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Evaporation of Metals by High-Density (10^7 A · cm^{-2}) Electrical Currents[1]

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In the present work, the problem of time evolution of pressure and temperature profiles across a wire through which an electrical current with a density of the order of 10^7 A · cm^{-2} flows is solved. The correct boundary conditions for a metal surface are obtained for the case when this metal is rapidly evaporated as a result of high-power Joule heating. The pressure profile appears under these conditions due to pinch-effect and inertia of thermal expansion of the metal; the temperature profile arises because of intensive evaporation from the surface of the wire. The conditions under which a liquid metal is superheated are formulated. On the basis of the analysis of the experimental results on exploding wires, the conclusion is drawn that decay of the metastable state takes place near the binodal. It is shown that the distribution of fine dispersed vapor bubbles is strongly nonuniform across the wire and the process of expansion of the two-phase mixture is very similar to the motion of a wave.

KEY WORDS: evaporation; exploding wires; high temperatures; liquid–vapor mixture; pulse-heating; superheated metal.

1. INTRODUCTION

The study of evaporation resulting from the passage of a high-density electrical current through a metal has a long history and is related to such phenomena as exploding wires [1]. The development of an adequate theory for this process is of great importance for the solution of many problems [2, 3]. In particular, it is of interest for research on metals in the

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area of phase diagram, including the critical point of liquid–vapor phase transition, and in the field of metal–insulator transitions [4, 5].

Despite the long history of research on exploding wires, a generally accepted theory does not exist. The problem is that the nature of evaporation of metals under these conditions is not completely understood. In one of the earliest works devoted to the systematic research of exploding wires [6], experimental dependences (resistance of wire versus enthalpy) were interpreted to be the result of the motion of an evaporation wave, separating the liquid metal from the liquid–vapor mixture which was assumed to be electrically not conducting. The wave starts at the surface of the wire at the beginning of explosion, i.e., when the resistance as a function of the enthalpy begins to increase rapidly. In Ref. 6 it was established that the dependence of the speed of the evaporation wave on temperature is qualitatively very similar to that of sound speed in the two-phase mixture on the binodal. Nevertheless, the investigator [7] postulated, based on the idea of shock regime of superheating [8], maximum superheating of a metal and phase decay on the spinodal. In so doing, isobaric expansion of the metal was assumed.

In investigations devoted to the numerical simulation of exploding wires [9–11], a one-dimensional magnetohydrodynamics model was used; evaporation from the surface of the metal was not taken into account. As a result, the model did not describe the formation of a temperature profile near the surface and real superheating of the metal, caused by evaporation [12].

2. ONE-DIMENSIONAL MAGNETOHYDRODYNAMICS MODEL

We consider the passage of an electrical current pulse through a thin rectilinear metal wire, placed in vacuum. Let the length of the wire l be much larger than its radius a and t a characteristic time $t \ll l/c_s$ (c_s is the sound speed). Then the particles of the material move exclusively in the radial direction, and distribution of current and electromagnetic field can be obtained in the approach of thin wire [13]: the intensity of magnetic field H has only an azimuthal component, and the density of current j has only a component along the axis of the wire. It should be noted that applicability of the one-dimensional model for the regimes with the characteristic current density of the order of $10^7\ \mathrm{A \cdot cm^{-2}}$ and higher is confirmed by many experiments [14–18].

In a cylindrical coordinate system, axis z of which is directed along the axis of the wire, laws of conservation of mass, momentum, and energy can be written as

$$\frac{\partial \rho}{\partial t}+\frac{1}{r}\frac{\partial}{\partial r}(r\rho v)=0 \tag{1}$$

$$\frac{\partial v}{\partial t}+v\frac{\partial v}{\partial r}=-\frac{1}{\rho}\frac{\partial p}{\partial r}-\frac{jH}{c_{\mathrm{f}}\rho} \tag{2}$$

$$\rho T\left(\frac{\partial s}{\partial t}+v\frac{\partial s}{\partial r}\right)=jE+\frac{1}{r}\frac{\partial}{\partial r}\left(r\kappa\frac{\partial T}{\partial r}\right) \tag{3}$$

According to Eq. (1), density ρ varies only due to the radial motion of parts of the wire with the speed v; r is the radius-vector. The Euler equation, Eq. (2), contains the Lorentz force, having in this case only a radial component (c_{f} is the speed of light in vacuum). The law of conservation of energy, Eq. (3), is rewritten for the specific entropy s; here E is the electrical field strength, T is the temperature, and κ is the thermal conductivity coefficient.

We assume that the skin effect is of no importance and therefore the electrical field is a function only of time $E=E(t)$ [19]. Then the Maxwell equations are reduced to the single equation:

$$\frac{1}{r}\frac{\partial}{\partial r}(rH)=\frac{4\pi}{c}j \tag{4}$$

The current density obeys Ohm's law: $j=\sigma E$ (σ is the electrical conductivity).

We are interested in the region of the phase diagram of a metal where temperature is at least below the critical temperature of the metal. In this range of temperature, conductivity of the vapor arising due to evaporation from the surface of a wire can be neglected. In this case, the problem of motion of the vapor can be solved separately from that one of condensed matter; the solution of the problem is necessary to obtain the boundary conditions on the surface of the wire.

3. BOUNDARY CONDITIONS

As it is well-known, near the surface of an evaporating body a thin vapor layer is formed, where the local thermodynamic equilibrium is infringed. The size of this layer is of the order of the mean free path of the molecules in the vapor (Knudsen layer).

We designate a metal by index 1 and vapor by index 2 and write the laws of conservation of mass, momentum, and energy on a plane phase boundary as follows:

$$\rho_1(v_1 - v_s) = \rho_2(v_2 - v_s) \tag{5}$$

$$p_1 - p_2 = \rho_2(v_2 - v_s)^2 - \rho_1(v_1 - v_s)^2 \tag{6}$$

$$q = \rho_1(v_1 - v_s)[w_2 - w_1 + (v_2 - v_s)^2/2 - (v_1 - v_s)^2/2] \tag{7}$$

where v_s is the speed of the boundary in the laboratory frame of reference, $q = -\kappa(\partial T/\partial x)$ is the heat flow in the frame of reference which is in rest in relation to the boundary, and w is the specific enthalpy; the x-axis is directed perpendicular to the boundary and the metal occupies the region $x > 0$.

In Ref. 20, the gas-kinetics problem of evaporation of a metal in vacuum was considered and equations relating the thermodynamics quantities for the vapor near the surface (but beyond the limits of the Knudsen layer and a condensation jump) to the temperature of the surface were obtained:

$$T_2 = 0.79T_1, \qquad p_2 = 0.37P_{st}(T_1), \qquad v_2 = 0.55c_{s2}(T_2) \tag{8}$$

where $P_{st}(T)$ is the saturation vapor pressure, and c_{s2} is the sound speed in the vapor.

Using Eqs. (8), we obtain the following equations from Eqs. (5) and (7), which will play the role of the boundary conditions on the surface of the metal:

$$(v_1 - v_s)\,\rho_1 = -0.30P_{st}(T_1)[\mu/(RT_1)]^{1/2} \tag{9}$$

$$p_1 = 0.56P_{st}(T_1) \tag{10}$$

$$\kappa_1 \frac{\partial T_1}{\partial x} = 0.30\lambda_{ef}(T_1)\,P_{st}(T_1)/(\mu RT_1)^{1/2} \tag{11}$$

where R is the molar gas constant, μ is the molar mass, and $\lambda_{ef}(T) = \lambda(T) - 0.325RT$ (λ is the molar latent heat of evaporation). These relations contain only the linear contributions of the ratio of density of the vapor to density of the liquid. The approach suggests a small difference between the inner energy and the enthalpy of the condensed matter. It is also assumed that vapor is described by the equation of state of ideal gas.

For the difference between enthalpies of the liquid and the vapor, one can write

$$w_2(T_2) - w_1(T_1) = \lambda(T_1) - 0.525RT_1/\mu$$

the dependence of λ on T is given by

$$\lambda = \lambda_0 + (C_{p2} - C_{p1})\, T \tag{12}$$

where $C_{p2} = 5R/2$ and λ_0 is the molar heat of evaporation at $T = 0$ K; we suppose that the molar heat capacity of a liquid metal C_{p2} is also a constant [4].

The temperature dependence of the saturation vapor pressure has the form

$$P_{st}(T) = A \exp[-\lambda_0/(RT)] \tag{13}$$

where A is a constant [21, 22].

4. PRESSURE AND TEMPERATURE PROFILES

Let us seek the solution of Eqs. (1)–(3) with a slightly nonuniform density profile:

$$\rho(r, t) = \rho_h(t) + \rho'(r, t)$$

where $|\rho'| \ll \rho$ (to be short, here and below the index 1 is dropped).

Equation (1), after substitution of such a profile, is easily integrated. As a result, for the speed of a particle of material, we obtain

$$v = -\frac{r}{2} \frac{d \ln \rho_h}{dt} \tag{14}$$

Using Eq. (14), from Eq. (2) the following equation for the pressure profile can be derived:

$$p = p_s + \frac{1}{c_f} \int_r^a jH\, dr + \frac{\rho_h}{4} \left[\frac{1}{2} \left(\frac{d \ln \rho_h}{dt} \right)^2 - \frac{d^2 \ln \rho_h}{dt^2} \right] (a^2 - r^2) \tag{15}$$

where p_s is the pressure on the surface of the wire.

As one can see, the contribution to the pressure by the magnetic field $\int_r^a jH\, dr/c_f \sim u_A^2$ ($u_A = H/\sqrt{4\pi\rho}$ is the Alfven speed), and the inertial contribution is of the order of $\rho a^2(\alpha_p\, dT/dt)^2$, where α_p is the coefficient of thermal expansion. Since $\rho' \sim p/c_s^2$, the inequality $|\rho'| \ll \rho$ is satisfied when $u_A^2 \ll c_s^2$, and $(a/t_T)^2 \ll c_s^2$, where t_T is the characteristic rise time of temperature.

Next, we discuss the temperature profile. The estimations show that the contribution of thermal expansion at $T < 0.9T_c$ (T_c is the critical temperature) is relatively small and the size of a layer where the temperature has an appreciable gradient due to evaporation of the metal from the surface of the wire is much smaller than a. Hence, the equation for the temperature profile can be presented for a flat surface.

Based on these remarks, and from Eq. (3), one may write

$$\frac{\partial T}{\partial t} = \chi \frac{\partial^2 T}{\partial x^2} + (v_s - v) \frac{\partial T}{\partial x} + \frac{\mu \sigma E^2}{\rho C_p} \tag{16}$$

where $\chi = \kappa\mu/(\rho C_p)$ is the thermal diffusivity. The frame of reference in rest in relation to the surface is used.

We also represent the temperature profile in the form

$$T(x, t) = T_h(t) - T'(x, t)$$

($|T'| \ll T$); the limits of applicability of such an approach are established below. Equation (16) can be brought to a system, consisting of a uniform equation for T':

$$\frac{\partial T'}{\partial t} = \chi(T_h) \frac{\partial^2 T'}{\partial x^2} + (v_s - v) \frac{\partial T'}{\partial x}$$

with the boundary condition, Eq. (11), and the ordinary differential equation for T_h:

$$\frac{dT_h}{dt} = j^2 \mu/(\sigma \rho C_p)$$

We consider the so-called fast regimes of exploding wires when the total current through the sample I on the stage from the end of melting of the sample up to the beginning of explosion is changed relatively little, i.e., we can assume that $I = \text{constant}$. In this case, the equation for T_h can be easily integrated.

The formal solution of the equation for T' can be obtained by means of the Green functions method [23]. Taking into account the temperature dependence of χ, the method gives, for $T'_s(t) \equiv T'(a, t)$

$$T'_s = \mu C_p^{-1} \int_0^t \exp[-x_s^2(t, t')/4\xi(t, t')] \frac{\lambda_{ef}(t')\, \nu_s(t')}{\sqrt{\pi \xi(t, t')}}\, dt' \tag{17}$$

where $x_s = \int_{t'}^{t} [v_s(t_1) - v]\, dt_1$, $\xi = \int^{t\ t'} \chi(t_1)\, dt_1$, and $\nu_s(t) = v_s(t) - v$ is the speed of evaporation.

Temperature T' in a thin layer near the surface can be represented by the expansion $T'(x, t) \approx T'_s + (\partial T'_s/\partial x)\, x$. As the boundary condition, Eq. (11), relates the gradient of temperature on the surface with the temperature of the surface, for the description of the profile we need only to obtain the dependence $T'_s(t)$, where the index s denotes surface quantities.

The estimations show that $x_s^2 \ll 4\xi$ and it is possible to neglect the displacement of the surface in Eq. (17). It is convenient to pass in this equation from the variable t to T_h. After this transition, Eq. (17) takes the form

$$\vartheta = B \int_0^\tau \frac{\exp\{-1/[\tau' - \vartheta(\tau')]\}}{[(\tau' - \vartheta(\tau'))\,\tau'(\tau - \tau')]^{1/2}}\, d\tau' \tag{18}$$

with

$$B = \frac{\nu_0 R\sqrt{t_T \tau_b/(\pi\chi_m)}}{C_p}, \qquad t_T = \frac{\rho_m C_p T_m \sigma_m}{j_m^2 \mu}, \qquad j_m = \frac{I\rho_m}{\pi a_0^2 \rho_0}$$

index m denotes the quantities in the liquid phase at the melting point. We have introduced the following dimensionless variables: $\tau = T_h R/\lambda_0$, $\vartheta = T'_s R/\lambda_0$ ($\tau_b = T_b R/\lambda_0$, where T_b is the boiling temperature at atmospheric pressure). In Eq. (18) we have neglected the temperature dependence of λ. For the speed of evaporation we make use of the expression

$$\nu_s = \nu_0 (T_b/T)^{1/2} \exp[-\lambda_0/(RT)]$$

which results directly from Eq. (9); $\nu_0 = 0.3\text{A}/(\rho_m \sqrt{RT_b/\mu})$.

The nonlinear integral equation, Eq. (18), completely defines the function $\vartheta(\tau)$. We consider first the case $\vartheta \ll \tau^2$. Expanding the right-hand part of Eq. (18) in a power series on ϑ, in linear approximation on ϑ/τ^2, it follows that

$$\vartheta(\tau) \simeq B \int_0^\tau \exp(-1/\tau')\, \frac{1 - \vartheta/\tau'^2}{\tau'(\tau - \tau')^{1/2}}\, d\tau' \tag{19}$$

It should be noted that τ obeys the inequality $\tau \ll 1$. For this reason, the terms proportional to ϑ/τ were dropped. We solve Eq. (19) by the step-by-step method, and its solution is given by

$$\vartheta/\tau^2 = \eta(1 - \eta/\sqrt{2} + \eta^2/\sqrt{6} - \eta^3/\sqrt{24} + \cdots) \tag{20}$$

with $\eta = \vartheta^{(1)}/\tau^2$, and

$$\vartheta^{(1)} = \text{B}e^{-1/2\tau} K_0[1/(2\tau)]/\sqrt{\tau}$$

where $K_0(z)$ is the McDonald function. The series converges at all η, and its sum $S(\eta) < 1$. The function $K_0(z)$ can be presented for $z \gg 1$ by $K_0(z) \approx [\pi/(2z)]^{1/2} \exp(-z)$. Hence, for $\vartheta^{(1)}$, we obtain

$$\vartheta^{(1)} \approx \sqrt{\pi}\, B \exp(-1/\tau) \tag{21}$$

The error, arising due to breakdown of the expansion of Eq. (18), becomes considerable at $\eta \sim 1$. The estimations show that it is necessary to have a formula for these values also. To obtain such a formula one has first to take into account the dependence of the exponent $\exp[-1/(\tau - \vartheta(\tau))]$ on τ more correctly. Let us pass in Eq. (18) to the variable $\varphi = \tau - \vartheta(\tau)$:

$$\vartheta = B \int_0^{\varphi} \frac{\exp\{-1/\varphi'\}(d\tau/d\varphi')\, d\varphi'}{[(\tau - \tau(\varphi'))\, \tau(\varphi')\, \varphi']^{1/2}} \tag{22}$$

We again make use of the step-by-step method. For the first approximation we take $\tau = \varphi$ and $d\tau/d\varphi = 1$ ($\vartheta \ll \tau$). As a result, we have

$$\vartheta = \mathrm{B}e^{-1/2\varphi} K_0[1/(2\varphi)]/\sqrt{\varphi}$$

When $2\varphi \ll 1$, we can write

$$\vartheta \approx \sqrt{\pi}\, B \exp\{-1/[\tau - \vartheta(\tau)]\} \tag{23}$$

As one can see, the expression turns into Eq. (21) at $\vartheta \ll \tau^2$. The next term given by the step-by-step method can be found if one uses the functions $\tau(\varphi)$ and $d\tau/d\varphi$ in Eq. (22) according to Eq. (23). It may be shown that the term differs from those given by Eq. (23) by a factor of the order of one, weakly depending on T.

Next we determine the thickness of the layer near the surface inside which the temperature has a marked gradient. We define the thickness δ_T of the layer by $\delta_T = T'_s/(\partial T_s/\partial x)$. After substitution in the formula the expression for T'_s according to Eq. (23) and $\partial T_s/\partial x$ according to Eq. (11), we obtain

$$\delta_T \approx \sqrt{t_T \chi \varphi} \tag{24}$$

Therefore the limit of applicability of the approach used for the calculation of the temperature profile is expressed by $\sqrt{t_T \chi \tau_b} \ll a$. Besides that, the requirement $T'_s \ll T_h$ leads to

$$\sqrt{\pi}\, B \ll \varphi \exp(1/\varphi) \tag{25}$$

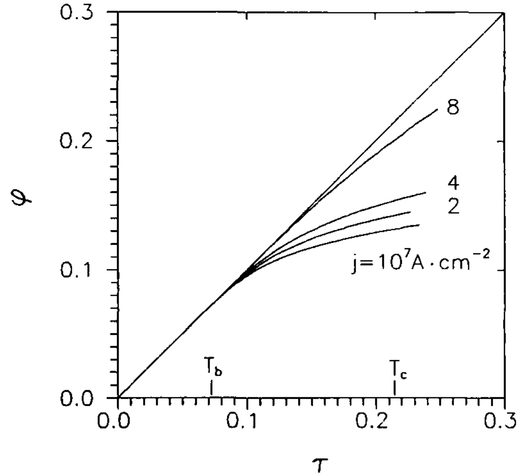

Fig. 1. Reduced temperature φ as a function of τ. The curves present the dependences at various values of the electrical current density j for Al.

In Fig. 1, the temperature on the surface φ as a function of the temperature far from the surface τ for several values of the electric current density j is shown.

5. CONDITION OF SUPERHEATING

As follows from Eq. (10), a metal near its surface is superheated at all temperatures; $P_s < P_{st}(T_s)$. The value of relative superheating $\delta T_s/T_s \equiv [T_s - T_{st}(P)]/T_s$, is given by

$$\delta T_s/T_s = 0.58\varphi/(1 + 0.58\varphi)$$

The function monotonously increases and its value at $T_s = T_c$ for Al is equal to 0.11. The parameters of the critical point we take according to Ref. 24 (for Al $T_c = 8000$ K).

Let us find the conditions under which a metal is superheated far from the surface. When calculating the thickness of the superheated layer δ, we assume first that $\delta \ll \delta_T$. We expand the equality $P(\delta, t) = P_{st}(T(\delta, t))$ in a Taylor series in δ. In the linear approximation, we obtain

$$\delta \approx 0.44 P_{st}(T_s)/\{(\partial P_s/\partial x) - [dP_{st}(T_s)/dT](\partial T_s/\partial x)\} \tag{26}$$

As a result of Eq. (24), δ increases infinitely when

$$\partial P_s/\partial x \to [P_{st}(T_s)/dT]\,\partial T_s/\partial x$$

We introduce the characteristic temperature T^* at which the denominator in Eq. (26) becomes zero. As one can see, till $T < T^*$ the superheating of the metal is maximum near the surface. As soon as $T > T^*$, the maximum is reached in the volume of the metal. Thus, the temperature T^* characterizes that value of temperature at which superheating of metal in volume becomes appreciable (greater than on the surface) and, in addition, the size of the superheated layer becomes greater than δ_T. We find the pressure gradient on the surface by differentiating Eq. (15) and the temperature gradient from Eqs. (23) and (11). As a result, we have the following equation for the temperature $\varphi^* = RT_s^*/\lambda_0$:

$$\frac{\exp(-1/\varphi^*)}{(\varphi^*)^{5/4}}\,[\lambda_{\mathrm{ef}}(\varphi^*)]^{1/2} \approx Dg \tag{27}$$

where $g = j\sqrt{\alpha}$, and D is a constant depending only on the thermophysical parameters of a metal. The temperature far from the surface τ^* corresponding to φ^* can be determined by $\tau^* = \varphi^* + \vartheta(\varphi^*)$. When deriving Eq. (27), we have neglected the temperature dependences of κ and ρ. Besides that, in Eq. (15) we have neglected the inertia contribution, which is small up to $j \sim 10^8\ \mathrm{A \cdot cm^{-2}}$.

After reaching the temperature T^*, the further evolution of the temperature and pressure profiles depends on the kinetics of the decay of the metastable state. In the case of homogeneous nucleation, the metal is strongly superheated and decay takes place near the spinodal. However, at typical pressures for the fast regimes of exploding wires, $P \sim 10^2$–10^3 bar, the temperature on the spinodal $T_{\mathrm{sp}}(P)$ is close to T_c, while the temperature, corresponding to the beginning of explosion, can achieve values $T_e \geqslant 0.5T_c$. So low values of T_e, in our opinion, indicate heterogeneous nucleation under conditions created by exploding wires.

Nevertheless, heterogeneous nucleation cannot happen in an appreciable volume suddenly after reaching temperature T^*. Indeed, the speed of sound c_{mx} in a liquid metal with finely dispersed vapor bubbles distributed in it at temperature $T = T_b$ is of the order of $10\ \mathrm{cm \cdot s^{-1}}$, i.e., it is a very small value. Therefore, the bubbles arise only in a layer of the thickness of the order of $\int_0^t c_{\mathrm{mx}}\,dt$. We introduce the characteristic temperature T_v, at which $\int_0^t c_{\mathrm{mx}}\,dt = a$. Performing the integration, we obtain

$$\frac{1 + \alpha'(\tau_v - \tau_m)}{1 + \beta'(\tau_v - \tau_m)}\sqrt{\tau_v}\exp(-1/\tau_v) \approx Fg^2 \tag{28}$$

where F is a constant depending only on the thermophysical parameters of a metal. The sound speed in the two-phase mixture (at infinitesimal mass fraction of vapor) is equal to [13]

$$c_{\mathrm{mx}}(T) = P_{\mathrm{st}}(T)\,\lambda_0/[\rho RT\sqrt{(C_{\mathrm{p}}T/\mu)}]$$

We have used the following temperature dependences of the conductivity and density for the liquid metal:

$$\sigma = \sigma_{\mathrm{m}}/[1+\beta(T-T_{\mathrm{m}})], \qquad \rho = \rho_{\mathrm{m}}/[1+\alpha(T-T_{\mathrm{m}})] \tag{29}$$

where $\beta' = \beta\lambda_0/R$, $\alpha' = \alpha\lambda_0/R$ (T_{m} is the melting temperature).

Thus, there are two cases: $T^* > T_{\mathrm{v}}$ and $T^* < T_{\mathrm{v}}$. The first is characterized by a relatively uniform expansion of the wire. In the second case, expansion is nonuniform. At first, the vapor bubbles arise in a thin layer near the surfaces, and only after reaching the temperature T_{v} (do) bubbles arise near the axis of the wire. It is of interest to find out what situation takes place for the typical regimes of exploding wires.

In Fig. 2, dependences of characteristic temperatures τ_{v} and τ^* on the parameter $j\sqrt{a}$ for Al, W, and Pb are presented. From Fig. 2, it may be seen that for the regimes presenting a practical interest ($g \sim 0.1$–10), the situation $\tau^* < \tau_{\mathrm{v}}$ is realized. The parameters used in these calculations are listed in Tables I and II.

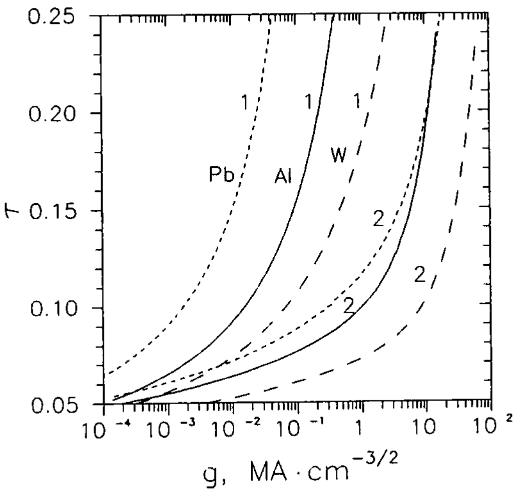

Fig. 2. Dependences of τ_{v} (curves with index 1) and τ^* (curves with index 2) versus the parameter g for Al (——), W (– – –), and Pb (- - - -).

Table I. Parameters for Some Selected Metals[a]

	λ_0 ($kJ \cdot mol^{-1}$)	T_b (10^3 K)	T_m (10^3 K)	C_p ($J \cdot mol^{-1} \cdot K^{-1}$)	σ_m^{-1} ($\mu\Omega \cdot cm$)
Al	113	2.79	0.93	32.4	26.1
Cu	338	2.82	1.36	30.5	21.5
W	838	5.95	3.65	47.8	146
Pb	197	2.00	0.60	33.2	95.0

[a] λ_0 is the molar heat of evaporation at $T = 0$ K; T_b is the boiling temperature at atmospheric pressure; T_m is the melting temperature at atmospheric pressure; C_p is the molar heat capacity of liquid metal; σ_m is the electrical conductivity of liquid metal at the melting temperature.

In Fig. 3, the dependences $\tau^*(g)$ and $\tau_v(g)$ for Cu are presented with marked experimental valued of the temperature $\tau_e = RT_e/\lambda_0$. As one can see, τ_e is close to τ^* and is always lower than the temperature on the spinodal $RT_{sp}/\lambda_0 \approx 0.21$ (at pressure $P \sim 10^2$–10^3 bar), which confirms the conclusion regarding heterogeneous nucleation and small superheating of metals under conditions created by exploding wires. The critical temperature for *Cu* according to [24] is equal to 8390 K.

The temperature τ^* increases with increasing g and can reach even the critical temperature. The noted situation is in good agreement with the well-known fact of a rise in energy imparted to the wire at the beginning of explosion with an increase in current density [2, 3]. At large values of g ($g \sim 10$) superheating and, therefore, the generation of vapor bubbles are impossible.

Table II. Parameters for Some Selected Metals[a]

	χ_m ($cm^2 \cdot s^{-1}$)	A (GPa)	β (10^{-4} K^{-1})	α (10^{-4} K^{-1})
Al	0.32	96.4	4.5	1.4
Cu	0.40	144	3.8	1.0
W	0.14	3520	0.0	0.93
Pb	0.091	12.2	3.2	1.3

[a] χ_m is the thermal diffusivity of liquid metal at the melting temperature; A is the coefficient in Eq. (13) for the temperature dependence of the saturation vapor pressure; β is the coefficient in Eq. (29) for the temperature dependence of electrical conductivity; α is the thermal expansion coefficient of liquid metal at the melting temperature.

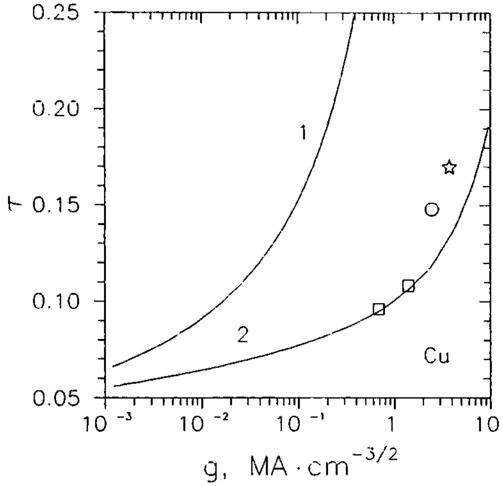

Fig. 3. Temperature τ_v(1) and τ^*(2) as functions of g for Cu. The symbols present the experimental values of τ_e: circle [15]; asterisk [16]; squares [25].

6. CONCLUSIONS

Superheating of a metal for typical regimes of exploding wires ($j \sim 10^7$ A·cm^{-2}) arises due to intensive evaporation from the surface of the wire. The values of superheating are relatively small; the decay of the metastable state takes place near the binodal. The distribution of the vapor bubbles across the wire is strongly nonuniform. The process of expansion of the two-phase mixture is very similar to the motion of a wave. The regimes, distinguished by a large value of the parameter $j\sqrt{a}$, exist, when superheating and generation of vapor bubbles are impossible.

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