produce nonlinear distortion of a small-amplitude harmonic wave. However, one should take special care here, because the formal use of Eq. (8) (quadratic approximation) without taking the smallness of $|Y|$ into account can result in significant distortions of the solution. For example, a false soliton is exhibited when $H = 3[D/(1+\gamma)]^3$, but in this case $Y_2 = 2D/(1+\gamma) \geqslant (2/3)(c_1/c_2)^2 \sim 1$ and Eq. (8) becomes invalid. In conclusion, explaining the nature of the solutions obtained, we note that the excess pressure p on the right-hand side of Eq. (2) can be expressed with the help of the equation of state in the form of the equation $p = \rho_0 nVc^2 c_0^2/[(1-z)c^2 - 1]$, which results from Eq. (1) for the case of a stationary wave. As is evident, the sign of p/V varies as a function of the relation between c^2 and c_1^2, and when $c^2 > c_1^2$, the excess pressure p increases as V increases, while when $c^2 < c_1^2$, a negative value of p (the pressure decreases) corresponds to an increase in V. The first case corresponds to pulsations of the bubble at frequencies higher than resonance, when the bubble represents a massive impedance and the liquid is elastic; at the same time, an increase in the volume of the bubble is accompanied by compression of the elastic element (a pressure increase). In the opposite case, when $c^2 < c_1^2$, the bubble is the elastic element, and its increase implies a dilatation of the elastic element (i.e., a pressure decrease). We note that the elasticity of the gas in the bubble, which is described by the first term on the right-hand side of Eq. (2), always opposes the expansion, i.e., this term is always negative upon an increase in V. Therefore, when $c^2 < c_1^2$ in the case $(-p) > 0$, the terms on the right-hand side of the equation have unlike signs and compensation of them is possible, which corresponds to the formation of a soliton.

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ASYMPTOTIC ANALYSIS OF THE PROBLEM OF IGNITION OF REACTIVE MATERIAL BY A HEATED SURFACE

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UDC 536.46

INTRODUCTION

Due to the Arrhenius dependence of the rate of a chemical reaction on temperature in the statement of many problems of macrokinetics, several relaxation lengths (usually two) are present whose ratio forms a small parameter (for example, the ratio of the chemical reaction and heating zones). Problems of this type pertain to special perturbation problems, for whose solution the method of spliced asymptotic expansions (SAE) is most suitable. The solution of a number of steady-state problems of slow burning and detonation (see [1] and the bibliography in it) has been found with the help of SAE. The attempt to apply SAE to problems of macrokinetics formulated within the framework of partial differential equations* is still very limited [1-3]. Upper and lower limits are found in this paper for the heating time in

*V. S. Berman, "Some problems in the theory of the propagation of a zone with exothermic chemical reactions in gaseous and condensed media." Dissertation in Competition for the Scientific degree of Candidate of Physico-Mathematical Sciences, Institute of the Problems of Mechanics, Academy of Sciences of the USSR, Moscow (1974).

Tomsk. Translated from Zhurnal Prikladnoi Mekhaniki i Tekhnicheskoi Fiziki, No. 6, pp. 96-102, November-December, 1976. Original article submitted December 2, 1975.

the thermal theory of ignition by a heated surface; a comparison is given of the analytic formulas with a numerical calculation by a computer; the asymptotic nature of B. Zel'dovich's formula [4] is established; the divergence of the results (in the case in which the temperature head θ_0 tends to infinity) is shown [5, 6].

1. Qualitative Estimates. For the assumptions usually used in ignition theory, the nonlinear thermal conductivity equation

$$\partial T/\partial t = \varkappa\partial^2 T/\partial r^2 + (Qz/c)\exp(-E/RT), \quad t > 0, \quad 0 < r < \infty \tag{1.1}$$

with the conditions

$$T(r, 0) = T_-, \ T(0, t) = T_+, \ \partial T(\infty, t)/\partial r = 0 \tag{1.2}$$

is the reference.

In agreement with [4], the heating time t_0 is defined as the time to reach thermal equilibrium between the reacting medium and the heated surface:

$$\partial T(0, t_0)/\partial r = 0.$$

The ignition time t_I is comprised of the heating time and the chemical induction time:

$$t_I = t_0 + t_{in}$$

and

$$t_{in} \ll t_0, \ t_I \approx t_0$$

only in case of large $\theta_0 = E(T_+ - T_-)/RT_+^2$.

The following notation was employed in the exposition above: T is the temperature, $\varkappa = \lambda/c\rho$ is the thermal diffusivity, λ is the thermal conductivity, c is the heat capacity, ρ is the density, Q is the thermal effect, z is the preexponent, E is the activation energy, r is the spatial coordinate, t is the time, T_- is the initial temperature, and T_+ is the temperature of the hot surface.

Let us cite some qualitative estimates which reveal the characteristic scales both for the dependent and the independent variables.

If $T_+ \gg T_-$, and $E \gg RT_+$, then the contribution of the nonlinear term in (1.1) is significant only in the vicinity of r = 0, when T lies in the interval $T_+ - RT_+^2/E \leqslant T \leqslant T_+$. Therefore, it is possible to distinguish two spatial regions of the solution: the chemical width r_{ch}, where heat generation is significant, and the thermal width r_q, where the reaction is practically inactive.

Writing the balance equation of N. N. Semenov for the reaction zone, we obtain

$$\rho Qzr_{ch}\exp(-E/RT_+) = q_*, \tag{1.3}$$

where q_* is the specific thermal flux at the coupling boundary of the reaction and heating zones.

The order of magnitude of the quantities r_{ch} and r_q is estimated in the following way:

$$r_{ch} \sim \lambda RT_+^2/Eq_*, \ r_q \sim \lambda(T_+ - T_-)/q_*, \tag{1.4}$$

so that the ratio

$$r_{ch}/r_q = 0(\varepsilon), \ \varepsilon = RT_+^2/E(T_+ - T_-) = 1/\theta_0 \ll 1 \tag{1.5}$$

is a small parameter of the original problem (1.1) and (1.2). In agreement with dimensionality considerations,

$$r_q \sim \sqrt{\varkappa t_0}. \tag{1.6}$$

Equations (1.3)-(1.6) lead to the obvious estimates

$$t_0 \sim \frac{E(T_+ - T_-)^2 c}{RT_+^2 Qz} \exp(E/RT_+);$$
$$r_q \sim \left[\frac{\lambda}{\rho Qz} \frac{E(T_+ - T_-)^2}{RT_+^2} \exp(E/RT_+)\right]^{1/2};$$
$$r_{ch} \sim \left[\frac{\lambda}{\rho Qz} \frac{RT_+^2}{E} \exp(E/RT_+)\right]^{1/2}; \tag{1.7}$$
$$q_* \sim \left[\frac{E(T_+ - T_-)}{\lambda \rho Qz RT_+^2} \exp(E/RT_+)\right]^{-1/2}.$$

Estimate (1.7) agrees with the results obtained in [4, 7] to an accuracy of a constant factor.

It follows from a qualitative analysis of the problem that the difference $T_+ - T_-$ is the temperature scale in the exterior region of the solution and r_q is the distance scale. In the interior region the characteristic scales are RT_+^2/E and r_{ch}, respectively. The time scale for both regions is equal to t_0.

2. Asymptotic Analysis of the Solution. The initial statement of the problem in the exterior variables is written in the form

$$\begin{gathered}\partial u/\partial\tau = \partial^2 u/\partial x^2 - (1/\varepsilon)\exp[-(1/\varepsilon)u/(1-\sigma u)];\\ u(x, 0) = 1,\ u(0, \tau) = 0,\ u(\infty, \tau) = 1,\\ u = (T_+ - T)/(T_+ - T_-),\ x = r/r_q,\ \tau = t/t_0,\\ 0 < \sigma = 1 - T_-/T_+ < 1.\end{gathered} \tag{2.1}$$

In the interior variables we have

$$\varepsilon^2 \partial U/\partial\tau = \partial^2 U/\partial X^2 + \exp[U/(1+\beta U)]; \tag{2.2}$$
$$U(X, 0) = -1/\varepsilon,\ U(0, \tau) = 0,\ U(\infty, \tau) = -1/\varepsilon, \tag{2.3}$$
$$U = E(T - T_+)/RT_+^2,\ X = r/r_{ch},\ \tau = t/t_0,\ \beta = RT_+/E \ll 1.$$

Since the small parameter ε in Eq. (2.1) has entered into the exponent, it is possible to neglect nonlinearity in the exterior region to any order of accuracy with respect to ε. Therefore, the solution of the exterior problem which satisfies the conditions $u(x, 0) = 1$, and $u(\infty, \tau) = 1$ is of the form

$$u = A + (1-A)\Phi(x/2\sqrt{\tau}),\ \Phi(z) = \frac{2}{\sqrt{\pi}}\int_0^z \exp(-y^2)\,dy. \tag{2.4}$$

Here the parameter A is determined by splicing.

The solution of Eq. (2.3) in the interior region is sought in the form of a series

$$U(X, \tau, \varepsilon) = \mu_1(\varepsilon)U_1(X, \tau) + \mu_2(\varepsilon)U_2(X, \tau) + \ldots, \tag{2.5}$$

where $\mu_n(\varepsilon)$ is an asymptotic sequence, $\mu_{n+1}/\mu_n \to 0$, $\varepsilon \to 0$. Substituting (2.5) into (2.2) and (2.3), we obtain $\mu_1(\varepsilon) = 1$ and $\mu_2(\varepsilon) = \varepsilon^2$, and the interior problem reduces to the equation

$$\partial^2 U_1/\partial X^2 + \exp U_1 = 0. \tag{2.6}$$

with an accuracy out to $O(\varepsilon^2)$ when $\beta \ll 1$.

All further discussion is limited by the error $O(\varepsilon^2)$; therefore, the equations for U_2, U_3,... are not written out. The solution of (2.6) is of the form

$$U_1(X, \tau) = \ln 2a \pm \sqrt{2}aX + 2b - 2\ln[1 + \exp(\pm\sqrt{2a}X + 2b)],$$

where a and b are functions of the time τ. The nearer boundary condition gives the value $a = \cosh^2 b$. It is sufficient to know the temperature field distribution at the instant τ_0 in order to reply to a particular question. Therefore, instead of $a(\tau)$ and $b(\tau)$, which are determined by splicing in the general case and do not contradict the divergence of the initial conditions, we will take their values at the instant τ_0.

The solution of (2.6) which satisfies the nearer boundary condition $U_1(0, \tau) = 0$ and the thermal equilibrium condition $\partial U_1(0, \tau_0)/\partial X = 0$, is of the form

$$U_1 = 2\ln 2 - \sqrt{2}X = 2\ln\ [1 + \exp\ (-\sqrt{2}X)]. \tag{2.7}$$

The time to establish thermal equilibrium τ_0 is sought in the form of a series

$$\tau_0 = D_0Q_0(\varepsilon) + D_1Q_1(\varepsilon) + D_2Q_2(\varepsilon) + \ldots, \tag{2.8}$$

where $Q_n(\varepsilon)$ is an asymptotic sequence, $Q_{n+1}/Q_n \to 0$, and $\varepsilon \to 0$. The coefficients D_n and the asymptotic sequence are found as a result of splicing the interior (2.7) and exterior (2.4) expansions. The rule of boundary splicing consists of the requirement of identical asymptotic behavior of the interior and exterior expansions written in the very same variables, i.e.,

$$\lim_{\varepsilon\to 0} \frac{u - \varepsilon|U_1|}{\gamma(\varepsilon)} = 0, \tag{2.9}$$

where as $\varepsilon \to 0$ the order of magnitude of γ is chosen to be equal to the order of magnitude of the spliced term.

After substitution of (2.8) into (2.4) and the expansion of $\Phi(z)$ into a series in the vicinity of $z = 0$, we arrive for small x at the expansion

$$u = A + \frac{2(1-A)}{\sqrt{\pi}}\left\{\frac{1}{2\sqrt{D_0Q_0}}\left(1 - \frac{x^2}{12D_0Q_0} + \frac{x^4}{160D_0^2Q_0^2} + \ldots\right)x - \frac{D_1}{4D_0^{3/2}Q_0^{3/2}}\left(1 - \frac{x^2}{4D_0Q_0} + \frac{x^4}{32D_0^2Q_0^2} + \ldots\right)x\ Q_1 - \right.$$
$$\left. - \frac{D_2}{4D_0^{3/2}Q_0^{3/2}}\left(1 - \frac{x^2}{4D_0Q_0} + \frac{x^4}{32D_0^2Q_0^2} + \ldots\right)xQ_2 + \frac{3D_1^2}{160D_0^{5/2}Q_0^{5/2}}\left(1 - \frac{5x^2}{12D_0Q_0} + \frac{7x^4}{96D_0^2Q_0^2} + \ldots\right)xQ_1^2 + \ldots\right. \tag{2.10}$$

The first term of the expansion (2.8) is determined by the usual means of splicing: the substitution of (2.10) and (2.7) into (2.9) in the case $\gamma(\varepsilon) = 1$ gives

$$A = 0,\ D_0 = 1/2\pi,\ Q_0(\varepsilon) = 1.$$

The standard splicing procedure does not take place for the determination of the succeeding terms of (2.8), since it is necessary here to know the rate at which the exterior variable x tends to zero [with respect to the appropriate terms of the asymptotic sequence $Q_n(\varepsilon)$]. It has not proven possible in the present approach to find an accurate value of the rate at which x tends to zero; however, it appears possible to obtain useful estimates.

The splicing condition is satisfied in two cases:

$$x^2/Q_1(\varepsilon) \to 0 \tag{A}$$
$$Q_1(\varepsilon)/x^2 \to 1, \tag{B}$$
$$\varepsilon \to 0.$$

Since τ_0 is determined by an asymptotic series constructed on the sequence $Q_n(\varepsilon)$, then when $Q_1(\varepsilon) \gg x^2$, we will have in case A the upper, and when $Q_1(\varepsilon) \sim x^2$, in case B the lower limit to the time τ_0.

The case $Q_1(\varepsilon)/x^2 \to 0$ is not possible, since it is impossible to satisfy the condition (2.9).

We will derive formulas for the upper limit to the time τ_0. Substituting (2.7) and (2.10) into (2.9) and setting $\gamma(\varepsilon) = \varepsilon$ with satisfaction of (A) taken into account, we will have

$$xQ_1(\varepsilon) = \varepsilon,\ D_1 = \sqrt{2}\ln 2/\pi,\ Q_1(\varepsilon) \gg \varepsilon^{2/3}. \tag{2.11}$$

Upon comparison of the subsequent terms of the expansion, three cases are already possible:

$$x^2/Q_1^2(\varepsilon) \to 0; \tag{A$_a$}$$
$$Q_1^2(\varepsilon)/x^2 \to 1; \tag{A$_b$}$$
$$Q_1(\varepsilon)/x^2 \to 0, \tag{A$_c$}$$

which are arranged in the order of decrease of the respective limits of the series (2.8).

In case (A_a), $Q_1(\varepsilon) \gg \varepsilon^{1/2}$ in agreement with (2.11), and we have from (2.9)

$$Q_2(\varepsilon) = Q_1^2(\varepsilon),\quad D_2 = 3\ln^2 2/\pi\ ;$$

consequently, one of the upper limits is written in the form

$$\tau_0 = 1/2\pi + (\sqrt{2}\ln 2/\pi)\,Q_1(\varepsilon) + (3\ln^2 2/\pi)\,Q_1^2(\varepsilon) + \ldots,$$

where $\varepsilon^{1/2} \ll Q_1(\varepsilon) \ll 1$.

In case (A_b), $Q_1(\varepsilon) = \varepsilon^{1/2}$ and upon continuation of splicing we find the coefficients

$$Q_2(\varepsilon) = \varepsilon,\ D_2 = 3\ln^2 2/\pi - 1/6$$

and the corresponding expansion

$$\tau_0^+ = 1/2\pi + (\sqrt{2}\ln 2/\pi)\,\varepsilon^{1/2} + (3\ln^2 2/\pi - 1/6)\,\varepsilon + \ldots. \tag{2.12}$$

Similarly, in case (A_c) we will have

$$Q_2(\varepsilon) = Q_1^{-2}(\varepsilon)\,\varepsilon^2,\ D_2 = -1/6,\ \varepsilon^{2/3} \ll Q_1(\varepsilon) \ll \varepsilon^{1/2};$$

$$\tau_0 = 1/2\pi + (\sqrt{2}\ln 2/\pi)Q_1(\varepsilon) - (1/6)Q_2(\varepsilon) + \ldots.$$

The lower limit of τ_0^- is found by successive splicing of (2.9) with the condition that case (B) is fulfilled. Leaving out the intermediate steps, we write out the final answer:

$$\tau_0^- = 1/2\pi + (\sqrt{2}\ln 2/\pi - 1/6)\varepsilon^{2/3} + (3\ln^2 2/\pi - \pi/60)\,\varepsilon^{4/3} + \ldots\ . \tag{2.13}$$

TABLE 1

θ_0	3	6	10	25	50	100	300
τ_* (Computer)	3,7	11	25,6	128,8	464	1744	14960
τ_*^+ (2.12)	3,9	12,1	28,7	145,8	523	1933	16033
τ_*^- (2.13)	2,9	8,7	20,9	113,6	430	1668	14634
$\langle\tau_*\rangle$ (3.1)	3,4	10,4	24,8	129,7	476	1800	15333
τ_{*2} (3.3)	3,1	10,4	27,1	161,1	628	2519	22550
τ_{*1} (3.2)	2,9	11,5	31,8	198,9	796	3183	28648
τ_{*7} (3.4)	1,4	5,7	15,9	99,4	398	1592	14324
τ_*^+, %	6,2	9,7	12,1	13,2	12,7	10,8	7,17
τ_*^-, %	−21,4	−21,3	−18,2	−11,8	−7,3	−4,4	−2,2
$\langle\tau_*\rangle$, %	−7,6	−5,8	−3,0	0,6	2,7	3,2	2,5
τ_{*2}, %	−14,9	−5,1	5,7	25,1	35,4	44,4	50,6
τ_{*1}, %	−22,6	4,2	24,34	54,5	71,5	82,5	91,50
τ_{*7}, %	−61,3	−47,9	−37,8	−22,8	−14,3	−8,8	−4,3

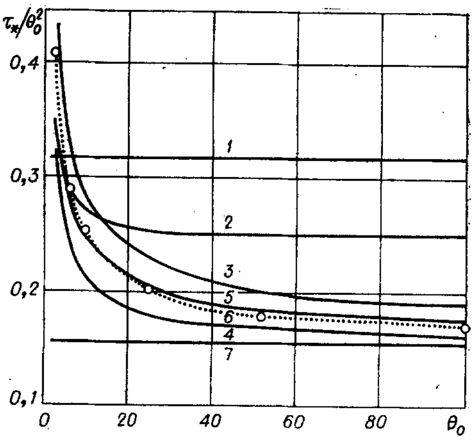

Fig. 1

3. Discussion of the Results. The numerical solution of the problem (2.2) and (2.3) was accomplished on an M-220 computer on the basis of an implicit difference scheme with selection of the optimal Courant number and extrapolation of the calulated values to zero step [8]. Comparison of Eqs. (2.12), (2.13), and the numerical calculation of [9] (see the table in which the conversion of τ_0 to $\tau_* = \theta_0{}^2\tau_0$ is done) shows that satisfactory agreement (with an error no worse than 3%) is given over a wide interval of variation of the temperature head ($10 \leqslant \theta_0 \leqslant 300$) by the arithmetic mean formed from the upper and lower limits:

$$\langle\tau_*\rangle = (\tau_*^+ + \tau_*^-)/2. \tag{3.1}$$

A comparison of the results obtained by different authors with the numerical calculation and the asymptotic Eqs. (2.12), (2.13), and (3.1) is given in Fig. 1 and Table 1:

curve 1 corresponds to Enig's formula [5],

$$\tau_{*1}/\theta_0^2 = 1/\pi; \tag{3.2}$$

curve 2 corresponds to Grishin's formula [6],

$$\frac{\tau_{*2}}{\theta_0^2} = -\frac{\theta_0}{6}\ln\left(1 - \frac{3}{2\theta_0}\right) \approx \frac{1}{4} + \frac{3}{16}\frac{1}{\theta_0} + \frac{3}{16}\frac{1}{\theta_0^2} + \frac{27}{128}\frac{1}{\theta_0^3} + \ldots; \tag{3.3}$$

curve 7 corresponds to Zel'dovich's formula [4],

$$\tau_{*7}/\theta_0^2 = 1/2\pi; \tag{3.4}$$

curves 3-5 correspond to Eqs. (2.12), (2.13), and (3.1), respectively; the results of the computer calculation are shown by open circles (curve 6). It is obvious that Eqs. (3.2) and (3.3) do not give the correct asymptote, in contrast to (3.4).

The authors express their gratitude to O. B. Sidonskii and G. V. Aleinikova for supplying some results of the numerical calculation.

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USE OF ELECTRIC EXPLOSION OF WIRES IN A HIGH-PRESSURE GAS TO BREAK A CURRENT CIRCUIT

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UDC 537.527.5;533.9.07

The high energy densities stored in the magnetic field of inductive storage devices have promising applications in experimental physics. The greatest energy storage levels are achieved in superconducting storage facilities and pulsed facilities, operating with explosive-magnetic generators (currents up to $3 \cdot 10^8$ A) [1].

To use the energy stored in a magnetic field one must cut the current in the storage circuit and switch it to the load circuit. One method of doing this is to use a switch based on electrical explosion of wires (EEW) [1]. There are several difficulties in creating current cut-off devices of this type: After the electric explosion a column of metal vapor forms in which breakdown can occur; then the cut-off process is slowed and the energy-transfer efficiency is decreased. The problem is that the wire material is instantly vaporized, i.e., it is a dielectric subject to stresses arising when the inductive storage device is switched to the load.

As the pressure of the surrounding medium is increased it becomes more difficult to create shunting arcs in EEW devices and to produce conditions for more complete vaporization of the wire material. A series of tests has been conducted with different materials in order to elucidate the possible use of EEW in a high-pressure gas for current switching. The equipment contained a high-pressure chamber with inserted electrodes, between which a wire of the test material was attached. The chamber was filled with argon at a pressure in the range 1-750 atm. The wire diameter in the various tests was 0.5-1 mm. A control switch was used to discharge a condenser bank into the wire, of capacity 200 μF, voltage 3-6 kV, circuit inductance ∿ 1 μH, and with length of the first half-period current 50-100 μsec. The current and the voltage were measured using a shunt and a voltage divider, from which the signals were recorded on a type S1-29 oscilloscope. The results obtained are shown in Table 1, where l is the wire length; p is the inert gas pressure in the working chamber; I_{max} is the maximum current before explosion of the wire; I_{max}' is the maximum discharge current arising after EEW; $U_{c.b}$ is the voltage on the capacitor bank; and τ is the closure time, i.e., the time for the current to fall from its maximum value to zero. It can be seen from Table 1 that for a certain pressure (different value for the different metals) a closure effect occurs, i.e., after the current increases the EEW occurs and the electric circuit is broken. In this case part of the stored power remains in the capacitor bank. Of the metals tested, the most suitable for current interruption are Li and Al. At a pressure of more than 300 atm, wires of these

Khar'kov. Translated from Prikladnoi Mekhaniki i Tekhnicheskoi Fiziki, No. 6, pp. 102-105, November-December, 1976. Original article submitted March 29, 1976.