INSTABILITY OF A LIQUID SURFACE WITH LASER HEATING

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UDC 530.182

We consider two mechanisms for occurrence of surface instability for a liquid with heating by means of laser radiation. Nonuniform heating along the liquid surface leads to thermocapillary instability. In contrast to other works a nonstationary heating regime is studied for the substance and estimates are provided for the minimum threshold for occurrence of instability of this type. With evaporation of the substance nonuniform heating of the liquid through the depth leads to occurrence of instability of another type. In this case estimates are obtained for the maximum increment in increasing surface disturbances with the minimum of assumptions.

1. Surface waves are excited with the action of powerful laser radiation on the free surface of a liquid. Excitation of capillary waves by a pressure output pulse is considered in [1], stimulated by laser radiation under conditions when the duration of the effect is small compared with the period of the capillary wave. A reverse situation is also possible when excitation of surface waves occurs as a result of development of instability during interaction of laser radiation with a substance. Here a build-up of surface disturbances may markedly change conditions for the absorption and scattering of light and affect development of the process of light interaction with a liquid [2]. One of the possible mechanisms for instability is connected with nonuniform heating of the surface layer. With nonuniform temperature distribution along the surface a thermocapillary force arises [3] proportional to the gradient of surface tension α. This tangential force operates in the surface layer of a substance and causing movement of it leads to formation of craters, i.e., deflection of the liquid surface in an area of increased temperature [4].

As is shown in [5], with nonuniform heating of a liquid at its surface two types of waves may be activated: gravitational-capillary and thermocapillary waves. Occurrence of instability of the first type is connected with distortion of the shape of the interface of two media $z = \zeta(x, t)$, and below the main attention will be devoted to self-activation of gravitational-capillary waves.

It is well known [6] that a change in mechanical vibration energy of an arbitrary type is described by the equation

$$\frac{dE_m}{dt} = -T\int \rho\dot{s}\,d\mathbf{r} + \langle vf\rangle, \tag{1.1}$$

where $E_m = \langle\int \rho v^2 d\mathbf{r}\rangle$ is vibration energy (in the case surface waves) averaged for vibration period $2\pi/\omega$. The energy is worked out most simply in an approximation of a low-viscosity liquid when movement of the substance may be considered almost potential. In this case liquid movement velocity v is determined exclusively by the scalar potential and the energy of surface vibrations arriving in a unit area equal to $\rho\omega^2|\zeta_0|^2/2k$ [ρ is liquid density, k is wave number of a propagating surface wave $\zeta(x, t) = \zeta_0 \exp(ikx - i\omega t)$].

If an increase in entropy of the substance $\dot{s}$ depends mainly on viscous losses, then the first term to the right in equality (1.1) is transformed to the expression $(1/2)\left\langle\int \sigma_{ik}^2\,d\mathbf{r}\right\rangle$ and with a plane monochromatic surface wave it equals $2\eta k\omega^2|\zeta_0|^2$ [3] (σ_{ik} is viscous stress tensor, $\eta = \nu\rho$ is dynamic viscosity). However, a change in vibration energy E_m in the general case may occur both as a result of an increase in entropy of the substance, and also as a result of completing additional work $\langle vf\rangle$. The same as in stationary flow, a build-up of vibrations is possible due to the energy of external flow reciprocation [6], and as in the case of a thermal source there is possible self-activation of surface waves due to the

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Moscow. Translated from Zhurnal Prikladnoi Mekhaniki i Tekhnicheskoi Fiziki, No. 3, pp. 28-34, May-June, 1991. Original article submitted February 17, 1989; revision submitted December 4, 1989.

work of thermocapillary force proportional to $\partial\alpha/\partial x$ (x is coordinate along the liquid surface). Assuming that the surface tension coefficient is a well-defined function of temperature and with an increase in it $\alpha(T)$ falls, we find that the power of this force referred to a unit of surface S is $\langle v_x f\rangle/S = \left|\frac{d\alpha}{dT}\right| \left\langle v_x\left(\frac{\partial T}{\partial x} + \zeta_x' \frac{\partial T}{\partial z}\right)_{z=\zeta}\right\rangle$. Here by substituting $v_x(z = \zeta) = -\omega\zeta$ and amplitude values of temperature we find that in a quadratic approximation with respect to ζ the power sought is

$$\langle v_x f\rangle = \frac{\omega k S}{2}\left|\frac{d\alpha}{dT}\right| \mathrm{Re}\left(iT_1\zeta_0^*\right) \tag{1.2}$$

(T_1 is amplitude of liquid temperature oscillation linear with respect to ζ).

In the case of a stationary flow of heat, zero with respect to ζ, the right-hand part of Eq. (1.1) may be reduced to the expression γE_m, where the value of γ which is constant with respect to time has the physical meaning of a decrement (or increment with $\gamma > 0$) of vibrations. By equating γ to zero it is possible to find the threshold for occurrence of instability [5, 7]. However, the calculation scheme provided relates to the case of stationary distribution of temperature in the liquid surface layer. In contrast to this we consider the question of self-activation of surface disturbances under conditions when the distribution of temperature is found from solving the problem of heating the surface layer of a liquid by laser radiation, and as we shall see it appears to be exclusively nonstationary.

2. Let at the surface of a liquid from direction $z < 0$ radiation fall with intensity $I(t)$ and it be absorbed in a thin layer of substance with thickness $\sim 1/\mu$. For simplicity we shall assume that the reflection coefficient for radiation equals zero, and the diameter of the laser spot d is large compared with thickness $1/\mu$. The temperature distribution in the liquid is described by the normal thermal conductivity equation

$$\rho c_p\left(\frac{\partial T}{\partial t} + v\nabla T\right) = \varkappa\Delta T + \mu I(t)\exp\{-\mu[z - \zeta(x, t)]\}, \quad z \geqslant 0, \tag{2.1}$$

where c_p is liquid specific heat capacity; $\kappa = \chi\rho c_p$ is its thermal conductivity coefficient. The condition $d \gg 1/\mu$ makes it possible in an approximation of a plane phase interface $\zeta = 0$ to consider the "zero" problem to be unidimensional and in particular to substitute the Laplace Δ in (2.1) by $\partial^2/\partial z^2$. With a moderate intensity of electromagnetic radiation when heating of the substance is not accompanied by evaporation of it, Eq. (2.1) is supplemented by a boundary condition

$$(\mathbf{n}\nabla T)_{z=\zeta(x,t)} = 0 \tag{2.2}$$

[$\mathbf{n}$ is the normal to surface $z = \zeta(x, t)$]. We note that condition (2.2) is accurate for a liquid bounded by a vacuum, and it is approximate for a liquid bounded by a real gas. Since gas thermal conductivity κ' is normally small, then heat flow from the direction of a gas $\kappa'(\nabla T')_{z=\zeta}$ in the right-hand part of (2.2) may be ignored.

By placing in Eq. (2.1) a disturbance ζ equal to zero, and also $v = 0$, and substituting Δ by $\partial^2/\partial z^2$ we find the zero solution of thermal conductivity equation (2.1). Here boundary condition (2.2) is simply reduced to equality to zero of derivative $\partial T_0/\partial z$ at surface $z = 0$. In this case the solution of the problem may be written in quadratures:

$$T_0(z, t) = T_\infty + \frac{\mu}{\rho c_p}\int_0^t dt' I(t')\exp\left(\frac{t - t'}{t_1}\right)\int_{(t-t')/t_1}^{\infty}\frac{d\eta}{\sqrt{\pi\eta}}\exp\left(-\frac{z^2\mu^2}{4\eta} - \eta\right). \tag{2.3}$$

Here T_∞ is initial liquid temperature; $t_1 = 1/\chi\mu^2$ is time (in which the thermal wave reaches the "end" of the absorption zone) determined from the equality $\sqrt{\chi t_1} = 1/\mu$. Analysis of expression (2.3) shows that the stationary regime in the unidimensional case and with absence of liquid evaporation is almost never accomplished. In fact, from solution (2.3) obtained it follows that at first (with $t \ll t_1$) the temperature at each point of space increases in proportion to the amount of heat released and it is approximately equal to

$$T_0(z, t) = T_\infty + (\mu/\rho c_p)\exp(-\mu z)\int_0^t dt' I(t').$$

With constant radiation intensity this means that temperature T_0 increases linearly with time. With $t \gtrsim t_1$ thermal conductivity has a marked role. Heat starts to flow from the liquid surface in porportion to $\sqrt{t}$, and the increase in T_0 with time slows down somewhat. In the limiting case $t \gg t_1$ and with the condition that the laser pulse was deexcited, the temperature T_0 at each point of space by reaching a maximum falls to the initial value T_∞. Asymptotic (2.3) acquires a very simple form and T_0 changes by the rule

$$T_0(z, t) = T_\infty + \frac{Q/\rho c_p}{\sqrt{\pi\chi t}}\exp\left(-\frac{z^2}{4\chi t}\right) \tag{2.4}$$

$\left(Q = \int_0^\infty dt' I(t')\right.$ is total laser radiation energy absorbed in the liquid).

Whence it follows that the normal approach to the problem of occurrence of instability of thermocapillary waves does not operate here. Calculation of the power of thermocapillary force in this case by Eq. (1.2) is still possible in an approximation of "frozen" temperature T_0. For this it is necessary that the time for existence of a change in T_0 is much greater than the inverse vibration frequency. Here the power of thermocapillary force $\langle v_x f\rangle$, averaged with respect to vibration period $2\pi/\omega$, appears to be a function of "slow" time. As a consequence of this there is a function of time and vibration decrement γ, and with it also combined frequency corresponding to a dispersion equation since $\mathrm{Im}\,\omega(k) = \gamma(t)/2$. Thus, finding the threshold for occurrence of instability of thermocapillary forces according to a traditional scheme causes certain difficulties. In view of this we discontinue the search for an accurate value of vibration decrement $\gamma(t)$ as a function of time, and by considering the problem exclusively close to the instability threshold $\gamma = 0$ we limit ourselves to finding the minimum value of laser radiation starting from which occurrence of instability is possible.

As can be seen from Eq. (1.2), in order to find the power of thermocapillary force it is necessary to determine the value of temperature T_1 oscillation amplitude, linear with respect to ζ at point $z = 0$. By linearizing the thermal conductivity Eq. (2.1) it is possible to find the corresponding value for T_1. Here boundary condition (2.2) after linearizing becomes nonuniform: $(\partial T_1/\partial z)_{z=0} = -\zeta(x, t)(\partial^2 T_0/\partial z^2)_{z=0}$.

By introducing subsidiary function $\theta(x, z, t) = T_1(x, z, t) + \zeta(x, t)(\partial T_0/\partial z)$ the problem may be reduced to solving an inhomogeneous equation for θ with a similar boundary condition of the previous form: $(\partial\theta/\partial z)_{z=0} = 0$. Here the equation for subsidiary function θ will be as for a normal thermal conductivity equation with a source:

$$\frac{\partial\theta}{\partial t} = \chi\left(\frac{\partial^2\theta}{\partial x^2} + \frac{\partial^2\theta}{\partial z^2}\right) + \left(\dot{\zeta} - v_z - \chi\frac{\partial^2\zeta}{\partial x^2}\right)\frac{\partial T_0}{\partial z}. \tag{2.5}$$

In solving Eq. (2.5) we shall proceed from the fact that activation of the surface $\zeta(x, t)$ has the form of the same plane monochromatic wave $\zeta_0\exp(ikx - i\omega t)$. Therefore, the dependence of temperature oscillation amplitude $\theta_1(x, z, t)$ on longitudinal coordinate x is determined by the same exponential multiple $\exp(ikx)$ as that for $\zeta(x, t)$, and $\partial^2\theta/\partial x^2 = -k^2\theta$, $\partial^2\zeta/\partial x^2 = -k^2\zeta$. However, the dependence of $\theta(t)$ on time is more complicated than for $\zeta(x, t)$, and it cannot be reduced to a similar time multiple $\exp(-i\omega t)$. In addition, in solving Eq. (2.5) as previously we shall assume for simplicity that liquid movement is potential and therefore velocity v_z equals $\dot{\zeta}\exp(-kz)$. Equation (2.5) together with its boundary condition is resolved accurately and its solution may be written in quadratures. Since we are interested in the value of temperature T_1 only at zero, then here we provide an expression for $\theta_1(z = 0, t) = T_1(z = 0, t)$ written with $z = 0$:

$$T_1(x, z=0, t) = 2\chi k^2 \int_0^\infty \frac{dt'\zeta(x, t')}{\sqrt{4\pi\chi(t-t')}} \exp[-\chi k^2(t-t')] \int_0^\infty dz' \frac{\partial T_0}{\partial z'} \left\{1 - \frac{i\omega}{\chi k^2} + \frac{i\omega}{\chi k^2}\exp(-kz')\right\} \exp\left[-\frac{z'^2}{4\chi(t-t')}\right]. \quad (2.6)$$

The gradient of "zero" temperature in the inner integral of (2.6) may be found by means of accurate solution of (2.3) for temperature T_0. By substituting this value of $\partial T_0/\partial z'$ in Eq. (2.6) we also obtain an accurate expression for linear temperature T_1. However, this expression for T_1 is very complicated and cumbersome, and it is expressed in the general case in terms of a fourfold integral. In order to facilitate analyzing the behavior of $T_1(x, z = 0, t)$ with time we use the following approximation. First, in finding the derivative $\partial T_0/\partial z$ we use for $T_0(z, t)$ not accurate Eq. (2.3), but its asymptotic (2.4). In addition, in calculating the inner integral in Eq. (2.6) the linear exponent $\exp(-kz')$ in braces is substituted by the quadratic component $(-k^2z'^2)$. All of this makes it possible to carry out integration in elementary functions by retaining sufficient accuracy for the calculations and the qualitative nature of the behavior of temperature T_1 with time.

In carrying out integration in (2.6) taking account of the comments made we find that

$$T_1(x, z=0, t) = -\frac{k^2 Q\zeta}{\pi\rho c_p}\Psi\left(\omega t, \frac{\chi k^2}{\omega}\right), \quad (2.7)$$

where dimensionless quantity Ψ is a function of two dimensionless parameters ωt and $\chi k^2/\omega$. In calculating the power of thermocapillary force the value of Ψ itself is not necessary, but only its imaginary part

$$\operatorname{Im}\Psi = \int_0^1 d\xi \sqrt{\frac{\xi}{1-\xi}} \sin(\omega t\xi)\, e^{-\chi k^2 t\xi} - 4\omega t \int_0^1 d\xi \frac{\xi^{3/2}(1-\xi)^{1/2}\cos(\omega t\xi)}{1 + 4\chi k^2 t\xi(1-\xi)} e^{-\chi k^2 t\xi}. \quad (2.8)$$

Now we consider the behavior of $\operatorname{Im}\Psi$ as a function of ωt. Here we note that parameter $\chi k^2/\omega$ reaches its greatest value with limiting values of $\omega_{\max}$ and $k_{\max}$ which are found from the condition that $\operatorname{Im}\omega \simeq \operatorname{Re}\omega$. Taking account of the dispersion relationship for gravitational-capillary waves it follows that the greatest value of $(\chi k^2/\omega)_{\max}$ is $\simeq\chi/2\nu$, which for example for water is $\simeq 0.08$ in the temperature range up to 100°C. Therefore, in the region for a change in variable ωt from zero to a value of the order of one the combination $\chi k^2 t\xi$ and $4\chi k^2 t\xi(1 - \xi)$ also appeared to be of small value and it is possible to ignore them in Eq. (2.8). In this case expression (2.8) is considerably simplified and it may be rewritten in the form

$$\operatorname{Im}\Psi(\omega t) \simeq \int_0^1 d\xi \sqrt{\frac{\xi}{1-\xi}}(7 - 8\xi)\sin(\omega t\xi). \quad (2.9)$$

By directing argument ωt to zero we find that $\operatorname{Im}\Psi(0) = \omega t\pi/8$. With an increase in ωt the linear growth of (2.9) slows down and somewhere in the region of $\omega t \simeq \pi/2$ the imaginary part of Ψ reaches a maximum. Then there is a drop in $\operatorname{Im}\Psi$ to zero after which the value in question becomes negative. With a further increase in the argument there oscillation of $\operatorname{Im}\Psi$ about zero and the vibration amplitude itself first falls as $(\omega t)^{-3/2}$, and then in an exponential way. Below we shall be interested in the greatest positive value of $\operatorname{Im}\Psi$. Simple estimates of expression (2.9) show that $(\operatorname{Im}\Psi)_{\max}$ does not exceed $\simeq 0.5$. In addition, it is noted that if for the zero temperature the accurate expression (2.3) is used and not its asymptotic (2.4), then the imaginary part of Ψ appears to be a function of yet another additional variable t/t_1. However, it is possible to show that with $t \lesssim t_1$, $\operatorname{Im}\Psi$ in modulus is much less than the maximum value found above. Therefore, the maximum $\operatorname{Im}\Psi$, equal approximately to $\simeq 0.5$, appears to be correct for the whole region of the change in time t. Physically this analysis corresponds to the fact that the oscillation amplitude for temperature T_1 at first increases from zero, then it reaches its maximum value, and then it falls again.

Now we substitute expression (2.7) for linear temperature in Eq. (1.2) and we find the average power of thermocapillary force. As might be expected, $\langle v_x f\rangle$ appears to be a function of time. Without going into a detailed study of the behavior of the whole picture of the phenomenon in time we make a rough estimate of stability in relation to occurrence of surface waves. For this we equate the maximum value $\langle v_x f\rangle_{max}$ to losses in viscosity and we find the minimum threshold for self-activation of capillary waves

$$Q_{\min} = \frac{8\pi\rho^2 \nu c_p \omega}{k^2\,|\,d\alpha/dT\,|\,(\mathrm{Im}\,\Psi)_{\max}}. \tag{2.10}$$

From the relationship obtained it follows that with $Q < Q_{min}$ liquid heating by laser radiation occurs stably and activation of surface waves is only possible with $Q > Q_{min}$ or significantly greater. If the estimate $(\mathrm{Im}\,\Psi)_{max} \simeq 0.5$ found is used and we draw attention to dispersion relationship $\omega^2 \simeq \alpha k^3$, then relationship (2.10) may be rewritten in the form

$$Q_{\min} \simeq 16\pi\rho^2 \nu c_p \sqrt{\frac{\alpha}{k\rho}}\left|\frac{d\alpha}{dT}\right|^{-1}. \tag{2.11}$$

In contrast to [5] relationships (2.10) and (2.11) superimpose limits not on the intensity of laser radiation I(t), but on the total energy of absorbed radiation. Since the power of thermocapillary force is determined exclusively by the temperature reached at the liquid surface, then in the stationary case, which was only considered in [5], temperatures T_0 and T_1, and with also power $\langle v_x f\rangle$, appeared to be functions of constant radiation intensity. In the nonstationary case this is not so and temperature, and with it the power of the force $\langle v_x f\rangle$ sought, are not determined by the intensity, but the whole radiation energy absorbed to instant of time t: $Q_t = \int_0^t dt'\, I(t')$.

As an example of using Eqs. (2.10) and (2.11) we provide a numerical value of threshold Q_{min} with heating of water by a CO_2-laser radiation with a wavelength of 10.6 μm. Data for water are taken for a temperature in the region ≃50°C: $c_p \simeq 4.2$ J/g, $\alpha \simeq 70$ mJ/m², $\nu \simeq 0.01$ cm²/sec, $|d\alpha/dT| \simeq 0.17$ mJ(m²·deg), $\mu \simeq 870$ cm⁻¹. In this case the minimum threshold for occurrence of instability will be proportional to the square root of the disturbance length and equal $Q_{min} \simeq 90/\sqrt{k}$ J/cm².

3. When temperature T_0 of the liquid surface reaches the boiling temperature, then apart from the mechanism of surface wave activation considered yet another mechanism of surface instability arises. The new mechanism is connected with nonuniform temperature distribution, but not along the liquid surface as previously, but with depth. Instability of this type has been studied in detail in [7]. With evaporation of substance boundary condition (2.2) changes radically and in ignoring thermal flow in the direction of the gas $\kappa'(\nabla T')_{z=\zeta}$ it is written in the form

$$\varkappa(\mathbf{n}\nabla T)_{z=\zeta(x,t)} \simeq jL. \tag{3.1}$$

Here L is heat of evaporation; j is flow of evaporating mass. From equality (3.1) it follows that with evaporation of a substance the maximum temperature shifts from the surface of the liquid into the depth and boundary condition (3.1) is now written for the moving surface $z = z_0(t)$ even in the absence of surface waves. Here the rate of change in the position of the boundary as a result of evaporation is connected with mass flow as $j \simeq \rho\dot{z}_0$. As far as the value of the mass flow j itself is concerned, then it may be found in the simplest case from the Hertz–Knudsen equation

$$j = \frac{p_s(T_0) - p'}{\sqrt{2\pi RT}}\,\beta, \tag{3.2}$$

where β is condensation coefficient; $p_s(T_0)$ is saturated vapor pressure at temperature T_0 at the liquid surface; R is gas constant.

Nonuniform temperature distribution through the depth of a liquid leads to the situation that with random deflection of the phase interface there is a reduction in boundary layer thickness, i.e., the length over which a marked change in temperature occurs. Then according to equality (3.1) there is an increase in mass flow as a result of additional evaporation of the substance and as a consequence of this deflection of the surface increases even more. Solution of thermal conductivity Eq. (2.1), taking account of new boundary condition (3.1), makes it possible in principle to find the increment for increasing surface disturbances as a function of disturbance length. The dependence $\gamma(k)$ in a stationary evaporation regime (I = const) was found in [7]. However, some simple estimates of γ may be made in this case without recourse to complex calculations and serious limitations.

As follows from Eq. (3.2), the velocity of the boundary $\dot{z}_0(t)$ is a function of temperature T_0 of the liquid surface, i.e., $\dot{z}_0 = \Phi(T_0)$. It follows from solving the set of Eqs. (2.1), (3.1), and (3.2) that temperature T_0 itself should be a function of $z_0(t)$. Therefore, for a surface disturbance ζ of infinite length, i.e., random deflection $\delta z_0(t)$, the corresponding equation is found simply as a result of expanding function $\Phi(T_0(z_0))$ into a series with respect to δz_0:

$$\frac{d}{dt}\,\delta z_0 = \frac{d\Phi}{dT_0}\left(\frac{\partial T_0}{\partial z} + \frac{1}{\dot{z}_0}\frac{\partial T_0}{\partial t}\right)\delta z_0 \tag{3.3}$$

($\Phi = j/\rho = \beta[p_S(T_0) - p']/\sqrt{2\pi R T_0 \rho^2}$). Considering that derivative $\partial T_0/\partial t$ is small compared with $\dot{z}_0(\partial T_0/\partial z)$, and also by using boundary condition (3.1) for finding spatial derivative $(\partial T_0/\partial z)_{z=z_0}$, from Eq. (3.3) we obtain the expression sought for the increment

$$\gamma \simeq \frac{\dot{z}_0^2 L^2 \rho'}{p_s(T_0)\,\chi c_p T_0}\left(1 - \frac{p_s(T_0)}{2\rho' L}\right). \tag{3.4}$$

Here ρ' is vapor density; in addition, in deriving expression (3.4) use was made of the Clausius—Clapeyron equation: $dp_S/dT \simeq \rho' L/T_0$. Up to the temperature for explosive boiling T_L [8] the multiple curved brackets in the right-hand part of (3.4) with sufficient accuracy equals one. In this case expression (3.4) coincides with the maximum increment $\gamma(k)$ obtained in [7] assuming a steady-state stationary evaporation regime with I = const. Further expansion of function Φ into a series with respect to δz_0 with an accuracy to square terms makes it possible to find coefficient γ_2 with δz_0^2 which appears to be negative. This indicates that the instability in question is mild in nature and after all there is establishment of a finite amplitude for increasing disturbances with $(\delta z_0)_{\max} = \gamma/\gamma_2$.

As can be seen from expression (3.4), increment γ increases with an increase in surface evaporation rate $\dot{z}_0$, which in turn is affected by the proportional intensity of radiation (in a quasistationary evaporation regime there is the equality $I = jL_{ef}$ [9]). Increment γ acquires the greatest value when T_0 reaches the explosive boiling temperature with which $\dot{z}_0$ is at a maximum with surface boiling. Thus, for water from Eq. (3.4) it follows that with $T_L \simeq 574$ K ($p \simeq 0.1$ MPa), $(\dot{z}_0)_{\max} \simeq 40$ m/sec and $I_{\max} \simeq 2.5\cdot10^5$ W/cm², the maximum increment $\gamma_{\max} \simeq 3\cdot10^6$ sec^{-1}. With the condition $1/\gamma > \tau$ this requires absorption of energy $I/\gamma \sim 0.1$ J/cm². With a further increase in intensity and energy of laser radiation liquid surface evaporation changes to volumetric and the instability of the phase interface acquires an absolute character.

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ELECTRIC-CURRENT-INDUCED RESTRUCTURING OF THE POROUS SPACE IN A MEDIUM

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UDC 534.222

It was shown in [1-3] that the transmission of an electric current through a porous saturated medium can be reduced to a change in its permeability and electrical conductivity. These changes are due to the restructuring of the porous space in a medium as a result of localization of the density of energy release in thin capillaries, limiting the conductivity of the medium. The structure of the porous space in rocks can be described within the framework of the percolation model of an inhomogeneous medium [4], using the capillary probability density function (CPDF) with respect to effective capillary radius. Here we report the results of an experimental study of the changes that occur in the structure of the porous space and, hence, in the permeability of sandy-argillaceous rocks when an ac electric current is passed through them.

The principal complication in carrying out such an experiment is that methods that would not affect the structure of the porous space of the medium must be developed for determining the CPDF. Accordingly, we used the electroporosimetric method developed especially for these purposes [5]; the essence of this method is as follows. The sample of rock under study is placed in a vertical tube (Fig. 1), whose end is immersed in a vessel of liquid; the steady-state distribution of saturation established in the medium by the capillary forces decreases as the height h of the liquid increases. As a result the electrical resistivity $\rho_e(h)$ of the rock decreases as h rises. If we measure $\rho_e(h)$ by the method of [5] at low voltages that do not result in a change in the CPDF, therefore, we can determine the CPDF before and after the passage of an electric current through the porous saturated medium.

As the electrically conducting liquid we used a $CaCl_2$ solution, which does not enter into an intensive ion-salt exchange with the rocks studied. The concentration of the solution was 0.1 N in all the experiments. Each tube of diameter 3 cm and length 1 m was filled with portions of dry rock, gradually permeated with solution entering through the end of the tube from the communicating vessel so that its level varied simultaneously with the level of the rock. We applied voltage to the outside electrodes to measure $\rho_e(h)$ and then measured $\Delta\varphi(h)$, the difference between the bottom electrode and electrodes placed equidistantly along the height of the tube. Knowing $\Delta\varphi(h)$, the current flowing in the tube, the cross-sectional area of the tube, and the distance between neighboring electrodes, we can find $\rho_e(h)$, the resistivity of the segment of the rock at a height h.

The experiment on electric treatment of sandy-argillaceous rock consisted of two stages. In the first stage the solution was filtered in the tube for 5-15 days (depending on the rock) until steady-state filtration was established. On the basis of the results of measurements of the liquid flow rate at a constant pressure gradient in the tube we measured K, the permeability of the medium as well as $\rho_{e0}(h)$, the resistivity of segments of the rock when the sample is completely saturated with the liquid. In the second stage, the flow of liquid through the tube was stopped until a steady-state saturation distribution was established, $\rho_e(h)$, was measured, and the CPDF was determined. Liquid was then again passed through the tube and the stability of K was checked.

Moscow. Translated from Zhurnal Prikladnoi Mekhaniki i Tekhnicheskoi Fiziki, No. 3, pp. 35-37, May-June, 1991. Original article submitted December 15, 1988; revision submitted August 7, 1989.