_{i} \) and \( k \), and \( S_{i} \) for ions with different outer electron shells are listed in Table I. The headings \( 2s - 2p \) in this table pertain to ions with ground state configurations \( 1s^{2}2s^{n} \) \( (m = 1 - 2) \) and \( 1s^{2}2s^{2}2p^{n} \) \( (m = 1 - 6) \), and

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1. Ionization, Recombination, and Thermal Conductivity

The physical processes that occur in the plasma-formation region are determined by ionization, recombination, and thermal conductivity.
3s – 3p to the configurations 1s22s22p63s23p\((m = 1 - 2)\) and 1s22s22p63s23p\((m = 1 - 6)\). The values of \(E_0\) and \(k\) were determined by least-squares averaging of the ionization potentials cited in Ref. 5. For Ca and Ti we have also constructed a single continuous function that approximates the dependence of the ionization rate on the ion charge for the 2s–2p and 3s–3p shells jointly. The parameters \(E_0\), \(k\), and \(S_0\) were determined with the aid of this function. Their values for Ti are given in Table I in the column headed \((2s - 2p) + (3s - 3p)\), and for Ca in Table III below. In this case, the dependence of the ionization potential on the ion charge, \(E_i = E_0(z + 1)k\), is approximate. This dependence, however, enters in the theory only via the ionization rate, so that this manner of specifying \(E_0(z)\) allows us to use the theory parameters \(E_0\), \(k\), and \(S_0\) without incurring appreciable errors. The close values of \(S_0\) for both approximations (see, e.g., Table III) mean that consideration of two shells results in an error that at any rate is not larger than the case of one shell. At the same time, the accuracy obtained thereby is higher, since the ionization can be taken into account during its earlier stages.

We turn now to the thermal conductivity. According to Ref. 6, the reciprocal thermal-conductivity coefficient is

\[
\frac{1}{\chi} = \frac{m_i}{T} \frac{1}{T_i} - \frac{m_i}{T} \sum n_i(T_n, T_i)
\]  

(1.8)

Here \(m_i\) is the electron mass, \(n_i\) is the electron density, \(T_n\) is the time of electron scattering by the ions, \(T_i\) is the cross section for this scattering by ions with charge \(z_i\), and \(T_n\) is the density of these ions. Since \(T_n = T_i^{z_i}\), we can express in the form

\[
\frac{1}{\chi} = \frac{m_i}{T} \left( \frac{T_i^{z_i}}{T_i} \right) \frac{1}{T_i^{z_i}}
\]

(1.9)

where the mean values are

\[
\bar{z} = \frac{1}{n_z} \sum n_i (T_n, T_i) \bar{z}_n = \frac{1}{n_z} \sum n_i (T_n, T_i) \bar{z}_n
\]

(1.10)

and \(n_z\) is the total density of the heavy particles. By expression \(n_i\) in (1.8) in terms of \(\bar{z}_n\), we have taken into account the neutrality condition

\[
n_z = \sum n_i = \sum n_i
\]

(1.11)

If we assume approximately \(T_n = T_i\), we can write for \(n_i\)

\[
\bar{x} = n_i \bar{T}_n (T_i / E_i)^{z_i}
\]

(1.12)

where \(\bar{n}_i\) is the thermal conductivity at \(\bar{T} = 1\) and \(T_i = E_i\) (see Table I).

### 2. BASIC EQUATIONS OF STATIONARY LAYER

In the derivation of the equations of plasma motion and ionization in the region of the plasma formation we shall assume that all the plasma components have the same drift velocity \(v\). Together with the neutrality condition (1.11) this ensures the absence of electric currents. We consider first the mass, momentum, and energy transport equations:

The mass conservation law:

\[
\frac{d}{dx} \rho = 0, \quad \rho = M n_i
\]

(2.1)
where \( p \) is the density and \( M \) the mass of the heavy particle; we neglect the electron mass.

The momentum conservation law:

\[
\frac{dp}{dx} + \frac{dp}{dz} = 0, \quad p = n_e T_e + \sum_i n_i T_i.
\]

(2.2)

Here \( p \) is the pressure and \( T_e \) the kinetic temperature of ions with charge \( z \). Since the ions are heated only by the electrons, \( T_e < T_\gamma \). At the same time, at a high degree of plasma ionization we have \( n_e \gg n_i \). We can therefore neglect the ion pressure and put \( p = n_e T_e \).

The energy conservation law:

\[
\frac{d}{dx} \left( \frac{p z^2}{2} + \frac{p z^2}{\gamma - 1} + \sum_i n_i z_i^2 \frac{dT_i}{dx} \right) = 0,
\]

(2.3)

where \( \gamma \) is the effective adiabatic exponent that takes account of the energy loss to ionization.

From (2.1) and (2.2) we have

\[
\frac{du}{dx} = n_e \frac{T_e}{z} \frac{dT_e}{dx} = \frac{1}{\gamma - 1} \frac{p}{n_e T_e} \frac{dT_e}{dx}.
\]

(2.4)

Equations (2.7)-(2.10) constitute the complete system of equations for the variables \( u, n_e, z_e, \) and \( T_e \).

The plan for solving these equations is the following:

We specify the values of \( z_e \) and \( T_e \) at the point \( x = x_e \) and designate them \( z_e \) and \( T_e \). We know the value \( n_e \) or \( n \) at the point \( x \). The velocity \( u_e \) at this point can be expressed in terms of \( z_e \) and \( T_e \), using the definition \( u_e \) is the Jouguet point:

\[
u_e = \left( 1 + z_e \right) T_e/M = z_e T_e/T_\gamma.
\]

(2.11)

It follows now from (2.7) and (2.8) that

\[
\frac{p}{n_e^2} = n_e u_e/z_e,
\]

(2.12)

\[
M \frac{n}{z} = a^2 + n T = M \frac{n}{z} + n T_e.
\]

(2.13)

From this we can express \( u_e \) and \( n_e \) in terms of \( T_e \) and \( z_e \) and the parameters \( T_e \) and \( z_e \) (which are as yet unknown). Substituting these expressions in (2.9) and (2.10) we obtain the system

\[
dz/dx = H(u_e, T_e, z_e, T_e), \quad dT_e/dx = G(u_e, T_e, z_e, T_e).
\]

(2.14)

We are interested in a solution of the system \( z(x) \) and \( T(x) \) for which

\[
u_e(0) = 0, \quad T_e(0) = 0.
\]

(2.15)

The first condition is connected with the fact that near the target the degree of plasma ionization is much less than the mean value. The cause of the second condition is that near the target the kinetic energies of all the plasma particles are of the order of the atomic energy, i.e., much less than the average plasma temperature.

The system (2.14) determines a family of trajectories in the \((z, T)\) plane, along which the \( x \) varies. The equation for the trajectories is

\[
dT_e/du_e = H(u_e, T_e, z_e, T_e), \quad H = G/F.
\]

(2.16)

The condition (2.15) means that the trajectory corresponding to the sought solution must pass through the point \((0,0)\). At the same time, if \( z_e \) and \( T_e \) are chosen, the trajectory passing through the point \( J = (z_e, T_e) \) is uniquely defined, and it need not necessarily pass through the point \((0,0)\). For this to be the case, it is necessary to impose some connection between \( z_e \) and \( T_e \). That this connection is unique follows from the fact that there are three free parameters, \( z_e \), \( T_e \), and \( z_e \), but only two conditions (2.15). The connection between \( z_e \) and \( T_e \) is indeed the boundary condition for the region of the nonstationary plasma flow mentioned in the Introduction.

3. PHASE PORTRAIT OF THE FORMATION REGION

Let us determine the trajectories of the system (2.14) in the \((z, T)\) plane. It can be seen from (2.10) that in the entire \((u, T)\) plane we have \( dT_e/du_e > 0 \). We turn now to (2.9): the function \( y_e \) reverses sign on a certain curve \( C \) in the \((u, T)\) plane. We turn now to (2.9): the function \( y_e \) reverses sign on a certain curve \( C \) in the \((u, T)\) plane.
The characteristics of plasma formation region in the case of high (a) and low (b) thermal conductivity of the plasma.
The values of $M$ and $\lambda$ for different ions are given in Table I. The indicated transformations yield

$$\frac{d\beta}{d\xi}=\psi(\xi)\chi(\xi),$$

(4.10)

where we have introduced the functions

$$\psi(\xi)=\psi_{01}^{(\xi)}(\xi)=\psi_{01}^{(\xi)}(\xi),$$

(4.11)

$$\chi(\xi)=\chi_{1}^{(\xi)}(\xi)=\chi_{1}^{(\xi)}(\xi),$$

(4.12)

which depend on the variables

$$\beta=\beta_{01}^{(\xi)}, \quad \xi=\xi_{01}^{(\xi)}.$$

(4.13)

Since the calculations are semi-quantitative, the exact value of $\psi$ is not very important, and we assume $\psi=5/3$. The parameter $q$ is independent of $\xi$ and $\beta$ and is defined as

$$q=x^{2}g_{01}^{(\xi)}+n_{01}^{(\xi)}.$$  

(4.14)

The right-hand side of (4.10) contains, besides the variables $\xi$ and $\beta$, the two parameters $q$ and $\beta_{01}$. The condition for the existence of a trajectory that connects the points $(0,0)$ and $(1,1)$ in the $(\xi,\beta)$ plane yields the connection between $q$ and $\beta_{01}$, i.e., according to (4.14), the connection between $e$ and $\beta_{01}$. Substituting here (4.7) and (4.8), we obtain the sought connection between $T_{w}$ and $z_{k}$.

It will be made clear below that the experimental conditions correspond to $q<1$. We consider therefore only this case. An analysis of (4.10) at small $q$ shows that the trajectories of interest to us, which arrive in the vicinity of the point $(\xi=1, \beta=0)$ from the region of small $\xi$ and $\beta$, lie close to the curve $\beta=\beta_{01}$, i.e., near the parabola $\beta=\xi^{2}$. These trajectories are shown in Fig. 3. It can be approximately assumed that the trajectory leaving the point $(\xi=0, \beta=0)$ and arriving at the point $(\beta=1, \xi=1)$ coincides with the parabola $\beta=\xi^{2}$. In this approximation, the derivative along the trajectory at the critical point is

$$\frac{d\beta}{d\xi}_{\xi=1}=k.$$  

(4.15)

It follows from (4.15) and (4.10) that

$$q_{c}(\beta)=k.$$  

(4.16)

This indeed is the sought connection between $q$ and $\beta_{01}$ at small $q$. Returning to the variables $e$ and $\beta_{01}$, we can rewrite (4.16) in the form

$$\frac{q}{e}=-\frac{1}{e^{3/5}x^{1/10}}\phi_{01},$$

(4.17)

In the derivation of (4.17) we have neglected the quantity $\alpha \leq 1$ in (4.11). This is correct if $\beta_{01}$ is close to unity.

Equation (4.17) is the sought boundary condition that connects the temperature $T_{w}$ with the degree $z_{k}$ of the plasma ionization on the boundary of the region of formation of the plasma and the plasma jet. This equation contains the parameters of the plasma and of the target (via $\alpha, k_{w}$, and $k$), but not the radiation power. We show one more form, sometimes more convenient, of the boundary condition,

$$z_{k}^{1/3}e^{a_{01}}=x_{k}.$$  

(4.18)

which is obtained from (4.17) by expressing $e$ in terms of $\beta_{01}$ and $z_{k}$ with the aid of (4.7). The proximity of the trajectory to the $\beta=\beta_{01}$ curve means that in the region of plasma formation the ratio $E_{w}/T_{w}$ is a (large) constant.

We now obtain the ionization and temperature profiles. We transform over in (2.9) to the dimensionless variables and integrate along the trajectory from $x=0$ to a certain arbitrary $x$. Noting that this integration $\beta=\beta_{01}$ and $e=\xi^{2}$, we obtain after simple transformations

$$x=\frac{I_{\beta_{01}}}{n_{s},s},$$

(4.19)

where

$$I_{\beta_{01}}(\xi)-\frac{1}{2}b_{01}^{1/10}\frac{1-(1-x^{2})}{1+(1-x^{2})}.$$  

(4.20)

$$M_{01}^{1/3}=e_{s}.$$  

(4.21)

Equation (4.19) determines the function $\xi(x)$, i.e., the ionization profile. The temperature profile is $\theta(x)=\xi^{2}(x)$. We put in (4.19) $x=x_{k}$ and $\xi=1$, and use the boundary condition (4.18), in order to express $z_{k}$ in terms of $\beta_{01}$. As a result we obtain the thickness of the formation region

$$z_{k}=e_{s}^{1/10}I_{\beta_{01}}(1)\frac{I_{\beta_{01}}}{n_{s},s},$$

(4.22)

The integral $I_{\beta_{01}}(\xi)$ was obtained by numerical integration for the value $k=1.4$, which is typical of the 2s–2p shells of the ions from Ca to Cu (see Table I). Figure 4 shows the profiles of the ionization, temperature, electron density, and gasdynamic-velocity of the plasma. It must be noted that the derivative of the electron density with respect to the coordinate, $\partial n/\partial x$, and the acceleration, $du/dt=n_{s}^{1/3}dx/dt$, are infinite at the critical point. As indicated in Ref. 2, this is due to the assumption of a $\delta$-like absorption of the laser radiation on the critical surface.

We note that a degree of polarization amounting to 0.9 of the maximum is reached at a distance 0.35$x_{k}$.
5. DEPENDENCE OF THE TEMPERATURE AND DEGREE OF IONIZATION OF THE PLASMA ON THE GASDYNAMIC POWER OF THE PLASMA JET

To find the distribution of the temperature and of the degree of ionization of the plasma over the jet, we must solve the equation that describe the plasma expansion, with the boundary condition (4.18). However, the temperature and degree of ionization in the base of the jet (i.e., at the location of its junction with the formation region) can be obtained only by knowing the total gasdynamic power of the jet. The gasdynamic flux \( W \) can be obtained by calculating it at the section \( x = 0 \); using (2.1), we obtain

\[
W = \frac{M}{L} \left[ \frac{1}{2} R a + \frac{5}{2} T \right] M^{-1}. \tag{5.1}
\]

We express \( T \) in terms of \( \beta \); we introduce the characteristic flux \( W_0 = 3n_E u_0 \) (5.2) and the dimensionless parameter \( \omega = W/W_0 \). We have then from (5.1)

\[
\omega = \frac{L \rho}{2M} \beta^{-\alpha}. \tag{5.3}
\]

This equation together with (4.18) comprises a system from which we can, given \( W \), obtain \( x \) and \( \rho \); and then find \( \beta \), i.e., \( T \), with the aid of (4.7). Eliminating \( x \) from this system, we obtain an equation for determining \( \beta \) from \( \omega \):

\[
\beta = e^{-\frac{L \rho}{2M} \omega^{\frac{1}{\alpha}}}. \tag{5.4}
\]

Having obtained from this \( \beta \), we calculate \( x \) and \( T \) from the formulas

\[
x = \frac{L \rho}{2M} \omega^{\frac{1}{\alpha}} \beta^{-\alpha}, \tag{5.5}
\]

\[
T = e^{-\frac{L \rho}{2M} \omega^{\frac{1}{\alpha}}} \beta^{-\alpha}. \tag{5.6}
\]

By reducing the calculation results we can obtain simple power-law relations that connect \( T_\alpha, T_\beta, \rho_\beta \), and \( u_\beta \) with the gasdynamic flux \( W_\beta \), namely

\[
\Delta = \Delta' (W/10^{12} \text{ W/cm}^2)^{\alpha}, \tag{5.7}
\]

where the values \( \alpha = 1, 2, 3, 4 \) correspond to \( T_\alpha, T_\beta, \rho_\beta \), and \( u_\beta \), with the gasdynamic flux \( W_\beta \), namely

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\[
\Delta = \Delta' (W/10^{12} \text{ W/cm}^2)^{\alpha}, \tag{5.7}
\]
when only the 2s-2p shell is taken into account and in the case of averaging over the 2s-2p and 3s-3p shells. It can be seen from Tables I and II that in both approximations the values of $T$, and $z$, are the same at $W = 10^{12}$ W/cm$^2$. When two shells are taken into account, however, the $z$, ($W$) is less steep, and the thickness of the formation is approximately doubled.

6. EXPERIMENT

We have irradiated bulk flat targets of different materials, in vacuum, by pulses of energy from 1 to 15 J and duration 2.5 nsec from a 1.06-μm neodymium-glass laser. The focused-spot diameters ranged in the different cases from 100 to 1500 μm. The density of the multiply charged ions in the produced plasma was determined from the emission spectra of those ions which were registered in a direction parallel to the target plane. The spectra were obtained with a spatial resolution along the direction normal to the target, and without a temporal resolution. The Ti and Cu plasmas were investigated with the aid of the spectra in the x-ray region, while the Fe and Ni plasmas were investigated in the vacuum ultraviolet region.

A Johann-type spectrograph with a KHP crystal, circularly bent at a radius 150 mm, was placed in a vacuum chamber. The target was inside the spectrograph. The wavelength region in which the spectra of Li-like and Ne-like ions were observed was 15-25 Å. UF-VR x-ray film was used for the photography. The arrangement of the spectrograph and the experimental setup are similar to those described in Refs. 8 and 9.

We determined from the spectrograms the relative intensities of the spectral lines corresponding to transitions of like type in the shells of the different ions, which yielded in turn the relative densities of the ions with different charges. Account was taken of the fact that in a number of cases these lines are optically thick. The necessary probabilities of the radiative transitions and rates of excitation by electron impact were calculated in accord with Ref. 4. From the relative ion densities we obtained the average charge of the plasma ions. The spectra of the multiply charged Ca, Fe, and Ni ions in the 80-200 Å range were registered with a glancing-incidence vacuum spectrograph ($R = 1$ m, $p = 600$ lines/mm, $\alpha = 0^\circ$), equipped with a gold-coated grating. The average charge was determined from the spectral lines corresponding to the transitions between the configurations 2$\ell$2$\ell^\prime$ and 2$\ell^\prime$2$\ell$ ($m = 5 - 0$), which are characterized by small optical thicknesses ($2 < r < 11$ for the most intense lines).

The registration of the vacuum-ultraviolet spectra with spatial resolution is described in Ref. 10.

In some cases (see e.g., Ref. 11), we tracked the behavior of the ion composition of the plasma with increasing distance from the target. It was found that, at least startig with distances 50-100 μm, the ion composition remains practically unchanged with further expansion of the plasma.

The spatial resolution of our spectrograms was mainly approximately 50 μm. We were therefore unable to measure directly the ion composition at the critical point. The experimental values $z$, of the average ion charge, given in Table IV and in Fig. 5, correspond to a distance 50-100 μm from the target surface. The described “freezing” of the ion composition in the expanding plasma, due to the rapid decrease of the plasma electron density and plasma temperature, should make these values of $z$, close to those at the critical point.

The experimental values of the average plasma ion charge are compared in Table IV with the theoretical ones calculated from Eqs. (5.7) with the parameters from Tables II and III. In individual cases, to reconcile the theoretical results with the experimental ones a coefficient $\eta = W/P$ was introduced for the conversion of the laser radiation flux $P$ into the gasdynamic flux $W$. Figure 5 shows the experimental dependence of the average charge $z$ on the laser radiation flux incident on the target for a titanium plasma and the theoretical dependence calculated under the assumption that the total laser-radiation energy was converted into the gasdynamic flux ($\eta = 1$).

It can be seen from the figure and from Table IV that the theoretical values agree well with experiment at laser-radiation fluxes $P \approx 10^{12}$ W/cm$^2$. At larger fluxes, the theoretical

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TABLE IV.

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545 Sov. Phys. JETP 56 (3), September 1982 Levinson et al.
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10 V. I. Bayanov, S. S. Guladov, A. A. Mak, G. V. Perel'skov, I. I. Sobel'man, A. D. Starikov, and V. A. Chirkov, ibid. 3, 2253 (1976) [7, 1226 (1976)].

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