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IMPACT OF AN EVAPORATING DROP ON A HEATED WALL

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UDC 536.24:536.423.1

The interaction of a superheated surface with a drop impinging upon it is treated in the quasistationary approximation, neglecting dissipative effects.

The cooling of heat-transfer surfaces by jets of dispersed liquid and by other streams carrying drops is widely employed in contemporary power engineering, metallurgy, cryogenics, and other branches [1-5], and has therefore been studied intensively. Primary attention has been paid to experimental study of the motion and heat transfer of a dispersed flow with the surfaces being cooled under conditions approximating those of industry [6-11]. Many fewer papers have been published on the analysis of elementary interactions of single drops with walls, which determine the main features of real cooling processes. Aside from numerous studies of the evaporation of drops lying on a surface, and their impact with walls without a phase transition, we point out that the dynamic and thermal interaction of drops with a surface complicated by evaporation was studied in [12-18].

As a drop approaches a superheated wall, the part of the surface of the drop facing the wall absorbs heat, which is expended mainly in evaporating liquid. The vapor which is formed is squeezed out of the thin layer between the drop and the wall, which leads to an increase in the pressure in it and to the production of forces which impede the motion of the drop. If other conditions remain the same, and the initial kinetic energy of the drop is high enough or the wall temperature is low enough, the thickness of this layer may become less than the height of the roughness protuberances on the wall or the range of the forces of molecular attraction, so that the drop comes into direct contact with the wall. For a low kinetic energy or a high wall temperature the drop is slowed down before this layer becomes sufficiently thin; as a result the drop (or the smaller drops which result from its fragmentation) rebounds from the surface. The critical values of the energy and temperature correspond to the well-known heat-transfer crisis [1-18]. The analysis of the conditions for the onset of this crisis is of primary interest for applications.

Basic Assumptions. The combined nonlinear and unsteady problems of hydrodynamics and heat-conduction theory with unknown boundaries which arise in the theoretical study of this phenomenon are exceptionally complex, and a constructive analysis is possible only by making very substantial simplifications. We completely neglect the initial underheating of the drop up to the saturation temperature, the radiation heat flux to it, ordinary hydrodynamic resistance to its motion, the change in mass of the drop in the collision process, and the corresponding reaction. We assume that all the mechanical and thermophysical parameters of the liquid and vapor or the vapor–gas mixture are constant and uniform, corresponding, i.e., to the saturation temperature or to a certain mean temperature in the vapor layer. We assume that the wall temperature is independent of the evaporation of the drop, which is admissible if the heat capacity and thermal conductivity of the wall material are high enough [16, 18]. In addition, in analyzing the motion and heat transfer in the vapor layer we neglect the de-

Ural State University, Sverdlovsk. Moscow Institute of Chemical Machinery. Translated from Inzhenerno-Fizicheskii Zhurnal, Vol. 43, No. 6, pp. 945-955, December, 1982. Original article submitted September 11, 1981.

 0022-0841/82/4306-1362$07.50

pendence of the equilibrium temperature at the evaporation surface on the pressure in it, the velocity and temperature jumps at its boundaries, and the forces of molecular interaction between them. It can be shown that under normal conditions this is quite admissible if the surface is rough and $h \gtrsim 10^{-4}$ cm. None of these assumptions are basic, in the sense that none determine the principal features of the impact of evaporating drops on superheated surfaces, and appropriate corrections for them can be introduced into the analysis without particular difficulty.

The following assumptions are more serious: In order to obtain fundamental results in visible form and not to complicate the basic idea by details of the calculations, we first neglect nonstationary effects in calculating both the heat flux to the evaporation surface and motion in the vapor layer, treating the hydrodynamics and heat-conduction equations in the quasistationary approximation. Second, we neglect dissipative effects related to motion within the drop, and the additional pressure force resulting from the slowing down of the drop, since this would necessitate either the escape of part of the vapor (or vapor–gas mixture) from the vapor layer during its thinning down, or the flow of vapor into the layer during its thickening, even when there is no evaporation. The third simplification concerns the shape of the drop, and is introduced later; practically, it reduces to the modeling of the actual deformation of a drop by a liquid disk of varying radius. It is clear then that it is not necessary to consider motion within the drop. In a number of cases these assumptions may appreciably distort the real physical picture of the collision process. We discuss this briefly below.

Basic Equations. Let us consider a drop of volume V impinging normally on a wall with a velocity v_0. Under the assumptions made, the equation of conservation of energy is written in the form

$$\delta E + \delta F = \delta A, \tag{1}$$

where δA is an element of work performed on the drop by pressure forces due to the evaporation of the drop as it approaches the wall; δA is generally not a total variation, just as dA is not a total differential.

The free "potential" energy of the surface tension is equal to σ times the surface area of the drop, and is a functional of the unknown shape of the deformed drop, which is uniquely determined, as is the position of the center of mass of the drop, by a function ψ in the laboratory coordinate system. The variation of the work also depends functionally on this function. The form of ψ depends on the time.

Obviously the approach and subsequent withdrawal of the drop from the wall may, in principle, correspond to each of an infinite set of trajectories in function space $\{\psi\}$. In a certain sense the limiting trajectories correspond to a change in time of only the shape of the drop or only its position (the drop is deformed but does not approach the wall, or it approaches the wall but is not deformed). Actually, a completely definite trajectory must be realized, and the problem arises of choosing it from all which are possible in principle. In significantly dissipative statistical systems such a choice is usually made within the framework of a phenomenological description by applying one of the extremum principles of the thermodynamics of irreversible processes [19]. In the nondissipative system considered here it is natural to use as a variational principle a variant of the general Le Chatelier principle, according to which the real variation $\delta\psi$ must be such that the external action of the wall on the drop is minimum. This leads at once to the problem of finding a function ψ for which $\delta E/\delta\psi = 0$, $\delta^2 E/\delta\psi^2 > 0$, where $\delta/\delta\psi$ symbolizes the variational derivative. Formally, this problem can be written as follows:

$$\frac{\delta}{\delta\psi}(-F + A) = 0, \quad \frac{\delta^2}{\delta\psi^2}(-F + A) > 0, \tag{2}$$

where $\delta A/\delta\psi$ and $\delta^2 A/\delta\psi^2$ are to be understood as limits of the ratios of the elementary work δA and its variation $\delta^2 A$ to $\delta\psi$ and $\delta\psi^2$ as $\delta\psi \to 0$.

In view of the very serious difficulties involved in the analysis of this kind of problem, even in considerably simpler cases, it is natural to convert from the function space $\{\psi\}$, characterizing the actual system with an infinite number of degrees of freedom, to ordinary phase space of finite dimensionality. In principle this can be done by the well-known method of collocations, specifying ψ a priori as a sum of a finite number of specially selected functions of spatial coordinates (e.g., spherical harmonics) with time-dependent co-

efficients, playing the role of phase-space coordinates. The simplest version of such a method corresponds to representing the drop as a disk of radius R and thickness $V/\pi R^2$, separated from the surface by a plane parallel vapor layer of thickness h. In this case, which we consider in this paper, the phase space is the (R, h) plane.

The free energy and the value $R = R_0$ at which it is minimum are expressed in the form

$$F = 2\pi\sigma\left(R^2 + 2\frac{R_0^3}{R}\right), \quad R_0 = \left(\frac{V}{2\pi}\right)^{1/3}. \tag{3}$$

To calculate the force f acting on the drop we consider quasistationary problems of motion and heat transfer in the vapor layer, assuming that it is thin and that the Reynolds and Peclet numbers are small. From the hydrodynamics problem, assuming no slipping at both boundaries of the layer, we calculate an expression for the mass of evaporating liquid $j = -(h^3/6\nu r)(\partial p/\partial r)$, and from the heat-conduction problem, an expression for the heat flux to the drop $q = (\lambda/h)\Delta T$, $\Delta T = T_W - T_S$ (the quantities j and q refer to unit area of the evaporation surface); setting q = jL, we obtain an equation for the pressure in the vapor layer; solving it and integrating the result over the evaporation surface, we obtain

$$f = \frac{3\pi}{2}\frac{\nu\lambda\Delta T}{L}\frac{R^4}{h^4}. \tag{4}$$

The quantity $\delta A = f\delta h$, and problem (2) takes the following form in the variables R and h:

$$\frac{d}{dR}\left(-\frac{dF}{dR} + f\frac{dh}{dR}\right) = 0, \quad \frac{d^2}{dR^2}\left(-\frac{dF}{dR} + f\frac{dh}{dR}\right) > 0. \tag{5}$$

Integrating the first equation and using (3) and (4), we obtain

$$\frac{3H}{h^4}\frac{dh}{dR} = \frac{8}{R^4}\left(-C + R - \frac{R_0^3}{R^2}\right), \quad H = \frac{\nu\lambda\Delta T}{\sigma L}. \tag{6}$$

The initial condition for this equation is $h = \infty$ at $R = R_0$. We determine the constant C by using the fact that in the present case Eq. (1) shows that $dE = -dF + f dh$. Integrating this equation with the initial condition $E = E_0$ at $R = R_0$, and requiring that $E = E_* = 0$ for a certain $R = R_* > R_0$, we obtain

$$C = \mathrm{We}\frac{R_0^2}{R_* - R_0}, \quad \mathrm{We} = \frac{\rho v_0^2 R_0}{4\sigma}. \tag{7}$$

Introducing the dimensionless variables

$$\eta = h/(HR_0^2)^{1/3}, \quad x = R/R_0, \quad x_* = R_*/R_0, \tag{8}$$

we obtain from (6) and (7) the problem

$$\frac{3}{\eta^4}\frac{d\eta}{dx} = 8\left(-\frac{\mathrm{We}}{x_* - 1}\frac{1}{x^4} + \frac{1}{x^3} - \frac{1}{x^6}\right); \quad \eta \to \infty, \ x \to 1, \tag{9}$$

whose solution has the form

$$\eta(x) = \left[\frac{8}{3}\frac{\mathrm{We}}{x_* - 1}\left(1 - \frac{1}{x^3}\right) - 4\left(1 - \frac{1}{x^2}\right) + \frac{8}{5}\left(1 - \frac{1}{x^5}\right)\right]^{-1/3}. \tag{10}$$

The velocity of the center of mass is

$$-v = \frac{dz_c}{dt} = \frac{dh}{dt} + \frac{1}{2}\frac{d}{dt}\left(\frac{V}{\pi R^2}\right) = \left(\frac{dh}{dR} - 2\frac{R_0^3}{R^3}\right)\frac{dR}{dt}, \tag{11}$$

which yields an equation for determining the $R = R_*$ at which dz_c/dt, and therefore also E, vanish. By using (6) and (7), this equation can be written in the variables of (8) in the form

$$\frac{4}{3}\varkappa\left(-\frac{\mathrm{We}}{(x_* - 1)x_*} + 1 - \frac{1}{x_*^3}\right) = \frac{1}{\eta_*^4}, \quad \eta_* = \eta(x_*), \quad \varkappa = \left(\frac{H}{R_0}\right). \tag{12}$$

Equation (12) determines x_*, and thereby closes Eq. (10) for the dimensionless η of the vapor layer, describing the trajectories of the system in the phase plane.

Phase-Plane Diagram of the Collision Process. The state of the system at any instant is uniquely characterized by the values of R and h (or x and η), which depend on t as on a parameter. The trajectories of the system in the (x, η) phase plane are completely determined by Eqs. (10) and (12), which depend on the Weber number introduced in (7), and on the parameter $\varkappa$ in (12).

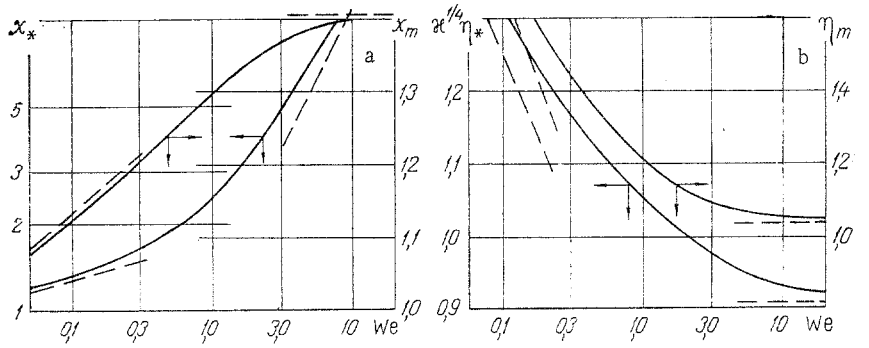

Fig. 1. a) Relative radius of vapor layer; b) its dimensionless thickness in states of maximum spreading and minimum thickness of vapor layer as functions of Weber number for $\varkappa \approx 0$; dashed lines represent asymptotic forms (13)-(16).

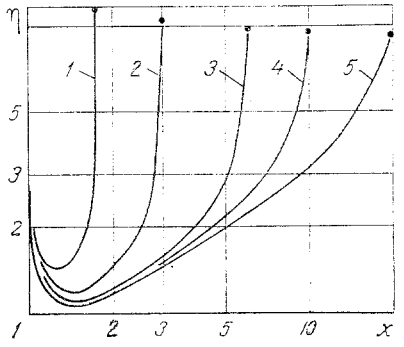

Fig. 2. Phase trajectories of the system for $\varkappa \approx 0$: 1) We ≈ 0.41; 2) 1.53; 3) 4.3; 4) 8.0; 5) 17.0; the end points of the trajectories correspond to the state of maximum spreading.

Under normal conditions, $\varkappa \sim 10^{-3} \ll 1$ for water, hydrocarbons, and most other liquids. Therefore, the value of x_* can be estimated for $\varkappa \approx 0$. Formally, the equation for it is $1/\eta_* = 0$ from (12). We obtain the following asymptotic expressions for large and small Weber numbers:

$$x_* = \frac{10}{9}\,\mathrm{We} + 1 + O\left(\frac{1}{\mathrm{We}}\right), \quad \mathrm{We} \gg 1;$$

$$x_* = 1 + \sqrt{\frac{2}{3}\mathrm{We}} + O(\mathrm{We}), \quad \mathrm{We} \ll 1. \tag{13}$$

The value $x = x_m$ for which the thickness of the vapor layer is minimum is of interest. It is obtained from the equation $d\eta(x_m)/dx_m = 0$. We have

$$x_m = 1.406 + O\left(\frac{1}{\mathrm{We}^2}\right), \quad \mathrm{We} \gg 1;$$

$$x_m = 1 + \sqrt{\frac{\mathrm{We}}{6}} + O(\mathrm{We}), \quad \mathrm{We} \ll 1. \tag{14}$$

The asymptotic forms (13) and (14) are plotted in Fig. 1a together with x_* and x_m as functions of We for $\varkappa \approx 0$. It is clear that the spreading of the drop over the vapor layer increases without bound with increasing We, but the minimum thickness of the vapor layer is reached for relatively insignificant spreading.

The dependence of η_* on We for various $\varkappa$ follows directly from (12). Using (13), we have, approximately for $\varkappa \ll 1$,

$$\eta_* = \left(\frac{3}{4\varkappa}\right)^{1/4}\left[1 + \frac{0.2025}{\mathrm{We}} + O\left(\frac{1}{\mathrm{We}^2}\right)\right], \quad \mathrm{We} \gg 1;$$

$$\eta_* = \frac{1}{\varkappa^{1/4}}\left[\left(\frac{3}{8\mathrm{We}}\right)^{1/8} + O(\mathrm{We}^{3/8})\right], \quad \mathrm{We} \ll 1. \tag{15}$$

In a similar way we obtain from (10) and (14) for the minimum dimensionless thickness of the vapor layer when $\varkappa \ll 1$

$$\eta_m = 1.048 + O(\mathrm{We}^{-2}), \quad \mathrm{We} \gg 1;$$

$$\eta_m = (2\mathrm{We})^{-1/3}[1 + O(\sqrt{\mathrm{We}})]. \quad \mathrm{We} \ll 1. \tag{16}$$

A comparison of (15) and (16) shows that the expression for η_m for We $\ll$ 1 is actually suitable only for $\varkappa^{1/4} < \mathrm{We}^{5/24}$. The values of η_* and η_m as functions of the Weber number are shown together with their asymptotic forms in Fig. 1b. The phase trajectories of the system for various We and $\varkappa \approx 10^{-4}$ are plotted in Fig. 2.

Since $d\eta/dx \neq 0$ at $x = x_*$, the derivatives dR/dt and dh/dt, together with $v = dz_c/dt$, vanish simultaneously in accordance with (11); i.e., x_* and η_* correspond to the maxima of R and h reached at time t_*. At later times the dimensionless phase trajectories are also determined by the solutions of Eq. (9) for the initial condition $\eta(t_*) = \eta_*$, but in the range $x < x_*$, and in view of the uniqueness of the solution of such a problem, coincide with the phase trajectories for times $t < t_*$. Thus, after reaching the point (x_*, η_*) in the phase plane, the system returns to the initial state $(1, \infty)$ on the same trajectory along which it reached the indicated point. The ratio of the kinetic energy of a drop moving away from the wall to the initial kinetic energy is identically equal to unity; i.e., the drop experiences an elastic rebound. Since dissipation was neglected, this is quite understandable.

It follows from Fig. 2 that a drop moving toward the wall is practically undeformed until it nears the point of closest approach. Then the deformation begins, accompanied at first by a certain thinning of the vapor layer, and later by its thickening, which is greater the smaller the value of the parameter $\varkappa$. This pattern is then repeated in the reverse order. Near the state of maximum spreading the drop "bounces" for a short time on the wall. The transformation of a drop into a very thin liquid film and the indicated bouncing are frequently observed in experiments [12-15]. The rebounding drop does not begin to break up until We $\sim$ 20 [15].

The critical conditions under which elastic rebound first occurs are of partricular interest. These conditions can be determined from the equations $\eta_m \approx \Delta$, corresponding to the concept that the drop comes into direct contact with the wall if the thickness of the vapor layer becomes less than the height of the roughness protuberances. Thus, by using (6), (7), and (16), the critical temperature drop for high and low We is

$$\Delta T_L \approx 0.87\,\frac{\sigma L \Delta^3}{\nu\lambda R_0^2}, \quad \mathrm{We} \gg 1;$$

$$\Delta T_L \approx \frac{1}{2}\,\frac{\rho L v_0^2 \Delta^3}{\nu\lambda R_0}, \quad \mathrm{We} \ll 1. \tag{17}$$

For We $\gg$ 1 the value of ΔT_L is practically independent of the velocity with which the drop impinges on the wall, and for We $\ll$ 1 it is independent of the surface tension of the liquid. Relations (17) are in qualitative agreement with experiment.

Dynamic Properties of the System. From Eq. (1) and the determination of the constant C in accordance with (5)-(7), we have

$$\frac{1}{2}\rho V \frac{d}{dt}\left(\frac{dz_c}{dt}\right)^2 = -4\pi\sigma\,\mathrm{We}\,\frac{R_0^2}{R_* - R_0}\,\frac{dR}{dt}. \tag{18}$$

Using (11), and integrating (18) with the initial condition $dz_c/dt = v_0$ at $R = R_0$, we obtain, after some simple transformations

$$\frac{dR}{dt} = \pm\left(\frac{2R_0^3}{R^3} - \frac{dh}{dR}\right)^{-1}\left(1 - \frac{R - R_0}{R_* - R_0}\right)^{1/2} v_0, \tag{19}$$

where the upper and lower signs refer to the forward and backward directions of motion along the phase trajectory. Substituting the expressions for h and dh/dR as functions of R into (19), we obtain the actual equation with separable variables for determining R(t).

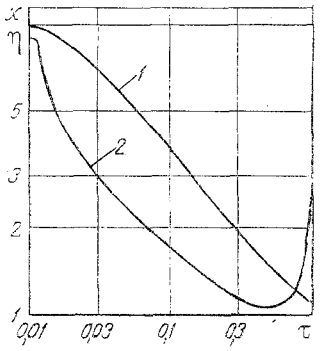

Fig. 3

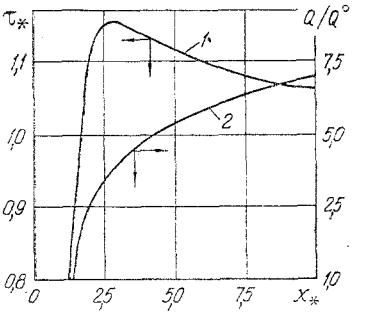

Fig. 4

Fig. 3. 1) Relative radius and 2) dimensionless thickness of vapor layer as functions of dimensionless time for $\varkappa = 10^{-4}$ and We ≈ 8.0 ($x_* \approx 10$).

Fig. 4. 1) Dimensionless collision time and 2) relative amount of heat transferred to drop as functions of maximum radius of vapor layer for $\varkappa \approx 0$.

It follows from the results presented above that dh/dR is of the order $\varkappa$, and for $\varkappa \ll 1$ it can be neglected in comparison with $R_0^3/R^3 = x^{-3}$ in (19) for all x except in the immediate neighborhoods of the points x = 1 and $x = x_*$. Considering for definiteness the backward motion along a phase trajectory, measuring time from the instant the drop is in the state of maximum spreading ($x = x_*$), and integrating (19), we obtain in the dimensionless variables of (8)

$$\frac{v_0}{R_0} t = \tau = \frac{1\sqrt{x_* - 1}}{x_*} \left[\frac{\sqrt{x_* - x}}{x} \left(\frac{1}{x} + \frac{3}{2x_*} \right) + \right.$$
$$\left. + \frac{3}{4x_*^{3/2}} \ln \frac{\sqrt{x_*} + \sqrt{x_* - x}}{\sqrt{x_*} - \sqrt{x_* - x}} \right]. \qquad (20)$$

This formula is displayed in Fig. 3, together with η as a function of τ. It can be seen that the "bouncing" of the drop near the state of maximum spreading actually occurs for a relatively short time.

The dependence of the dimensionless time τ_* for x to change from 1 to x_*, which follows from (20), is shown in Fig. 4. The quantity R_0/v_0 enters as the characteristic time of the collision process, as should be expected from dimensionality considerations. For $x_* \sim$ We $\to \infty$, $\tau_* \to 1$; for $\mathrm{We} \ll 1$ $\tau_* \approx 4(x_* - 1) \approx 3.266\sqrt{\mathrm{We}}$. The corresponding dimensional collision times are R_0/v_0 and $1.633(\rho R_0^3/\sigma)^{1/2}$; the latter agrees to within a constant with the known period of the principal mode of free vibrations of a drop.

Let us estimate the total heat flux to the drop during the time of collision, assuming that $h \approx$ const in the collision process. Using the expression for dτ/dx from (19), and neglecting dh/dr, we have

$$Q \approx \frac{\lambda \Delta T}{h} \pi \int_0^{t_*} R^2 dt = \frac{\lambda \Delta T}{h} \frac{R_0^3}{v_0} \pi \int_1^{x_*} x^2 \frac{d\tau}{dx} dx =$$
$$= 2Q^\circ \left(1 - \frac{1}{x_*}\right)^{1/2} \ln \frac{\sqrt{x_*} + \sqrt{x_* - 1}}{\sqrt{x_*} - \sqrt{x_* - 1}}, \quad Q^\circ = \frac{\lambda \Delta T}{h} \pi R_0^2 \frac{R_0}{v_0}. \qquad (21)$$

The dependence of Q/Q° on x_* is also shown in Fig. 4. For large and small We we obtain from (21)

$$Q/Q^\circ \approx 2 \ln(4x_*) \approx 2(\ln \mathrm{We} + 1.492), \quad \mathrm{We} \gg 1;$$
$$Q/Q^\circ \approx 4(x_* - 1) = 4\tau_* \approx 3.266\sqrt{\mathrm{We}}, \quad \mathrm{We} \ll 1. \qquad (22)$$

The corresponding order of magnitude estimates for Q are

$$Q \sim (\lambda \Delta T)^{2/3} \left(\frac{\sigma L}{\nu}\right)^{1/3} \frac{R_0^{7/3}}{v_0} \left(\ln \frac{\rho v_0^2 R_0}{\sigma} + 0.105\right), \quad \mathrm{We} \gg 1;$$

$$Q \sim (\lambda \Delta T)^{2/3} \left(\frac{L}{\nu}\right)^{1/3} \frac{\rho^{5/6}}{\sigma^{1/2}} R_0^{19/6} v_0^{2/3}, \quad \mathrm{We} \ll 1. \tag{23}$$

The dependence of Q on ΔT, the volume of the drop, and the physical parameters is confirmed qualitatively in most experiments. With increasing R_0 the value of Q for small We is practically proportional to the volume of the drop, which agrees with the important practical conclusion that the heat removed from a hot surface is proportional to the spray density [11]. According to (23), as v_0 is increased, the flux Q at first increases and then begins to decrease. Unfortunately, the data obtained in different experiments for the thermal effect of a shock interaction of a drop with a superheated wall differ even in order of magnitude [12-14]. Therefore, strict comparison of theory and experiment is impossible, not only because of the severe assumptions made in deriving (22) and (23), but also because of insufficient experimental accuracy.

As a result of not taking account of the unsteady nature of the heat-conduction process, the heat flux to a drop as it approaches the surface turns out to be somewhat smaller than the flux to it as it moves away. Therefore, the resisting force experienced by the drop as it approaches the surface is smaller than the force accelerating it as it moves away. As a result, the kinetic energy of the drop after collision should be larger than the initial energy, which actually occurs if the velocity v_0 is not too large [12]. On the other hand, the varying thickness of the vapor layer and its effect on the dynamics of the drop give rise to an additional resisting force which opposes the motion of the drop both toward and away from the wall. This effect contributes to the speeding up of the time rate of change of R and h while the drop is approaching the wall, and to the slowing down of this change as it recedes. The retardation of the drop is thus more rapid than its acceleration. As a result, the time dependences of x and η are not symmetrical with respect to the time $t = t_*$, and the phase trajectory of the withdrawal of the drop from the wall does not coincide with its trajectory of approach. Accordingly, the drop loses kinetic energy in the collision process. All the characteristics of the process noted are confirmed by experiment [12].

In conclusion, we note that obtaining reliable results which might be used with confidence in engineering calculations involves not only taking account of unsteady and dissipative phenomena and introducing radiation heat flux, which is important for high wall temperatures, etc., but also taking account of collective effects which can severely alter the expected cooling pattern. In addition to the purely hydrodynamic aspect of the effect of the constraint of the motion — the methods for investigating which were developed in [20], it is important to take account of collisions of drops in the immediate vicinity of the wall, which on the whole facilitate their direct contact with the wall, and also the effect of the flux of vapor from the wall, which is capable in a number of cases of slowing down drops as they approach it [21].

NOTATION

δA is the elementary work; C, constant introduced in (6) and (7); E, kinetic energy of drop; F, free energy of surface tension; f, pressure force; H, characteristic length scale introduced in (6); h, thickness of vapor layer; j, mass of liquid evaporating per unit time from unit area of evaporation surface; L, heat of vaporization; p, pressure; Q, Q^0, total amount of heat transferred to drop in collision process, and its characteristic value; q, heat flux to unit area of evaporation surface; R, radius of vapor layer; r, radial coordinate; T_s, T_w, saturation temperature and wall temperature; $\Delta T = T_w - T_s$; t, time; V, volume of drop; v, velocity of center of mass of drop; x, relative radius of vapor layer; z_c, coordinate of center of mass of drop; Δ, height of wall roughness; η, dimensionless thickness of vapor layer; $\varkappa$, parameter defined in (12); λ, thermal conductivity, ν, kinematic viscosity; ρ, density of liquid; σ, surface tension; τ, dimensionless time; ψ, function determining shape and position of drop; and We, Weber number defined in (7). Subscripts 0 and m denote the initial state and the state with minimum thickness of vapor layer, respectively; * denotes state of maximum spreading.

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