Model of stationary laser plasma corona with allowance for the effect of deflagration at the critical surface

E. N. Ragozin and P. V. Sasorov

P. N. Lebedev Institute of Physics, Academy of Sciences of the USSR
(Submitted 16 May 1980)
Zh. Eksp. Teor. Fiz. **80**, 1371–1382 (April 1981)

A model is constructed which admits of the existence of a rarefaction jump (deflagration) that is not connected with the effect of the light pressure. There occurs as a result of the assumption by the electronic thermal conductivity of anomalously low values in the vicinity of the critical point a hydrodynamic discontinuity (a jump in the density, temperature, etc.) accompanied by nonequilibrium ionization of the ions. Another manifestation of the deflagration effect is the production of hard x rays and high-energy ions in fluxes $q_L > 10^{13}$–10^{14} W/ cm^2. The theoretical calculations are compared with experiment within the framework of the proposed model.

PACS numbers: 42.55.Bi

I. INTRODUCTION

Theoretical and experimental investigations aimed at obtaining lasing in the far ultraviolet region of the spectrum on transitions of multiply charged ions in a laser plasma (LP) have been carried out in recent years.[1,2] Any well-founded assessment of the prospects of a particular approach and the optimization of the experiment turn out to be impossible if we do not have answers to a number of questions of a diagnostic nature. Among such key questions are the following: 1) What are the profiles of the electron density (N_e) and the electron temperature (T_e) in the LP corona? 2) What are the values of the ionization temperature (T_Z), and do the mean values of the charge Z in different regions of the laser flare correspond to the equilibrium charge values at given T_e? Therefore, the construction of an adequate model of the laser corona is a pressing problem. Moreover, the investigation of the LP is of undoubted spectroscopic and general physical interest. As will be seen below, the large-scale characteristics of the corona are determined, in particular, by the plasma-physical processes that occur in the vicinity of the point with the critical density

$$\rho^{cr}=N_e^{cr}\frac{Am_p}{Z}=\pi\frac{m_e m_p c^2 A}{\lambda_L^2 e^2 Z^{cr}},$$

where $m_{e,p}$ are the electron and proton masses, c is the velocity of light, λ_L is the wavelength of the laser radiation, and A is the mass number of the target nuclei. Therefore, the concretization of the model of the laser corona provides the prerequisites for the elucidation of the physical processes occurring in the zone of (collective) absorption of the laser radiation, which is of general physical interest.

The necessity to return to this problem after the appearance of Ref. 3, in which a stationary model of the corona of spherical laser targets is constructed, is due to the following circumstances. A number of experiments on the LP (the experimental situation is analyzed in Sec. III of the present paper) indicate that the plasma density in the region $l > l^{cr}$ (l is the distance from the target surface) is significantly lower than the value predicted by the theory.[3] This compelled one of the authors of the present paper to postulate[4] that a rarefraction jump exists in the neighborhood of the critical point.[1)] This is at variance with Ref. 3. The theoretical analysis, carried out below, of the situation shows that it is impossible to obtain a unique solution to the formulated problem within the framework of hydrodynamics without going deep into the physical processes that occur in the neighborhood of l^{cr}. There exists a three-parameter family of solutions to the hydrodynamic equations, to choose among which is possible only when we go over to a microscopic description in the neighborhood of the critical point. The solutions discussed in Ref. 3 belong to this family, and are the only continuous solutions.

In the present paper two of the three parameters are chosen on the basis of reasonable physical assumptions with allowance for the indications of experiment. The third parameter is selected by quantitatively fitting the results of the model to the experimental data.

II. A THEORETICAL MODEL OF THE CORONA

In the range of laser fluxes from 10^{11} to 10^{16} W/cm^2, the laser corona parameters for geometric target dimensions ~10^{-2} cm are such (see below) that the inverse bremsstrahlung can be neglected in the $\rho<\rho^{cr}$ region. In this case the main absorption occurs in the neighborhood of $\rho\sim\rho^{cr}$ because of the conversion of the electromagnetic wave into a plasma wave and the fairly high effective electron collision rate. The plasma wave dissipates, as result of the Landau damping, as heat, possibly via the strong-Langmuir-turbulence phase.

Our aim is to construct a large-scale model of the laser corona. For this reason, we shall construct it within the framework of hydrodynamics. This excludes from consideration the phenomena that occur on a scale smaller than the mean free path of the electrons. The small parameter that makes such a description possible is, as will become clear below, the ratio $m_e/M_i \ll 1$. As is well known, the equations of hydrodynamics with allowance for viscosity and thermal conduction do not allow the existence of discontinuous flows.[5] The problem under consideration, however, has the following specific features. The dispersion of the plasma in the supercritical ($l > l^{cr}$) region occurs at supersonic

 0038-5646/81/040701-07$02.40

velocities. The characteristic scale of the variation of the flow parameters (the distance from the critical surface to the target, l^{cr}) is determined by the thermal conductivity. Therefore, the Péclet number is of the order of unity ($Pe = N_e V l^{cr}/\varkappa \sim N_e c_s l^{cr}/\varkappa \sim 1$). Here V is the characteristic flow velocity; c_s is the isothermal velocity of sound; $c_s^2 = p/\rho$; p being the plasma pressure $\varkappa$ is the thermal conductivity coefficient: $\varkappa \sim N_e V_{Te} \lambda_{ei}$, V_{Te} and λ_{ei} being respectively the thermal velocity and mean free path of the electrons. On such scales the Reynolds number is much greater than unity, which indicates that the viscous forces can be neglected in the hydrodynamic equations ($\mathrm{Re} = \rho V l^{cr}/\eta \gg 1$; $\eta \sim \rho V_{Ti} \lambda_{ii}$ is the coefficient of viscosity, V_{Ti} is the thermal velocity of the ions and λ_{ii} is the mean free path of the ions with respect to the ion–ion collisions). On that ground, viscosity at such dimensions introduces only small corrections, and can (with a definite accuracy that rises with increasing Z) be neglected. Viscosity must be taken into consideration in the problem in question if the plasma parameters possess large gradients, so that the gradient lengths are comparable to the mean free path of the ions. The investigation of such effects are excluded from our treatment, since it requires a kinetic description.

Thus, we shall consider the model of the laser corona within the framework of hydrodynamics without allowance for viscosity. It is well known (see, for example, Ref. 6) that, without allowance for viscosity, the equations of hydrodynamics admit of the existence of discontinuities. Moreover, a Riemann wave of sufficiently high amplitude, for example, gets "broken up" in this case with the formation of a discontinuity. The width of the discontinuity is then determined by the viscosity. Thus, we cannot exclude the possibility of the existence of discontinuities in our solution. We shall require the conservation of matter, momentum, and energy at the possible discontinuities.

We shall consider the problem of laser-target evaporation in spherical geometry. Let Q_L be the energy flux of the laser radiation incident on a sphere of radius R from a unit solid angle, so that the intensity of illumination of the sphere is equal to $q_L = Q_L/R^2$. Let us further assume that a significant fraction, q, of this radiation is absorbed at the $r = r^{cr}$ surface, where $\rho = \rho^{cr}$. Let ρ^{cr} be much lower than the sphere's density, which we set equal to infinity. The characteristic distance over which the laser-radiation absorption occurs is determined, in particular, by the Debye distance, which, under our conditions, is much shorter than l^{cr}; therefore, we shall assume that the heat is released in a delta-function fashion. The heat that enters the sphere of radius $r = r^{cr}$ is wholly expended on the "evaporation" of the target material. The following conditions are fulfilled at the target surface:

$$\rho(R)=\infty,\ T_e(R)=0,\ \varkappa T_e'(R)=0,\ V(R)=0, \tag{1}$$

where V is the radial velocity of the plasma and $\varkappa$ is the electronic thermal conductivity coefficient. The plasma disperses into a vacuum, so that the following conditions must be imposed at large r:

$$\rho(\infty)=0,\ T_e(\infty)=0,\ r^2\varkappa T_e'|_{r\to\infty}=0. \tag{2}$$

We neglect the following effects:

a) the laser inverse bremsstrahlung in the region $\rho < \rho^{cr}$,

b) the collective absorption effects in the $\rho < \rho^{cr}$ region,

c) heat transfer by epithermal electrons,[7]

d) the radiative energy losses,

e) the effects of the ionization of the material, i.e., we shall assume that $Z = \mathrm{const}$ (see, however, below).

Then the stationary flow of the plasma is described by the following equations:

$$\rho V r^2 = \mathrm{const} = \dot{m}, \tag{3}$$

$$\rho V V' = -(\rho c^2)', \tag{4}$$

$$\frac{\rho V r^2}{2}(5c^2+V^2) - \tilde{\varkappa} r^2 c c' = \varphi(r), \tag{5}$$

where

$$\tilde{\varkappa} = \frac{2\varkappa A m_p}{Z},\quad c^2 = \frac{Z T_e}{A m_p}\quad \varphi = \begin{cases} 0, & r < r^{cr} \\ q, & r > r^{cr} \end{cases}.$$

The laws of conservation of matter, momentum, and energy are fulfilled at the possible discontinuities. Since the energy transport equation (5) and the equation of continuity (3) have been written in the integral form, there is no need to write out again the corresponding conditions at the discontinuities. Let us just give the condition for the conservation of the momentum flux across the discontinuity:

$$\rho(V^2+c^2)\,|_2^1 = 0. \tag{6}$$

The indices 1 and 2 denote the points respectively before and after the discontinuity (Fig. 1). The plasma flows from left to right.

Since the flow is subsonic in the region of small r and supersonic at large r, there exists a point where the flow velocity passes through the isothermal sound velocity value. A smooth passage through the isothermal sound velocity value in the $r < r^{cr}$ region is impossible, since otherwise the condition $q_T < q_{max}$, where q_T is the heat energy flux and

$$q_{max} = -5\Phi_e \rho c^3 \tag{7}$$

is the maximum thermal-conduction-induced energy flux[8] directed into the dense layers of the plasma in the $r < r^{cr}$ region, and realized upon the violation of the condition $\lambda_{ei} \ll L_T = T_e/|\nabla T_e|$, will not be fulfilled. The point is that the classical expression $\mathbf{q}_T = -\varkappa_{Sp}\nabla T$ ($\varkappa_{Sp}$

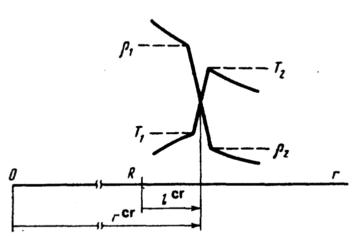

FIG. 1. Hydrodynamic "discontinuity" in the neighborhood of the critical point. The behavior of the quantities T_e and ρ at the discontinuity. The indices 1 and 2 denote the values of the quantities respectively to the left and right of the discontinuity.

is the Spitzer value of the thermal conductivity coefficient) is applicable under conditions when the Coulomb mean free path of the electrons is short compared to the temperature gradient length, i.e., when $\lambda_{ei} \ll L_T$. At the discontinuity λ_{ei} is not small in the indicated sense, and therefore the use of the expression (7), which describes the saturation of the heat flux, is justified. In the simple situation, $\Phi_s \sim 1$. If, however, the effective mean free path of the electron is made shorter than L_T (e.g., as a result of the development of strong magnetic and electric fields), the factor Φ_s is correspondingly reduced. There are experimental indications[4,9-11] that the electronic thermal conduction is strongly suppressed (i.e., $\Phi_s \ll 1$) in the vicinity of the critical point in a LP.

A simple analysis shows that the solution can admit of only a rarefaction jump, otherwise we cannot satisfy the boundary conditions without violating the condition $q_T < q_{max}$. Such a discontinuity is possible only in the region where energy is released, and is called deflagration (see, for example, Ref. 6).

Let us make the additional assumption that the radius of the target is sufficiently large, so that $l^{cr} \ll R$. This assumption does not affect the physics of the matter, and we can get rid of it, though that will lead to somewhat unwieldy formulas. A full analysis of the experiments in which we encounter the opposite situations ($l^{cr} \lesssim R$) is carried out without allowance for this assumption.

A simple count of the number of unknowns in the system (1)–(6) shows that the solution to the system is characterized by three arbitrary parameters. As these parameters, it is convenient to take:

a) $M_1 = V_1/c_1 (M_1 < 1)$;
b) $M_2 = V_2/c_2 \geqslant 1$;
c) $\rho_1/\rho_2 (\rho_1 > \rho^{cr} > \rho_2)$.

Let us recall that the indices 1 and 2 denote the values of the quantities respectively before and after the discontinuity.

It may be inferred that the existence of the discontinuity is a consequence of the restriction imposed on the thermal flux so that the maximum thermal flux $5\Phi_s \rho c^3$ will be realized on the left of the discontinuity. Because of this, M_1 and Φ_s become connected by the relation $M_1(1 + M_1^2/5) = 2\Phi_s$, which expresses the equality of the heat and hydrodynamic energy fluxes under steady-state conditions.

Since we have used the equations of hydrodynamics throughout, the parameters M_1, M_2, and ρ_1/ρ_2 cannot be determined within the framework of hydrodynamics, and require for their determination a kinetic description of the phenomena occurring in the vicinity of the critical surface, in particular, a self-consistent kinetic description of the process of collective laser-radiation absorption.

We shall choose these parameters in the following manner. We shall determine the value of $M_1(\Phi_s)$ by quantitatively fitting the model to experiment. Further, we shall assume that the Chapman-Jouguet rule

$$M_2 = 1, \qquad (8)$$

which, as a rule, arises in the microscopic consideration of detonation waves,[6] is satisfied. Indications that (8) is valid are provided by experiments on LP diagnostics with the aid of the line emissions of multiply charged ions in the far infrared region of the spectrum.[4,12,13] We fix the third parameter by assuming that

$$(\rho_1 \rho_2)^{1/2} = \rho^{cr}. \qquad (9)$$

The condition (9) arises if we assume that the vital—to the hydrodynamics—rapid heat release accompanying the laser-radiation absorption occurs in the region where $\omega_L/k < \omega_p = (4\pi N_e e^2/m)^{1/2} < k\omega_L (k > 1)$. We realize that the conditions (8) and (9) need to be refined and justified on the basis of a self-consistent microscopic theory. An additional (and, perhaps, prime) justification of these hypotheses will be the agreement of all the available experimental data (see below) that are sensitive to these two assumptions.

We shall, in accordance with the condition $l^{cr} \ll R$, assume that the region in which the thermal conduction is appreciable is significantly smaller than the target region. For $r > r^{cr}$ this implies that

$$\varkappa_2 T_2 r^{cr}/Q_L \ll 1. \qquad (10)$$

The inequality (10) is fulfilled at not only moderate values of q_L, since as $\varkappa_2$ we should use not the Spitzer value of the thermal conductivity coefficient ($\varkappa_{Sp} = \varkappa_0 T_e^{5/2}$), but a significantly smaller value because of the strong suppression of the thermal conduction that obtains in the $r > r^{cr}$ region as well, owing, for example, to the presence of the strong magnetic fields produced in the corona,[14] ion-acoustic vibrations, etc. In this case the parameter (10) drops out of the results. Correcting small errors in the solution given in Ref. 3 for the $r > r^{cr} = r_{1,2}$ region, we have at distances where thermal conduction is no longer important the formulas

$$\frac{T}{T_2} = \frac{1}{5}\left[\frac{25}{4} - \left(\frac{V}{c_2}\right)^2\right], \quad \left(\frac{V}{c_2}\right)^2 \left[\frac{25}{4} - \left(\frac{V}{c_2}\right)^2\right]^3 = \frac{3^3 \cdot 5^8}{2^{16}} \left(\frac{r_2}{r}\right)^4,$$
$$\frac{\rho}{\rho_2} = \frac{c_2}{V}\left(\frac{r_2}{r}\right)^2, \qquad (11)$$

which at large r yield

$$\frac{T}{T_2} = \frac{15 \cdot 2^{1/3}}{2^5}\left(\frac{r_2}{r}\right)^{4/3} \quad \left(\frac{V}{c_2}\right)^2 = \frac{25}{4} - \frac{3 \cdot 25 \cdot 2^{1/3}}{2^5}\left(\frac{r_2}{r}\right)^{4/3}$$
$$\frac{\rho}{\rho_2} = \frac{2}{5}\left(\frac{r_2}{r}\right)^2 \left[1 + \frac{3 \cdot 2^{1/3}}{2^4}\left(\frac{r_2}{r}\right)^{4/3}\right]. \qquad (12)$$

In the formulas (11) we have not distinguished between the locations of the isothermal and adiabatic Jouguet points in view of their spatial proximity. Moreover, we have the relation

$$q = \frac{25}{8}\rho_2 c_2^3. \qquad (13)$$

The solution for $r < r^{cr}$ is much the same as the solution given in Ref. 3. The only difference is that we use another boundary condition at r^{cr}, where we have the relation:

$$\dot{m}=\rho_1 V_1 r_1^2=\frac{4}{25}\frac{Am_p\varkappa_1 r^{\mathrm{cr}2}}{Z_1 l^{\mathrm{cr}}}. \tag{14}$$

Solving (3), (6), (8), (9), (13), and (14) simultaneously, we obtain

$$c_1=2^{5/6}(8/25)^{1/3}(q/\rho^{\mathrm{cr}})^{1/3}M_1^{2/3}/(1+M_1^2)^{5/6},$$
$$c_2=2^{-1/6}(8/25)^{1/3}(q/\rho^{\mathrm{cr}})^{1/3}(1+M_1^2)^{1/6}/M_1^{1/3},$$
$$\rho_1=\rho^{\mathrm{cr}}(1+M_1^2)^{1/2}/2^{1/2}M_1,\quad \rho_2=\rho^{\mathrm{cr}}\cdot 2^{1/2}M_1/(1+M_1^2)^{1/2}, \tag{15}$$
$$l^{\mathrm{cr}}=\left(\frac{4}{5}\right)^5\frac{5^{1/3}}{2^{1/6}}\frac{\varkappa_0}{\rho^{\mathrm{cr}}}\left(\frac{Am_p}{Z}\right)^{7/3}\left(\frac{q}{\rho^{\mathrm{cr}}}\right)^{4/3}\frac{M_1^{8/3}}{(1+M_1^2)^{23/6}}.$$

Thus, we have obtained all the characteristic quantities, (15) describing the plasma flow as functions of a single unknown quantity, M_1. The flow in the $r>r^{\mathrm{cr}}$ region is described by the formulas (11) and (12). Let us write out here for reference the expression for N_e at $r\gg r^{\mathrm{cr}}$:

$$N_e=\frac{2}{5}N_e^{\mathrm{cr}}\left(\frac{r^{\mathrm{cr}}}{r}\right)^2\left[1+\frac{3\cdot 2^{1/3}}{2^4}\left(\frac{r^{\mathrm{cr}}}{r}\right)^{4/3}\right]\frac{2^{1/2}M_1}{(1+M_1^2)^{1/2}}. \tag{16}$$

III. EXPERIMENTAL FITTING OF THE MODEL

1. Our model is described by a single free parameter, M_1 (or Φ_s). The most sensitive to changes in M_1 and, perhaps, the most easily measured is the value of N_2 (the electron density immediately after the rarefaction jump). The N_e values in the region $r>r^{\mathrm{cr}}$ satisfy the relation (11) (with, evidently, ρ replaced by N_e). It can be seen from (11) that the characteristic scale of the variation of N_e at $r>r^{\mathrm{cr}}$ (i.e., outside the region where the thermal conduction is important) is equal to r^{cr}. From this it follows that if the rarefaction jump itself turns out to be unobservable in an experiment, the N_e values at distances from the critical surface much shorter than R are close to the values of $N_e=Z\rho_2/Am_p$.

In Table I we have collected the results of N_e measurements performed at distances of 10–15 μm (Refs. 12–20), which results can, in our opinion, be considered to be N_2 values. Only Attwood *et al.*[19] have directly observed the density jump. In all the above-cited investigations a neodymium laser with an emission wavelength of $\lambda_L=1.06$ μm ($N_e^{\mathrm{cr}}=10^{21}$ cm^{-3}) was used to excite the LP. The most reliable, in our opinion, are the recent N_e measurements performed by Peregudov *et al.*,[12] using the Stark broadening of the transitions of the Balmer series and the interferometric measurement performed by Raven *et al.*[14] These two measurements yielded the same value for $N_2\approx 0.16N^{\mathrm{cr}}$, although the corresponding q_L values differ by 2–3 orders of magnitude.

It can be seen from the values given in Table I that N_2 is appreciably smaller than N^{cr}, and that there is no apparent dependence of N_2 on q_L, even though the latter quantity varies through 5 orders of magnitude. To the value $N_2=0.16N^{\mathrm{cr}}$ corresponds $M_1\approx 0.11$ ($\Phi_s=0.055$). We arrive at the conclusion that the most realistic value of M_1 is 0.11, and that M_1 does not depend on q_L.[2)] For such values of M_1 we can justifiably replace $1+M_1^2$ by 1 in the formulas (15). It is possible that M_1 decreases somewhat with increasing A.

With the choice of the value of M_1, we complete the

TABLE I.

q_L, W/cm^2	Target	N_e, 10^{20} cm^{-3}	Measurement Technique;	Reference
10^{15}–10^{16}	–	<2.5	Dependence of reflection coefficient on the angle of incidence of radiation on target	[15]
10^{12}	Carbon	1,9	Stark broadening of the H_α and H_β transitions of the CVI ion (far UV region)	[16]
10^{11}–10^{13}	Calcium Titanium	0.3	Spectrum of oxygen-like ions (far UV region)	[17]
~10^{16}	Glass	1	Interferometry $\lambda=5300$ Å	[18]
10^{14} — 10^{15}	Glass	2–3	Interferometry $\lambda=2660$ Å	[19]
10^{11} — 10^{13}	Calcium Titanium Iron	0,3–0,6	Spectrum of oxygen-like ions (far UV region)	[20]
10^{13}	Nickel Calcium Nickel	1	$N_e(l)$ profile constructed with the aid of the equation of continuity (from the known $V(l)$ dependence) (far UV region)	[13]
10^{15}–10^{16}	Aluminum	1.6	Interferometry $\lambda=6330$ Å	[14]
10^{13}	Carbon	1.6	Stark broadening of Balmer-series transitions of the CVI ion (far UV region)	[12]

specification of our model. It is of interest to clarify how this model describes the rest of the experimental data, i.e., the data other than the density measurements, which have been used to concretize the model. An important characteristic parameter of the model is T_2.

2. The values of T_2 during the action of the laser radiation are significantly higher than the corresponding temperature values in the model proposed in Ref. 3 [by roughly a factor of $(\rho^{\mathrm{cr}}/\rho_2)\approx 3.4$]. At the same time, the LP volume at a temperature $T\sim T_2$ is relatively small (Fig. 2): the ratio $T(r)/T_2$ varies from 1 down to 0.65 over a thickness $\Delta l\approx 0.1R$. A typical value of R in an experiment is ~10^{-2} cm; therefore, to obtain the true spatial distribution of the temperature in a LP in the vicinity of the Jouguet point, we need an apparatus with a spatial resolution of ~10^{-3}–10^{-4} cm. The majority of the published measurements in which the temperatures in the LP were obtained from the x-ray continuum emission spectrum[9,18,22–24] were generally performed without spatial resolution. As a result of the spatial temporal integration that occurred in the course of the measurements, when x rays from LP regions with different T_e were integrated in each detector, the measured energy distribution of the quanta does not correspond to a single T_e value. The hardest x ray quanta are emitted by the near-target-surface region with a temperature ~T_2. The peripheral regions of the laser flare and also the LP that exists after the action of the laser pulse produce the long-wave part of the continuous x-ray spectrum.

The majority of the authors of Refs. 18, 22–24 un-

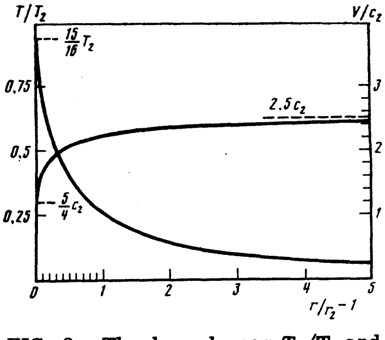

FIG. 2. The dependences T_e/T_2 and V/c_2 in the supercritical region.

conditionally relate the generation of hard x-ray quanta to the appearance of a relatively small number of "epithermal" electrons, characterized by a temperature T_h. Within the framework of the ideas developed in the present paper and in Ref. 4, the hard x-ray quanta, firstly, are a natural and inevitable consequence of the deflagration effect and, secondly, are generated precisely by the thermal electrons. For fluxes above ~10^{14} W/cm^2, T_2 attains, according to the formulas of the present paper, values ~10 keV. This should lead to the production of continuous-spectrum x rays with photon energies of up to ~100 keV. Estimates show that we can in this way explain not only the shape, but also the absolute intensity of the hard-x-ray spectra.[9]

Closely tied with the foregoing question is the question of the rate of dispersion and the kinetic energy of the ions. In the case of an "almost adiabatic" dispersion, the asymptotic kinetic energy, E_{as}, of the ions is connected with T_2 by the relation $E_{as} \approx 25/8\bar{Z}T_2 \approx 3\bar{Z}T_2$; consequently, it should, for fluxes $q_L \gtrsim 10^{13}$ W/cm^2, have a value of scores or hundreds of keV (depending on the A of the target). Thus, the irridation of calcium ($\bar{Z} \approx 13$) and nickel ($\bar{Z} \approx 20$) targets by nanosecond laser pulses in the case in which $q_L \sim 10^{13}$ W/cm^2 should lead to the production of ions with energies respectively equal to ~80 keV and ~120 keV, which is what is experimentally observed.[13] These ions are produced only during the action of the laser pulse, and their number, equal to ~$N_2 c_2 S_{foc} \tau_L / Z$, may be relatively low as compared to the total number of ions that leave the target, and are registered in the experiment.

The deflagration effect causes the kinetic energy of the dispersing ions to increase by a factor of $(\rho^{cr}/\rho_2)^{2/3}$, which allows us to account for the experimental data without recourse to nonhydrodynamic acceleration mechanisms. The term "fast ions" seems to us to be infelicitous, since it implies the action of some special acceleration mechanisms for the ions, and contrasts these ions with the thermal ions. The maximum ion energy in a LP has been observed by Siegrist *et al.*[25] in an irradiation of gold and tungsten targets, and was as high as 5.6 MeV for $q_L \sim (4-8)\times 10^{14}$ W/cm^2. The deflagration effect for these fluxes should lead to values of $T_2 \sim 25-40$ keV, while the high degree of ionization of the ions ($Z \sim 40-50$) ensures the acceleration of the ions to energies of 3–6 MeV.

Figure 3 shows the computed $T_2(q_L)$ dependences for $M_1 = 0.11$, different A/Z values, and a 50% and 100% laser-radiation power conversion into a hydrodynamic energy flux. For high q_L values the 50% conversion ratio is more realistic, in view of the increased reflection coefficient for the laser radiation.[15] We have plotted in Fig. 3: a) the T_h values obtained by Manes *et al.*[23] and Boiko *et al.*[24] from continuous x-ray spectrum measurements, b) the values of $E_{fi}/3\bar{Z}$ (where E_{fi} is the energy of the so-called "fast" ions) obtained by Decoste and Ripin[26] and McLean *et al.*[27] in time-of-flight measurements, c) the values of $E_{as}/3\bar{Z}$ (where E_{as} is the asymptotic kinetic energy of the ions) obtained by Peregudov and Ragozin[13] from measurements of the Doppler broadening of the spectral lines, and d)

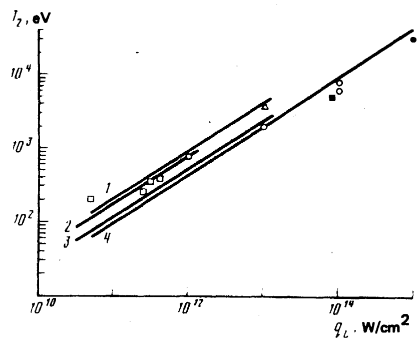

FIG. 3. Computed $T_2(q_L)$ dependences for $M_1 = 0.11$, $A/Z = 3$ (the straight lines 2 and 3), 4 (1), and 2 (4), and a 50% (the straight lines 3 and 4) and 100% (the lines 1 and 2) laser-radiation power conversion into hydrodynamic energy flux. ○) T_h values taken from Refs. 23 and 24; ■) and ●) $E_{fi}/3Z$ values from Refs. 26 and 27 respectively; △) $E_{as}/3Z$ values from Ref. 13; □) T_Z values.

the T_Z values experimentally obtained by us by a spectroscopic method, namely, from measurements of the relative intensities of the transitions, lying in the 100–200-Å region, of multiply charged calcium, titanium, and iron ions between the $2s^2 2p^k$ and $2s2p^{k+1}$ configurations.

Notice, in particular, the excellent agreement between the computed $T_2(q_L)$ values and the "epithermal"-electron temperature, T_h, obtained from continuous x-ray spectrum measurements in the region of high-energy cutoffs. It is noteworthy that Peregudov *et al.*[20] have established a limited increase of T_h with increasing mean size of the nuclear charge of the target. This fact is easily explained by the increase of T_2 if our observation that N_2 and, consequently, Φ_s have a tendency to decrease with increasing A is correct. The deviation of the temperature values from the computed curve for $q_L \lesssim 10^{11}$ W/cm^2 can be related to neglected (by us) processes that play an appreciable role in the general energy balance (e.g., to energy "losses" due to ionization and radiation emission), and should, according to our estimates, exert their influence at low q_L.

Thus, we see that our model satisfactorily accounts for the characteristic velocities (V_{as}), temperatures (T_2), and densities (ρ_2) of the flow. But a comparison of the spatial dependences $V(r)$, $T_e(r)$, and $\rho(r)$ with the accuracy required by us (i.e., to within ~1 μm) is at present impossible because of the inadequate spatial resolution in the corresponding experiments.

But some indirect verification of the dependences $N_e(r)$ and $T_e(r)$ is possible. It turns out that the density and temperature profiles obtained by us with the definite M_1 value of 0.11 allow us to explain the nonequilibrium character of the ionization in the LP corona.[12]

3. Let us, for definiteness, consider the experiment on the irradiation of a plane calcium target under conditions when $q \approx 2\times 10^{12}$ W/cm^2. We shall, allowing for the conical character of the LP dispersion in this case,[13] compare the dispersion with a spherical dispersion; we set R equal to the distance from the target surface to the point of convergence of the generatrices of

the dispersion cone, i.e., $R \approx r_{toc}/\tan\alpha$ ($r_{toc} \approx 200$ μm is the radius of the focal spot and α is the angle between the axis and a generatrix of the dispersion cone, and is $\approx 24°$; $R \approx 200$ μm). The most highly charged ions, abundantly present in the corona, are the Ca XVI ions, which corresponds to an ionization temperature $T_Z \sim 0.3$ keV.

According to our model, T_1 and T_2 in the indicated case are approximately 50 eV and 1 keV. The equilibrium charge value corresponding to a temperature of 50 eV is equal to ~10 (the neon-like Ca XI ions with an electronic-ground state configuration of $1s^2 2s^2 2p^6$ and an ionization potential of 595 eV are largely present). The ionization in the $l < l^{cr}$ region still has an equilibrium character ($T_e = T_Z$). Then the ions enter the region with the high electron temperature $T_2 \sim 1$ keV (equilibrium charge value $\bar{Z} \approx 18$; predominant at equilibrium at $T_Z \sim 1$ keV is the helium-like Ca XIX ion, which has an ionization potential of 5.13 keV). But the time of stay of the ions in the hot LP layer turns out, as we shall presently show, to be insufficient for the attainment of the equilibrium Z values corresponding to the electron temperature T_2.

Since $T_e > T_Z$ on the right of the discontinuity, the recombination of the ions can be neglected; the relative concentrations of the ions are given by the system of equations:

$$\begin{aligned} \dot{n}_1 &= -s_1 n_1, \\ \dot{n}_2 &= s_1 n_1 - s_2 n_2, \\ \dot{n}_3 &= s_2 n_2 - s_3 n_3, \\ &\ldots\ldots\ldots \end{aligned} \tag{17}$$

Here n_1, n_2, etc. denote the relative concentrations of the ions Ca XI, Ca XII, etc.; $s_Z = N_e \langle v\sigma_i \rangle_Z$. Let us, considering the fact that at $T_e \sim 1$ the ionization rate $\langle v\sigma_i \rangle_Z$ varies comparatively slowly in the Ca XI–Ca XVI ion series, replace the quantities $\langle v\sigma_i \rangle_Z$ by their Z-averaged value $\overline{\langle v\sigma_i \rangle}_Z$. Then the solution to the system (17) with the initial condition $n_k(r = r^{cr}) = \delta_{ik}$ is

$$n_k(r) = \frac{p^{k-1}}{(k-1)!} e^{-p}, \tag{18}$$

where the value of p is given by the integral

$$p = \int_{r^{cr}}^{r} N_e \overline{\langle v\sigma_i \rangle}_Z \frac{dr}{V}, \tag{19}$$

evaluated along the path of an element of volume of the plasma. The integrand decreases by roughly an order of magnitude over the layer from r^{cr} to $1.1 r^{cr}$, and the ionization stops almost completely after this. The estimation of the upper limit of the magnitude of the integral (19) yields $p \sim 4$ (the ionization rate was computed in accordance with Ref. 29). Let us note that the time of stay of the ions in the hot dense LP layer is equal to $\sim 0.1 r^{cr}/c_2 \sim 10^{-10}$ sec. For $p \approx 4$, the relative concentrations of the ions Ca XII, Ca XIII, Ca XIV, CA XV, CA XVI, and Ca XVII are, according to (18), respectively equal to 0.07; 0.15; 0.20; 0.19; 0.16; and 0.10 ($\bar{Z} \approx 13.5$), which are in very good agreement with the ionic composition in experiment.[4,13]

These calculations show that the displacement of the ionic abundance into the region of high ionization multiplicities is indeed limited by the time of stay of the ions in the hot LP layer where $T_e > T_Z$. There comes in the course of the subsequent dispersion a moment when T_e becomes less than T_Z, but because of the low density values the recombination fails to exert any significant influence on the ionic composition.[4]

A more exact computation of the ionic composition in the LP requires the consideration of the effect of the ionization kinetics on the hydrodynamics. A self-consistent hydrodynamic and ionization-kinetics calculation can be carried out in the general case only by numerical methods, and is a subject for a separate paper.

IV. CONCLUSION

In spite of its simplicity, the constructed LP model allows us to fit all the available experimental data. We are able to explain from a unified standpoint the unexpectedly low N_e values in the LP corona (including the value at the Jouguet point), the presence of the rarefaction jump in the vicinity of the critical point, the production of hard x-ray quanta, and the high rates of dispersion of the ions (including the "production" of the so-called "fast" ions). One of the manifestations of the deflagration effect is the ionization's nonequilibrium character, which comes out most distinctly for targets with small A and R and at high q_L values. In future, when x-ray spectral investigations with 1–10-μm spatial resolution become possible, a comparison of the theoretical and experimental profiles of N_e, T_e, and V in the vicinity of the critical point will have to be carried out.

The results of the present paper should stimulate investigations of the microscopic transfer processes occurring in the neighborhood of the critical point and the mechanisms responsible for the suppression of the electronic thermal conductivity. This area of the theory is not fully developed at present.

[1)]This rarefaction jump is accompanied by an increase in the specific entropy as a result of the absorption of the laser-radiation energy.

[2)]The data given in Ref. 21 indicate that M_1 is probably almost independent of ω_L as well.

[1]A. V. Vinogradov, I. I. Sobel'man, and E. A. Yukov, Kvantovaya Elektron. (Moscow) 4, 63 (1977) [Sov. J. Quantum Electron. 4, 32 (1977)].

[2]A. A. Ilyukhin, G. V. Peregudov, E. N. Ragozin, I. I. Sobel'man, and V. A. Chirkov, Pis'ma Zh. Eksp. Teor. Fiz. 25, 569 (1977) [JETP Lett. 25, 535 (1977)].

[3]Yu. V. Afanas'ev, E. G. Gamaliĭ, O. N. Krokhin, and V. B. Rozanov, Zh. Eksp. Teor. Fiz. 71, 594 (1976)[Sov. Phys. JETP 44, 311 (1976)].

[4]E. N. Ragozin, Kvantovaya Elektron. (Moscow) 7, 868 (1980) [Sov. J. Quantum Electr. 10, 493 (1980)].

[5]L. D. Landau and E. M. Lifshitz, Mekhanika sploshnykh sred (Fluid Mechanics), Fizmatgiz, Moscow, 1954 (Eng. Transl., Addison-Wesley, Reading, Mass., 1959).

[6]C. Chu and R. Gross, transl. in: Fizika vysokotemperaturnoĭ plazmy (High-Temperature Plasma Physics), Mir, Moscow, 1972, p. 262.

[7]A. V. Gurevich and Ya. N. Istomin, Zh. Eksp. Teor. Fiz. 77, 933 (1979) [Sov. Phys. JETP 50, 470 (1979)].

[8]L. L. Cowie and C. F. McKee, Astrophys. J. 211, 135 (1977).

[9]B. H. Ripin, P. G. Burkhalter, F. C. Young *et al.*, Phys. Rev. Lett. **34**, 1313 (1975).
[10]R. C. Malone, R. L. McCrory, and R. L. Morse, Phys. Rev. Lett. **34**, 721 (1977).
[11]B. Yaakobi and T. C. Bristow, Phys. Rev. Lett. **38**, 350 (1977).
[12]G. V. Peregudov, M. E. Plotkin, and E. N. Ragozin, Kvantovaya Elektron. (Moscow) **6**, 1284 (1979) [Sov. J. Quantum Electron. **9**, 754 (1979)].
[13]G. V. Peregudov and E. N. Ragozin, Pis'ma Zh. Eksp. Teor. Fiz. **28**, 27 (1978) [JETP Lett. **28**, 26 (1978)].
[14]A. Raven, O. Willi, and P. T. Rumsby, Phys. Rev. Lett. **41**, 554 (1978).
[15]B. H. Ripin, Appl. Phys. Lett. **30**, 134 (1977).
[16]E. N. Ragozin, Kvantovaya Elektron. (Moscow), **4**, 2262 (1977) [Sov. J. Quantum Electron. **7**, 1296 (1977)].
[17]A. V. Vinogradov, G. V. Peregudov, E. N. Ragozin, I. Yu. Skobelev, and E. A. Yukov, Kvantovaya Elektron. (Moscow) **5**, 1077 (1978) [Sov. J. Quantum Electron. **8**, 615 (1978)].
[18]H. Azechi, S. Oda, K. Tanaka, T. Norimatsu, T. Sasaki, T. Yamanaka, and C. Yamanaka, Phys. Rev. Lett. **39**, 1144 (1977).
[19]D. T. Attwood, D. W. Sweeny, J. M. Auerbach, and P. H. Y. Lee, Phys. Rev. Lett. **40**, 184 (1978).
[20]G. V. Peregudov, E. N. Ragozine, I. Yu. Skobelev, A. V. Vinogradov, and E. A. Yukov, J. Phys. D **11**, 2305 (1978).
[21]R. Fedosejevs, M. D. J. Burgess, G. D. Enright, and M. C. Richardson, Phys. Rev. Lett. **43**, 1664 (1979).
[22]V. W. Slivinsky, H. N. Kornblum, and H. D. Shay, J. Appl. Phys. **46**, 1973 (1975).
[23]K. R. Manes, H. G. Ahlstrom, R. A. Haas, and J. E. Holzrichter, J. Opt. Soc. Am. **67**, 717 (1977).
[24]V. A. Boĭko, O. N. Krokhin, S. A. Pikuz, A. Ya. Faenov, and A. Yu. Chugunov, Fiz. Plazmy **1**, 309 (1975) [Sov. J. Plasma Phys. **1**, 165 (1975)].
[25]M. R. Siegrist, B. Luther-Davies, and J. L. Hughes, Opt. Commun. **18**, 603 (1976).
[26]R. Decoste and B. H. Ripin, Appl. Phys. Lett. **31**, 68 (1977).
[27]E. A. McLean, R. Decoste, B. H. Ripin, J. A. Stamper, H. R. Griem, J. M. McMahon, and S. E. Bodner, Appl. Phys. Lett. **31**, 9 (1977).
[28]F. E. Irons and N. J. Peacock, J. Phys. B **7**, 2084 (1974).
[29]L. A. Vaĭnshteĭn, I. I. Sobel'man, and E. A. Yukov, Secheniya vozbuzhdeniya atomov i ionov élektronami (Cross Sections for Excitation of Atoms and Ions by Electrons), Nauka, Moscow, 1973, §§ 25 and 27.

Translated by A. K. Agyei