Laser thermonuclear fusion

A. M. Prokhorov, S. I. Anisimov, and P. P. Pashinin

P.N. Lebedev Physics Institute, USSR Academy of Sciences, Moscow
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A review is given of the current status of theoretical and experimental investigations of laser thermonuclear fusion. A wide range of physical problems encountered in fusion research and ways of solving them are discussed. Attention is concentrated on current problems such as fusion energetics, adiabatic target compression, conditions for development of hydrodynamic instabilities, absorption of laser radiation in plasma, and plasma heating. Detailed analysis is made of the results of numerical fusion simulations that have been carried out subject to various assumptions relating to the type of target, laser pulse parameters, and predominant processes in the laser plasma. Some results are given of recent experimental work on laser fusion.

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1. INTRODUCTION

The problem of controlled thermonuclear fusion is undoubtedly one of the central problems in modern physics. Its enormous attraction for scientists lies in the promise of utilization of unlimited thermonuclear resources of energy on the earth to be released by the fusion of light atomic nuclei to form heavier nuclei. The nuclear fuels would be mainly the heavy isotopes of hydrogen, deuterium and tritium. The reserves of deuterium are so enormous that they would be sufficient to serve humanity for millions of years even at a very high rate of energy consumption. However, exceptional difficulties will have to be overcome before controlled fusion is achieved. The difficulties arise because the fusion process can occur only if two nuclei approach each other to a distance of the same order as their size which is 10^{-13} cm. This is possible if positively charged nuclei overcome the mutual electrostatic repulsion, i.e., if they have a sufficiently high energy. The energy can be provided by heating matter to a very high temperature when the kinetic energy of the nuclei is sufficient to overcome the electrostatic repulsion in their collisions. In nature these conditions occur in the interiors of stars and man has been using for a long time the energy of one such thermonuclear reactor in the form of the radiant energy of the sun. Having investigated the nature of this energy source, man has quite rapidly reproduced this process on the earth in creating the most powerful weapon, which is the thermonuclear bomb. However, attempts to achieve controlled fusion have made it necessary to perform exceptionally extensive fundamental investigations in nuclear and plasma physics, to build unique equipment, and to carry out a very wide range of theoretical investigations.

The work on controlled fusion started 25 years ago when the first ideas for plasma containment were put forward. These investigations were headed by leading Soviet physicists, particularly by I. V. Kurchatov, who ensured an extensive approach to the problem and attracted talented young scientists. For many years the work was directed successfully by L. A. Artsimovich and the school of M. A. Leontovich tackled the theoretical aspects of the problem. The idea of magnetic confinement and thermal isolation of hot plasma was put forward in the Soviet Union[1]: this idea is the basis of steady-state systems in which the fusion of deuterium and tritium should occur in the form of slow "combustion." Enormous difficulties had to be overcome due to hydrodynamic instabilities of various equilibrium plasma configurations and anomalous plasma diffusion in a magnetic field. Determined research and development made it possible to accumulate a considerable amount of knowledge of the nature of the processes in plasma in such systems and to determine the most promising directions for further work. The Soviet Tokamak program, in which the main difficulties of the stationary systems are tackled in realistic ways, [2] occupies a leading position. Work in accordance with this program is now very extensive and has become one of the major directions of thermonuclear research also in the USA, Europe, and Japan.

In recent years a basically different approach, which avoids the above difficulties associated with magnetohydrodynamic instabilities and anomalous diffusion, has been increasing in importance. This approach was first proposed in the fifties and it rejects the idea of magnetic isolation (and confinement) but proposes that the fusion reaction be carried out in the pulse regime in which energy is evolved in a series of moderate-power explosions. Clearly, the main difficulty in this approach is purely quantitative and it arises from the question how to achieve a significant degree of utilization of the thermonuclear fuel by means of "microexplosions" of relatively low energy. This approach requires creating an extremely high energy density for a short time and in a small amount of matter. Attempts to solve this problem in the fifties by conventional electrical engineering methods have not been successful and, therefore, the pulse approach has for long been regarded as of little promise. The situation has changed drastically in the last decade when new efficient methods of energy concentration have been developed and these comprise high-power lasers, high-current pulsed relativistic electron beams, and cumulation method for creating megagauss magnetic fields and pressures amounting to millions of atmospheres. This has made it possible to formulate new promising approaches to the solution of the problem of controlled fusion which are now being investigated very intensively so that one may expect the solution to the principal physcal problems in the next five years.

Work on the generation of high-temperature dense plasma by lasers has been going on for over a decade. In the early sixties estimates have been obtained and it has been shown theoretically that, in principle, it should be possible to heat plasma to temperatures of the order of 10^7 °K by irradiation of solid deuterium targets with lasers.[3,4] Soon after investigations of laser breakdown in gases have shown that this method readily produces plasmas with temperatures of several hundreds of electron-volts and densities of 10^{20} cm^{-3}.[5-7] In the late sixties the first neutrons were recorded as a result of action of laser radiation on solid targets.[8]

Although heating by itself is a necessary rather than a sufficient condition for effective initiation of thermonuclear reactions, these investigations have provided a very strong stimulus for theoretical and experimental studies of the interaction of laser radiation with plasma. The greatest and most successful efforts have been made in the Soviet Union, USA, France, West Germany, and Japan. Experimental studies have been made of the mechanisms of the absorption of light in an inhomogeneous superdense plasma and of gasdynamic expansion of matter heated by a laser; semianalytic theories and numerical calculation methods have been developed. This has made it possible to estimate the conditions for reaching the break-even physical threshold of thermonuclear reactions, i.e., an estimate has been obtained of the minimum laser energy needed to obtain a thermonuclear energy equal to the laser energy. According to[9], the threshold laser energy required for a pellet composed of equal amounts of deuterium and tritium is 10^8 J, which means that laser fusion is hardly likely to be achieved by means of a simple laser scheme. The idea of heat-conduction heating of a superdense plasma by short laser pulses is considered in[10]. This approach has no significant advantages in the sense of reduction of the minimum threshold energy. It is proposed in[11,12] that inertial confinement of a dense plasma by a heavy cylindrical shell, ensuring one-dimensional plasma expansion, is needed for a significant reduction in the threshold value of the laser energy. It is proposed to reduce the losses due to radial heat conduction in a cylindrical plasma column by the use of longitudinal "moderate" magnetic field pulses of 10^6 Oe intensity. In this variant the necessary laser energy is about 1.5 orders of magnitude less but it is still very high. We shall not consider several other calculations which have been reviewed elsewhere.[13-15]

A very fruitful idea of using laser radiation not only to heat a target but also to produce simultaneously a very strong compression of a spherical fusion fuel target was put forward by J. Nuckolls *et al.*[16] 1) We shall show later than an increase in the density of the fusion fuel makes it possible to reduce considerably the threshold laser energy. Nuckolls *et al.* suggested a time-programed irradiation of a spherically symmetric target[16] and they showed that the reactive force of the expanding hot plasma could be used to compress almost adiabatically the central core of the target to a density 10^2–10^4 times higher than the density of a solid and to thus initiate a thermonuclear "microexplosion" with a positive energy yield. The threshold laser energy of this scheme is many orders of magnitude lower and lies in the range of 10^3–10^4 J for pulses of 10^{-10} sec duration, i.e., the required powers are 10^{13}–10^{14} W. In the same period a considerable progress has been made in developing laser systems emitting nanosecond and picosecond pulses of 10^2–10^3 J energy and high directionality. This very dramatic narrowing of the gap between the theoretical estimates of the threshold laser energy and technological capabilities has altered basically the direction of work carried out on controlled fusion in several countries. Special laser units have been built to check the physical principles underlying the new approach to laser fusion. Theoretical investigations, directed to developing adequate mathematical models of laser microexplosions, have been intensified. Many papers have been published essentially on the calculation of the thermonuclear yield under various conditions and on optimal target configurations as well as on optimal laser pulse parameters. Calculations of this kind are very difficult to carry out. Moreover, their correct formulation requires solution of a number of complex physical problems, particularly in plasma physics. Although not all these problems have yet been solved, we can now speak of the promise of laser fusion and its competitiveness with other approaches to controlled fusion. Naturally, much experimental work still remains to be done on all the physical effects resulting from laser compression and heating of matter. The experimental difficulties are not only due to the

1)It should be noted that the idea of using a laser to compress a target was put forward earlier by Daiber *et al.*[17] However, they did not carry out sufficiently reliable calculations and their work remained practically unnoticed.

exceptionally rigorous requirements that laser systems have to satisfy but are also due to the fact that all the important processes occur in time intervals of 10^{-9}–10^{-12} sec on a spatial scale of 10^{-2}–10^{-4} cm.

We shall review a number of the major physical problems which are facing laser fusion and consider ways of solving them, as perceived at present.

2. SIMPLE ESTIMATE OF THRESHOLD ENERGY

We shall begin a more detailed discussion of the energetics of laser fusion from a simple model. Essentially, this model is based on the same considerations that are usually employed in the derivation of the well-known Lawson criterion. We shall consider a homogeneous pellet made of a deuterium–tritium mixture, heated to a temperature of several kiloelectron-volt. The following reaction takes place in the pellet:

$$\mathrm{D} + \mathrm{T} \to \mathrm{He}^4\,(3.5) + n\,(14.1)$$

(the energies of the reaction products, in megaelectron-volts, are given in parentheses). This reaction produces energy per unit volume

$$dE_{\mathrm{TN}} = \varepsilon_0 n_{\mathrm{D}} n_{\mathrm{T}} \langle \sigma v\rangle\, dt,$$

where ε_0 is the energy liberated in one reaction event, $\sigma(v)$ is the reaction cross section, and $\langle\cdots\rangle$ denotes averaging over the Maxwellian distribution of velocities. We shall initially ignore the heating of the particles during this reaction. Then, assuming that the reaction time τ is of the same order of magnitude as the hydrodynamic expansion time of the pellet r_0/c_s (r_0 is the pellet radius and c_s is the velocity of sound), and assuming that the densities of the deuterium and tritium particles are $n_{\mathrm{D}} = n_{\mathrm{T}} = n/2$, we find that the density of the energy evolved as a result of this reaction is

$$E_{\mathrm{TN}} \approx \frac{1}{4}\varepsilon_0 n^2 \langle\sigma v\rangle \frac{r_0}{c_s}.$$

We shall introduce the concept of an energy gain G, equal to the energy E_{TN} obtained as a result of the thermonuclear reaction, to the energy E_L of the initiating laser pulse:

$$G = \frac{E_{\mathrm{TN}}}{E_L} = \frac{\varepsilon_0 n \langle\sigma v\rangle r_0 \eta}{12 k c_s T_0}, \tag{1}$$

where η represents the fraction of the laser energy transformed by absorption into the thermal energy of the plasma. It follows from Eq. (1) that

$$r_0 n = G\eta^{-1}\psi(T_0).$$

and, as is easily shown, the function $\psi(T_0)$ has a minimum at $T_0 \approx 2\times 10^8\,{}^\circ$K which is approximately 10^{22} cm^{-2} (ε_0 is assumed to be equal to the α-particle energy). Now, assuming that $r_0 n$ represents the lowest value of this product, we can calculate the minimum laser energy required to reach an energy gain G. The following formula is obtained for E_L:

$$E_L = G^3\eta^{-4}\left(\frac{n_s}{n_0}\right)^2\cdot 10\ \text{(MJ)} \tag{2}$$

where n_s is the density of particles in solid hydrogen.

3. ALLOWANCE FOR HEAT EVOLUTION DUE TO REACTION

A self-evident shortcoming of the above estimates is our failure to allow for the plasma heating as a result of the thermonuclear reaction occurring in it, i.e., the above estimates are obtained on the assumption that the process is isothermal. In fact, the reaction is almost adiabatic. The heating of the plasma by the energy evolved in the reaction clearly accelerates the reaction and increases both the degree of fuel burn-up and the thermonuclear yield. We have in fact considered the linear stage of the process when the thermonuclear energy is less than the initial thermal energy of the plasma. We must refine this approach by allowing for the change in the plasma temperature during the reaction. The change is due to the deceleration of the α particles in the reacting substance. We can assume approximately[18] that an α particle with a range $l(n, T)$ gives up, in a sphere of radius r, an energy of the order of $\varepsilon_0 r(r+l)^{-1}$. Then, the temperature equation can be expressed in the form

$$3nk\dot{T} = \frac{\varepsilon_0 r}{r+l}\langle\sigma v\rangle \frac{n^2}{4}\xi^2, \tag{3}$$

where the factor $\xi = 2n_{\mathrm{D}}/n$ is introduced to allow for the burn-up of the reacting nuclei.

The size of the region in which the reaction takes place varies with time. However, as before, we shall ignore this change and assume that $r = r_0$. If the deceleration of the α particles is due to their interaction with the target electrons, we obtain the following expression for the α-particle range (the details of the calculation and the range of its validity are given in[19]):

$$l(n, T) = l_0 T^{3/2} n^{-1},\qquad l_0 = 5\cdot 10^{10}\ (\mathrm{cm}^{-2}\cdot\mathrm{deg}^{-3/2}).$$

Replacing in Eq. (3) the variable in accordance with the relationship $dr = -c_s(T)dt$ and integrating, we obtain

$$\begin{gathered} nr_0 = \int\limits_{T_0}^{T_m} \frac{c_s(T)\,dt}{\langle\sigma v\rangle} F(T, T_0). \\ F(T, T_0) = \frac{12k(1+\beta T^{3/2})}{\{\varepsilon_0 - 3k[T - T_0 + 0.4\beta(T^{5/2} - T_0^{5/2})]\}^2}. \\ \beta = \frac{l_0}{nr_0}. \end{gathered} \tag{4}$$

We shall now calculate E_{TN}. Since $dE_{\mathrm{TN}} = 3k[1 + l(n, T)/r_0] \times(\varepsilon_\alpha/\varepsilon_0)dT$, we find that

$$\frac{E_{\mathrm{TN}}}{\eta E_L} = \frac{G}{\eta} = 5\left[\frac{T_m}{T_0} - 1 + 0.4\frac{\beta}{T_0}(T_m^{5/2} - T_0^{5/2})\right]. \tag{5}$$

Next, using Eq. (4), we obtain the following expression for the laser energy

$$E_L = \frac{4}{3}\pi r_0^3\cdot 3nkT_0\eta^{-1},$$

which—together with Eq. (5)—is the parametric equation for the energy E_L expressed as a function of η, G, T_0, and n. We shall not consider the detailed calculations of the values of this function but note that the reaction self-acceleration effect depends considerably on the value of the ratio G/η. If $G/\eta \approx 1$, the laser energy calculated allowing for this effect is practically identical with that given by Eq. (2). However, if $G/\eta \approx 100$, Eq. (2) underestimates the results by a factor of about 100. In the roughest approximation an estimate of the laser energy allowing for burn-up can be expressed in the form

$$E_L = 10G^2\eta^{-3}\left(\frac{n_s}{n_0}\right)^2 \text{ (MJ)}.$$

Two conclusions can be drawn from the above estimates. Firstly, a significant thermonuclear yield in a reasonable range of laser microexplosion energies can be obtained using a DT plasma compressed to densities much higher than the density of solid hydrogen n_s. Secondly, the laser energy required for the initiation of thermonuclear fusion depends strongly on the plasma absorptivity and becomes unacceptably large if $\eta \ll 1$ even if compression is employed. Thus, from the practical point of view, the laser initiation of thermonuclear fusion requires 1) compression of the reactive substance to a density of $n \gg n_s$ and 2) ensuring an efficient transfer of energy from laser radiation to the plasma.

4. ADIABATIC COMPRESSION

We shall first consider the problem of compression. We can easily show that compression of solid hydrogen to a density $10^3 n_s$ at a temperature of several kiloelectron-volts requires a pressure of the order of 10^{11} atm. Such high pressures can be produced, in principle, by the reactive impulse produced by the expansion of the outer absorbing layer of the target. However, we must bear in mind the following difficulty. It is well known that, in contrast to a rarefaction wave, a compression wave is always transformed into a shock wave and this happens in a finite time. Thus, when a target is compressed by the reactive impulse, strong shock waves are generated (they were indeed observed experimentally[20,21]) and these waves cause irreversible heating of the reacting medium which will obviously prevent its further compression. Therefore, high degrees of compression can be reached by avoiding formation of strong shock waves and making the process as close to isentropic as possible.

A reasonable approximation to an isentropic process may be obtained[22,23] by a suitable selection of the time dependence of the compressive stress (in the laser experiments this requires programing of the laser pulse shape). This is the approach considered by Nuckolls *et al.*[16] Another approach involves a special density profile in which strong shock waves do not penetrate into the inner layers of the target.[24] In particular, the use of multilayer shell targets, discussed frequently in recent years, should be very effective.

The possibility of adiabatic compression with a suitable supply of energy can be demonstrated most easily in the case of an ideal gas compressed by a spherical piston. We shall consider the particular self-similar solution of the gasdynamic equations which depends on the variable $\xi = r/t$ and has a spherical symmetry. We can easily show that this solution describes an adiabatic compression wave traveling from the piston to the center of symmetry. For a particular law of motion of the piston the flow is adiabatic right to the moment of total collapse and the degree of compression can be as large as we please. The law of motion of the piston can easily be obtained by integrating the equation for the piston radius $\dot{R} = v(R, t)$, where the velocity of the medium $v(\xi)$ found by solving the gasdynamic equations is substituted on the right-hand side. The asymptotic form of the dependence $R(t)$ is given in[26].[2)] This dependence can be obtained from qualitative considerations without solving the gasdynamic problem.

We shall now consider the final stage of compression when the average density of matter is much higher than the initial density. Clearly, during this stage the nature of the compression cannot depend on the initial parameters of the medium. In particular, there is now no parameter with the dimensions of velocity (all the velocities are much higher than the initial velocity of sound c_{s0}). Therefore, the velocity of sound c_s, rising during compression, should be proportional to r/t (the coordinates are selected so that the wave collapses at the point $r = 0$ at $t = 0$). Since during adiabatic motion the velocity of sound is $c_s = c_s(\rho) \propto \rho^{(\gamma-1)/2}$, it follows that $\rho \propto (R/t)^{2/(\gamma-1)}$. On the other hand, the law of conservation of mass gives $\rho \propto R^{-3}$. The last two relations yield the law of motion of the piston[26] $R \propto t^{\mu/(\mu+3)}$ and the time dependences of the average density and average pressure:

$$\rho \propto t^{-3\mu/(3+\mu)}, \quad p \propto t^{-3\gamma\mu/(3+\mu)},$$

where $\mu = 2/(\gamma - 1)$. The same results are naturally obtained by a rigorous solution of the model hydrodynamic problem.[27] They can be used to calculate the power needed for compression. It is assumed then that throughout the compression process the power represents a constant fraction of the total laser power, which gives the following estimate of the time dependence of the laser radiation flux:

$$\dot{Q} = p\dot{R}R^2 \propto t^{-(9+\mu)/(3+\mu)}. \tag{6}$$

In the case of spherically symmetric compression of an ideal gas with $\gamma = 5/3$ it follows from Eq. (6) that $\dot{Q} \propto t^{-2}$, which is practically identical with that obtained as a result of numerical optimization of the compression regime.[16]

A detailed study of the hydrodynamics of compression

2)It should be noted that one of the first references to plane centered compression waves was given in a monograph by Courant and Friedrichs.[22] The solution considered here can be obtained formally by time inversion of the well-known Zel'dovich equation for the flow behind a spherical detonation front.[25]

can also be made for more complex cases. In particular, it is possible to obtain self-similar solutions analogous to those described above in the case when the initial density is distributed in accordance with a power law. The flow involving converging spherical shock waves has been investigated in detail for this case.[28] The problem of collapse of a hollow spherical shell is of special interest in connection with laser fusion. A detailed study of this problem is reported in[29].

Interesting particular solutions of the gasdynamic equations, describing adiabatic compression of a finite mass of a gas with an inhomogeneous density distribution can be found in[30–33]. They also give the optimal shape of a laser pulse, which is close to that found as a result of numerical calculations.

5. HYDRODYNAMIC INSTABILITY OF TARGET COMPRESSION

It is clear that the model considered above, like many other analytic or semianalytic models, are far too simplified and are unsuitable for quantitative calculations of laser compression. Nevertheless, many important features of the dynamics of real compression can be described correctly by the simplest model of compression produced by a spherical piston. This applies particularly to the energy supply regime ensuring adiabatic compression. In this aspect the numerical calculation and the simplest estimates give approximately the same results. A more detailed comparison of one of the models with results of a numerical calculation[32] shows that the model solution describes quite correctly the motion of matter in the central compressed part of a laser target.

Another important feature of the compression process, which is also predicted as a result of analyses of simple models, is the possibility of hydrodynamic instabilities which destroy the optimal compression regime. It is found that compression produces hydrodynamic instabilities of two kinds. The first kind is associated with an increase in one-dimensional perturbations which have the same symmetry as the main motion. Examples of such an instability are described in[28,34]. This kind includes, in particular, an instability which displaces the point of collapse. The corresponding perturbations of hydrodynamic variables increase in the linear approximation as t^{-1}. Numerical calculations show that this one-dimensional instability does not alter basically the compression process. An instability of a different kind, which disturbs the spherical symmetry of compression, is much more dangerous. This is known in hydrodynamics as the Rayleigh–Taylor instability. Its simplest manifestation is the instability of a boundary between two media with different densities in a gravity field.[35,36] If a heavier liquid lies below a lighter one, small perturbations of the interface are gravitational waves with the dispersion law $\omega = \sqrt{gk}$ (g is the acceleration due to gravity). The opposite case, when a heavier liquid is above, corresponds to a reversal of the sign of g; the interface is then unstable and its perturbations rise exponentially with an increment proportional to g. In laser compression the gravity field is replaced by the acceleration of matter due to the pressure gradient and the interface is now the somewhat arbitrary boundary between the dense core and the rarefield corona of a laser target. Compression becomes unstable if the acceleration is directed parallel to the density gradient. In real situations the development of this instability is influenced strongly by the compressibility, viscosity, and thermal conductivity of the plasma, fluctuations of the acceleration, and the existence of a flux of matter across the boundary of the dense core of the target. All these factors are allowed for to a greater or lesser extent in[37–40], where a linear analysis is made of the Rayleigh–Taylor instability. The nature of the growth of perturbations depends strongly on the wavelength (or number) of a spherical harmonic. The perturbations with the longest wavelengths, comparable with the target radius, increase significantly in a time of the order of the total compression time. Estimates given in[35] demonstrate that the amplitude of a perturbation which transforms a sphere into an ellipsoid (second spherical harmonic) increases approximately by a factor of $\sqrt{m}$ when the average density is altered by a factor m. Thus, in the case of a sufficiently small initial amplitude, the long-wavelength perturbations do not present a serious danger in the compression process. The fastest-growing short-wavelength perturbations also do not disturb greatly the compression regime because they become stabilized by dissipative effects and because the thickness of the transition layer between the core and the corona is finite. Therefore, in the case of homogeneous targets the most serious problem is presented by perturbations whose wavelength lies in the range between the core size and the transition layer thickness. There have been several studies of the growth of a Rayleigh–Taylor instability by numerical solution of two-dimensional hydrodynamic equations.[41,42] Although this approach deals only with the growth of perturbations which are independent of the azimuthal angle, calculations of this kind give some information on the influence of such an instability on the processes in a laser target and, in the final analysis, on the thermonuclear yield. The majority opinion is that the hydrodynamic instability reduces very greatly the energy gain of simple targets and imposes very stringent requirements on their initial sphericity and homogeneity of irradiation. Experimental study of the Rayleigh–Taylor instability under laser compression conditions have not yet been made so that the correctness of this view has not yet been tested.

Recently, many investigators have turned to complex shell targets because the use of heavy shells makes it possible to relax significantly the requirements in respect of the time dependences of laser pulses. However, one must bear in mind that the Rayleigh–Taylor instability problem is likely to be particularly acute in the case of shell targets because of the steep density gradients existing in these targets right from the beginning. During the initial stage of compression a shell is accelerated by the recoil impulse produced by the material being evaporated. This acceleration gives rise to an instability of the outer part of the shell.

Near the compression maximum the accelerated heavy shell experiences strong braking and then an instability appears at the boundary between the shell and the thermonuclear fuel. Both instability stages influence considerably the dynamics of compression and thermonuclear yield. Some idea of the unstable compression of thin shells can be obtained by investigating the motion of zero-thickness shells under the action of an external pressure. Such a study[43] shows that there are always perturbations whose development causes the shell to break up in a time shorter than the collapse time.

A fairly extensive class of perturbations disturbing the symmetry of compression of a heavy spherical shell was investigated numerically by Thiessen *et al.*[41] They showed that the instability effects limit very considerably the energy gain of shell targets and they proposed a way of achieving strong compression in which a laser pulse is replaced by a series of ultrashort pulses of rising amplitude. We can easily see that during each such ultrashort pulse the amplitude of "dangerous" perturbations cannot increase too much, so that the whole process is more stable.

An interesting manifestation of the small-scale Rayleigh–Taylor instability is the turbulent mixing of two media separated by an unstable boundary. The nonlinear instability stage can be described then as a mutual diffusion of the media resulting in equalization of the density gradient. A theory of turbulent mixing based on these ideas is developed in[45] and the relevant diffusion coefficient is calculated. An attempt to allow for turbulent mixing in the hydrodynamic problem of laser target compression is made in[46]. Introduction in numerical calculations of the effects of turbulent mixing alters considerably the integrated characteristics of the compression process: it lowers the density, reduces the thermonuclear yield, and alters the characteristics of the x-ray radiation emitted from the target.

We can summarize by stating that the influence of the Rayleigh–Taylor instability on the processes occurring in laser targets is undoubtedly one of the most important. Estimates and numerical calculations show that this instability may influence strongly the behavior of a target. However, the quantitative results obtained so far have to be refined still further.

6. MECHANISMS OF ABSORPTION OF LASER RADIATION IN PLASMA

We shall now consider another important subject which is specific to laser fusion: we shall discuss the problem of absorption of laser radiation in a plasma corona.

When the intensity of laser radiation is not too high, its absorption in a plasma is due to Coulomb collisions between particles. The absorbed power is proportional to the electron collision frequency ν_{ei}, which decreases with rising electron temperature as $T_e^{-3/2}$.[47] Therefore, at high temperatures the usual collisional absorption mechanism becomes ineffective.

A plane electromagnetic wave incident normally on an inhomogeneous plasma layer is reflected at a point where the wave frequency is equal to the local electron plasma frequency $\omega_p = \sqrt{4\pi e^2 n_e/m}$.[48] For neodymium laser radiation the corresponding critical value of the electron density is $n_c = 10^{21}$ cm^{-3}; the absorption of light occurs only if $n_e \leq n_c$. The local absorption coefficient for moderate intensities can be presented in the form[48]

$$\varkappa = \frac{\omega_p^2 \nu_{ei}}{c\omega^2 \sqrt{1-(\omega_p^2/\omega^2)}}.$$

The optical thickness of a plane plasma layer is

$$\Lambda = 2\int_{-\infty}^{x_0} \frac{\omega_p(x)\,\nu_{ei}(x)\,dx}{c\omega^2\sqrt{1-(\omega_p^2(x)/\omega^2)}}, \tag{7}$$

where x_0 is the point at which the local plasma frequency is equal to the frequency of light $\omega_p(x_0) = \omega$. We can easily see that the main contribution to the integral (7) is made by the region in the vicinity of the point x_0. Expanding $\omega_p^2(x)$ and $\nu_{ei}(x)$ as series near the point x_0 and integrating, we obtain

$$\Lambda \approx \frac{2L\nu_{ei}}{c},$$

where $L = [d(\ln n_e)/dx]^{-1}$ and the values of L and ν_{ei} are taken at the point x_0. The radiation flux absorbed in the plasma is clearly $\dot{Q}_0(1-e^{-\Lambda})$, where $\dot{Q}_0$ is the incident flux.

In a rough estimate of the plasma temperature we shall assume that all the radiation is absorbed near the point where the density is critical. Equating the absorbed radiation flux to the hydrodynamic energy flow at the point x_0, we obtain

$$T_e \approx 1.3\cdot 10^{12} A^{1/3}\left[\frac{Z\dot{Q}_0(1-e^{-\Lambda})}{n_c(Z+1)}\right]^{2/3}, \tag{8}$$

where Z is the charge and A is the mass number of ions. Substituting Eq. (8) into the expression for the collision frequency, we obtain the following equation for Λ:

$$\Lambda(1-e^{-\Lambda}) = \frac{\dot{Q}_0^*}{\dot{Q}_0}, \quad \dot{Q}_0^* = \frac{5\cdot 10^{15} L(Z+1)}{\sqrt{A}}.$$

The plasma absorptivity thus depends on the relationship between $\dot{Q}_0$ and $\dot{Q}_0^*$. If $\dot{Q}_0 \ll \dot{Q}_0^*$ (which, in the case of a DT mixture and $L \approx 100\ \mu$, corresponds to $\dot{Q}_0 \ll 10^{14}$ W/cm^2), we have $\Lambda \approx \dot{Q}_0^*/\dot{Q}_0 \gg 1$, i.e., the radiation is absorbed almost completely as a result of Coulomb collisions between electrons and ions. However, if $\dot{Q}_0 \gg \dot{Q}_0^*$, then $\Lambda \approx \sqrt{\dot{Q}_0^*/\dot{Q}_0} < 1$ and we have a strong reflection of light.

This estimate is qualitative. A more accurate result can be obtained by numerical hydrodynamic calculations. Figure 1 shows, by way of example, the dependence of the optical thickness and absorptivity of a plane DT plasma layer on the laser radiation intensity (neodymium laser, pulse duration 1 nsec, only Coulomb collisions are allowed for). It follows from this discussion that at intensities of the order of 10^{15}–10^{16} W/cm^2 the usual inverse bremsstrahlung absorption mechanism becomes ineffective because of a strong reduction in the frequency of electron–ion collisions.

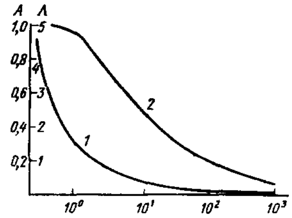

FIG. 1. Dependence of the optical thickness $\Lambda(1)$ and absorptivity $A(2)$ of a plane DT plasma layer on the intensity of laser radiation (expressed in units of $\dot{Q}_0/\dot{Q}_0^*$). The curves are plotted for $\lambda = 1.06\ \mu$ and pulse duration $\tau = 1$ nsec.

However, it should be pointed out that a reduction in the collision frequency reduces the damping and creates favorable conditions for the excitation of plasma waves. Thus, it is in the range of laser intensities where the collisional absorption ceases to operate that we can expect collective absorption associated with the excitation of plasma waves. Recent investigations indicate that this mechanism is very important to the whole laser fusion problem. We shall now consider the collective absorption mechanism in somewhat greater detail. It is known[49,50] that a plasma subjected to a strong electromagnetic wave is unstable. This instability is manifested by spontaneous growth of the amplitudes of plasma waves with wave vectors $\mathbf{k}_1$ and $\mathbf{k}_2$ satisfying the condition $\mathbf{k}_1 + \mathbf{k}_2 = \mathbf{k}_0$, where $\mathbf{k}_0$ is the wave vector of the incident (pump) electromagnetic wave. The frequencies of the excited plasma waves are related by the parametric resonance condition. In particular, if the pump amplitude is relatively low, the main result is the generation of waves whose frequencies are $\omega_1 + \omega_2 = \omega_0$ (ω_0 is the pump wave frequency). In this case the relationships governing ω and $\mathbf{k}$ may be interpreted as the laws of conservation of energy and momentum in the process of the decay of the original electromagnetic wave into two plasma waves. At high pump amplitudes a band of allowed values of $\omega_1 + \omega_2$ forms near ω_0.

Two variants of the parametric instability are of practical importance in laser fusion problems. In the range where the electron density is close to the critical value n_c, an electromagnetic wave may decay into a Langmuir wave and an ionic-sound wave: $\omega_0 = \omega_p + \omega_s$, $\mathbf{k}_p \approx -\mathbf{k}_s$ (since $\mathbf{k}_0 \ll \mathbf{k}_p$ at $\omega_0 \approx \omega_p$). The second process occurs in the vicinity of a point where the electron density is $n_e \approx n_c/4$ and it involves a decay of an electromagnetic wave into two Langmuir waves with similar frequencies $\omega_{p1} \approx \omega_{p2} \approx \omega_0/2$ and wave vectors $\mathbf{k}_{p1} \approx -\mathbf{k}_{p2}$. These two processes result in conversion of the energy of an incident electromagnetic wave into the energy of plasma oscillations. They are analogous to the usual light-absorption processes because the plasma waves resulting from such decay do not escape outside and their energy is finally transformed into the energy of the plasma particles.

This parametric excitation of waves occurs if the amplitude of the pump wave exceeds a certain threshold. The existence of this threshold is associated with damping of the excited waves. The thresholds of the principal instabilities have been calculated by many authors; a summary of the results can be found in[18,49,50]. A linear theory of the decay of an electromagnetic wave into a Langmuir wave and an ionic-sound wave in a nonisothermal plasma with $T_e \gg T_i$ gives the following value of the laser intensity threshold:

$$\dot{Q}_t = \frac{2cn_e kT_e \nu_s \nu_p}{\omega_s \omega_p},$$

where ν_s and ν_p are the damping decrements of the excited waves. When the threshold is exceeded by a small amount, the instability increment increases linearly with the intensity:

$$\gamma_i = \frac{1}{2}\frac{\nu_s \nu_p}{\nu_s + \nu_p}\left(\frac{\dot{Q}}{\dot{Q}_t} - 1\right).$$

A basically similar situation arises also in the decay of an electromagnetic wave into two Langmuir waves, which occurs when the electron density is $n_e \approx n_c/4$. In this case an important feature is that the instability increment depends on the directions of the wave vectors of the excited waves relative to the wave vector of the electromagnetic pump wave and it vanishes (for a spatially homogeneous electron density) in the limit $k_0 \to 0$. The formulas for the threshold and increment, and an analysis of the nonlinear growth of this instability can be found in[51].

It should be pointed out that, in addition to the instabilities considered above, which result in the generation of plasma waves and absorption of light, we can also have instabilities of the stimulated scattering type which generally result in energy losses. In the range of electron densities $n_e < n_c$ this may occur due to the stimulated Brillouin scattering $\omega_0 \to \omega_0' + \omega_s$, which represents the decay of an electromagnetic wave into an ionic-sound wave ω_s and a scattered electromagnetic wave ω_0'. This results in a frequency shift of the order of $c_s k_0$, where c_s is the velocity of sound. The threshold of this process is close to the threshold of the decay of a photon into a plasmon and a phonon.

Another scattering process which occurs at densities $n_e \lesssim n_c/4$ is the stimulated Raman scattering of laser radiation accompanied by the generation of plasmons of frequency $\omega_p \lesssim \omega_0/2$. The threshold of this process is of the order of the threshold of the two-plasmon decay of a photon.

So far we have considered only the threshold conditions and the behavior of instabilities in the vicinity of this threshold. A much more complex and more important in practice is the problem of nonlinear growth of parametric instabilities and of saturation mechanisms which govern the degree of turbulence and anomalous collisions frequency. This problem is discussed in many papers (see[49-57]), in which numerical simulation of plasma processes is used extensively[54-57] in addition to analytic methods. The instability associated with the decay of a pump wave into Langmuir and ionic-sound waves has been investigated most thoroughly. A one-dimensional numerical simulation of this decay in the strong pump case can be found in[57]. The saturation mechanism considered in[57] is the capture of electrons by unstable plasma waves. A similar problem is treated in[54,55]. The anomalous collision frequencies found in the strong-field limit in[54–57] are of the order of $2\omega_{pi}$. The weak pump case is considered also in the one-dimensional formulation in[58]. The plasma wave spectrum and the energy distribution of electrons are found to be quite different than in the strong pump case. It is concluded in[58] that the rate of transfer of energy to electrons from the pump field and the shape of the electron distribution "tails" are in agreement with the predictions of the quasilinear theory.

A one-dimensional numerical simulation of a parametric instability is carried out in[55,59] allowing for the Coulomb collisions in the Bernstein–Greene–Kruskal (BGK) model. The Coulomb collisions play an important role when the pump field is weak. In strong pump fields the electron heating is mainly due to their interaction with plasma waves and this interaction is usually described by introducing an "anomalous" collision frequency. In this case allowance for the Coulomb collisions makes it possible to utilize more satisfactorily the results of one-dimensional simulation in calculations of the anomalous collision frequency in a real three-dimensional plasma.

An anomalous absorption of electromagnetic waves associated with the parametric instability has been observed experimentally in the rf[60,61] and infrared[62] ranges. The results of some experimental studies of the reflection of high-power laser pulses from plasmas[63,64] are also in agreement with the assumption that the absorption is associated with the generation of plasma waves.

In analyzing the anomalous absorption of light due to the parametric instabilities we meet a serious problem relating to the dissipation of the energy of plasma waves. We shall not go into details of this problem (see[65,66]) but note that the transfer of energy from the waves to the particles is accompanied by the formation of a group of fast electrons with an average energy exceeding by an order of magnitude or more the average thermal energy of the bulk of the electrons. This effect may influence considerably the degree of compression of the core of a laser target because the fast electrons with a greater range can heat the core and increase the pressure without increasing the density. An analysis of the influence of the fast electrons on the compression process can be found in[67].

We shall conclude by noting that in addition to the above "reflection" instabilities, we find that when the laser radiation intensity of the order of 10^{15}–10^{16} W/cm^2, its wavelength is 1 μ, and the density gradient in the corona is low, a strong reflection of the laser radiation may occur because of the stimulated Compton scattering.[68,69]

7. THERMAL CONDUCTIVITY OF LASER PLASMA

Heat conduction is one of the most important processes of laser heating and compression of a thermonuclear plasma. The point is that for all the currently available high-power lasers the critical electron density is many orders of magnitude lower than the density in the target core so that the transfer of energy from the absorption zone to the boundary of the core is entirely due to heat conduction. High degrees of compression require, as mentioned earlier, that the matter ahead of the compression wave should remain relatively cold. This means that a heat wave should not overtake the compression wave, i.e., that the propagation of the heat wave should be subsonic. Such a heat wave acts like a piston: the acoustic perturbations emitted from it form a compression wave.

Correct allowance for the electron thermal conductivity, which is important in the selection of the optimal regime in numerical calculations of laser compression, meets with some difficulties. Since the targets are small and the electron temperatures are high, the electron range in cases of practical interest is frequently of the same order of magnitude as the corona dimensions and the usual representation of a heat flux by a product of the temperature gradient and the thermal conductivity becomes invalid. We can show that in the case of electron heat conduction a deviation from the Fourier law may appear even before the electron range becomes of the same order as the characteristic scale length of the problem. In fact, a heat flux between two points at different temperatures should not exceed the flow of energy into vacuum corresponding to the higher of the two temperatures $q_{\max} \approx n_e \bar{v}_e kT_e$, where $\bar{v}_e$ is the average thermal velocity of electrons. We can easily show that the flux given by the Fourier law is comparable with the vacuum flux exactly when the electron range becomes of the same order as the characteristic scale length of the problem.

Hence, we can use a simple scheme for describing heat conduction at high temperature gradients and this is the scheme which is frequently employed in numerical calculations: the heat flux is assumed to be proportional to the temperature gradients when these gradients are low and equal to the energy flux into vacuum if the gradients are high. An interpolation formula is sometimes convenient in the description of the intermediate situation. For example, a formula of the following kind has been used:

$$q^{-1} = q_c^{-1} + q_m^{-1},$$
$$q_c = -\chi(T_e)\nabla T_e, \quad q_m = \alpha n_e (kT_e)^{3/2} m_e^{-1/2}.$$

It follows from the above discussion that the dimensionless parameter α in the last expression should be of the order of unity. However, this discussion ignores an important feature of the transfer of energy by electrons, which can be described as follows. The loss of fast electrons from the hotter regions creates an excess positive charge which produces a compensating current of cold electrons flowing toward the hotter regions. This situation is encountered in many other plasma kinetics problems. It is known that this compensating current may give rise to an instability which is accompanied by the generation of plasma waves. The interaction of electrons with a fluctuating electric field of these waves may, in its turn, reduce the heat flow.

The thermal conductivity of a rarefied plasma subjected to a high temperature gradient is calculated in[70]. It is shown that in the case of temperature gradients corresponding to fluxes $q_c > q_m$ with $\alpha \approx 0.1$, the plasma becomes unstable when ionic-sound oscillations are excited. Thus, deviations from the classical Fourier law should occur at heat fluxes much lower than the energy flow into vacuum.

A calculation of the electron heat flux in the case of gradients exceeding the instability threshold is reported in[71]. This calculation is concerned with the case when the characteristic scale of the temperature inhomogeneity is much greater than the electron range governed by the scattering on plasma oscillations. In a strongly isothermal situation with $T_e \gg T_i$, the heat flux is given by the formula

$$q = n_e kT_e v_e \left[\frac{\nu_{ei}}{\omega_{pe}} \left(\frac{MT_e}{mT_i} \right)^{1/2} \right]^{1/5} \approx \sqrt{\frac{m}{M}} n_e v_e kT_e.$$

Thus, when the temperature gradient is high, the heat flux is independent of temperature and it is $\sqrt{m/M}$ times smaller than the flow of energy into vacuum. The plasma corona temperature calculated allowing for this factor is[71]

$$kT_e \approx mc^2 \left(\frac{\dot{Q}}{n_e m_e c^3} \right)^{2/3}.$$

There is direct experimental evidence that this situation does indeed occur when laser radiation interacts with a plasma. The x-ray emission spectra, and some other characteristics of a laser plasma found experimentally, are compared in[72] with the corresponding parameters calculated using the one-dimensional hydrodynamic model. The thermal conductivity is found allowing for limitations imposed on the flux and the maximum flux is taken to be q_m, whereas the parameter α is varied to obtain the best agreement between the calculations and experimental results. The value $\alpha \approx 0.03$ is obtained in[72] and this value is in good agreement with the calculations reported in[71].

However, it should be pointed out that, in principle, the experimental results obtained in[72] can be explained in a different way. In particular, a reduction in the thermal conductivity may be attributed to the generation of spontaneous magnetic fields by an expanding plasma. Such fields were observed for the first time by Korobkin and Serov[73] for relatively low laser radiation intensities. Later, a more detailed experimental study was made of the generation of magnetic fields[74,75] and a mechanism was suggested explaining their appearance. It was found that in laser experiments the main mechanism was the appearance of a current due to the nonparallel orientation of the density and temperature gradients in a plasma. An estimate of the magnetic induction gave

$$B \approx \frac{c\tau k}{en} [\nabla T \times \nabla n]. \tag{9}$$

If the characteristic scale was taken to be the scale of the main motion of the plasma, $h \approx 10^{-2}$ cm, it was found that for intensities of $\dot{Q} \approx 10^{15}$ W/cm^2, the fields should be of the order of 1 MG. These fields were in agreement with those found experimentally.[74-76] However, the spatial distributions of the fields have not been determined. In the majority of the calculations it is assumed that the scale of the field is the same as the main hydrodynamic scale. However, it is possible that regions of smaller scale with different directions of the field may appear in an expanding plasma.

Magnetic fields in the megagauss range have only a slight influence on the hydrodynamic motion. However, these fields alter considerably the thermal conductivity of a plasma. This effect is allowed for in[77]. It is shown that the x-ray emission spectrum calculated using a two-dimensional hydrodynamic model allowing for the generation of magnetic fields is in agreement with the experimental results without assumption that the maximum heat flux is weak.

Clearly, both the above factors play some role in the reduction of the thermal conductivity of a plasma. The difference between the one- and two-dimensional calculations dealing with the same process may indicate that some effective values of the transport coefficients should be used in the one-dimensional case.

8. NUMERICAL CALCULATIONS RELATING TO SHELL TARGETS

Quantitative information on the processes in laser targets can be obtained and laser fusion experiments can be optimized only on the basis of detailed numerical calculations based on models which are as realistic as possible. It is difficult to construct such models because of the extreme complexity of the investigated phenomena which makes it necessary to carry out lengthy numerical calculations. For example, considerable difficulties in numerical simulation are encountered in studies of the dynamics of compression in real three-dimensional formulation and this applies also to correct allowance for turbulent mixing in the motion of multilayer shells, to kinetic description of the evolution of the corona and heat-conduction zone where the particle range may be of the same order as the characteristic spatial scale, and to correct description of the plasma turbulence in the light-absorption zone. All the calculations of compression carried out so far have been based on hydrodynamic models. The most complex of these models allow for two-dimensional transient motion of a many-component quasineutral plasma transfer of energy by heat conduction, radiation, and fast particles, absorption and scattering of laser radiation, thermonuclear energy yield, generation of spontaneous magnetic fields, and viscous dissipation. However, it is more usual to employ a one-dimensional model and to ignore many of these processes. Calculations carried out using simplified models usually overestimate the degree of compression and thermonuclear yield.[41]

We shall now consider some of the results obtained by numerical calculations. We shall begin with calculations of the compression dynamics. Figure 2 shows the results of a one-dimensional hydrodynamic calcula-

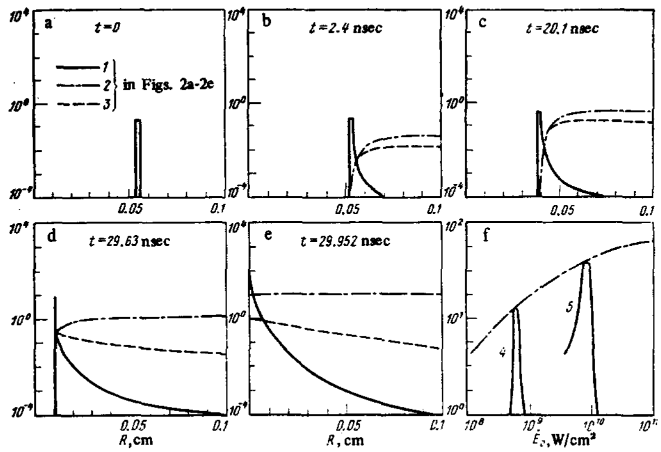

FIG. 2. Density and temperature distributions at various moments during compression of a thin solid DT shell. The results are plotted for a target of 7.5 μg mass irradiated with laser pulses of 5.3 kJ energy. The curves in Figs. 2a–2e give the distributions $\rho(R)$ (1), $T_e(R)$ (2), and $T_i(R)$ (3). The curves in Fig. 2f give the dependences of the energy gain G on the initial power $\dot{E}_0$ carried by a pulse of optimal shape: 4) $E_L=5.3$ kJ, $M=7.5$ μg; 5) $E_L=43$ kJ, $M=60$ μg.

tion of the distributions of the density and temperature at various moments during compression of a simple spherical shell.[78] The results apply to a DT shell with an initial radius of 0.54 mm and a mass of 7.5 μg irradiated with a suitably shaped laser pulse of 5.3 kJ energy. It follows from these calculations that the thermonuclear yields depends very strongly on the rate of supply of the energy to the plasma. For example, for the case shown in Fig. 2, the maximum energy gain close to 10 is obtained when the initial laser pulse power is $\dot{E}_0=6\times10^8$ W. When this pulse power is reduced, the gain for a target of the above mass falls steeply even in the case of the optimal shape of a laser pulse which is $\dot{E}\approx\dot{E}_0[1-(t/t_0)]^{-2}$. Thus, in the case of simple shells the thermonuclear yield is very sensitive to the compression regime.

Considerable advantages result from the use of complex multilayer targets, which have been the subject of intensive investigations in recent years. The main components of a complex target are a shell formed by the thermonuclear fuel (DT), a heavy shell of a high-Z material (for example, Au or U), and an ablation layer made of a material with a comparatively low value of Z (Be, polyethylene).

The use of a heavy high-Z shell ensures primarily screening of the thermonuclear fuel from heating as a result of electron heat conduction and fast particles, so that the compression regime can be made nearly adiabatic. It is very important from the practical point of view that this also leads to a considerable relaxation in respect of the shape of a laser pulse. The use of thin ($\Delta R/R\ll1$) heavy shells result in an accumulation of the kinetic energy for a fairly long time so that longer laser pulses of relatively low intensity can be used.[79] This also helps in the solution of various technical problems encountered in experiments. Other advantages can be obtained by the use of a complex system of shells of different densities, as suggested in[24].

The results of calculations and optimization of shell targets are reported in[38,78–82]. The distribution of the density during compression of a three-layer shell is shown in Fig. 3. (This figure gives the results of calculations reported in[38] for a shell whose initial radius is 0.5 mm and whose mass is 10 μg; it is assumed that the laser pulse has a Gaussian profile, 100 kJ energy, and half-width of 3.5 nsec.) We can see that, in spite of far from optimal shape of the laser pulse, the use of a heavy shell makes it possible to reach high degrees of compression. We shall now consider in greater detail the energetics of complex targets. In thermonuclear fusion under inertial containment conditions the analog of the Lawson criterion $n\tau$ is the product ρR or, more exactly, the integral $\mu'=\int_0^R\rho dR$, which governs the degree of burn-up of the DT mixture, loss of the α particles, and energy gain. An analysis of the numerical results shows that in the case of homogeneous targets the degree of burn-up is $\varphi\approx[(6/\mu')+1]^{-1}$, and the thermonuclear yield is given by the simple formula $E_{TN}\approx330M\varphi$ (here, the energy is in kilojoules and M is the mass of the DT mixture in micrograms). A heavy shell has a number of advantages, compared with a homogeneous target, not only during the compression but also during the reaction stage. The main effect is that the part of the shell remaining at the reaction ignition moment slows down (because of its inertia) the expansion of the fuel and helps to increase the degree of burn-up. The thermonuclear burn-up of targets with heavy shells is considered in detail in[80], where it is shown that in the case of sufficiently thin shells the degree of burn-up can be estimated employing a simple formula obtained for homogeneous targets except that μ' should be regarded as the sum of the integrals taken over the regions occupied by the fuel and the shell. For examples shown in Fig. 3 the integral with the shell gives $\mu_s=1.5$ g/cm², whereas the fuel, which is of lower density, gives $\mu_F\approx0.8$ g/cm². The degree of burn-up is governed by the sum $\mu_s+\mu_F=2.3$ and it is approximately 0.28.

Complex shell targets have a number of advantages but they are not free of certain shortcomings. The most obvious shortcoming is that a target of this kind

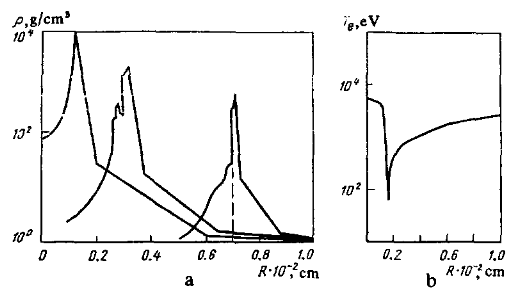

FIG. 3. (a) Distribution of the density during compression of a three-layer shell target (layer of solid DT, layer of high-Z material, and layer of low-Z material; initial target radius 0.5 mm, mass 10 μg; Gaussian laser pulse of 3.5 nsec duration and 100 kJ energy). b) Distribution of temperature at the moment of shell collapse.

consists mainly of an inert substance which does not participate in the thermonuclear reaction. The energy yield can be increased considerably by using a heavy shell made of a reactive substance (for example a suitable but still slightly futuristic material might be a DT metal). Another shortcoming of the shell targets is a strong Rayleigh–Taylor instability which appears because of the steep density gradients. Estimates indicate that this instability is a serious obstacle in the case of heavy shells whose ratio of the thickness to the radius, $\Delta R/R$, is much less than 10%. A rapid growth of short-wavelength perturbations of the accelerated moving boundary between two media of different densities results in turbulent mixing of these media. If a shell is sufficiently thin, such mixing results in its complete "dissolution" in the surrounding matter. Thus, there are fairly stringent limitations on the relative thickness of the shells which—in the final analysis—may exert a decisive influence on the maximum possible energy gain for a given type of target. The finding of the optimal target structure may be helped by various methods of stabilization of the Rayleigh–Taylor instability suggested in the literature.[81,82] Since the fabrication of microspheres for such targets presents considerable technological difficulties, it is suggested in[83] that a shell target be formed by point explosion in a very dense gas. A shock wave expending from a point explosion, initiated by a breakdown of gas with an auxiliary laser, collects almost all the matter into a shell near the front of a spherically diverging shock wave. The thickness of a shell of this kind should be of the order of $\Delta R \approx 0.1R$ and the distribution of the density in the shell should be relatively smooth, because this favors suppression of the Rayleigh–Taylor instability. Irradiation of such a shell with high-power laser pulse produces a spherically diverging optical detonation wave in the gas surrounding the shell[5] and the pressure exerted by this wave compresses the shell to very high densities and initiates thermonuclear reactions in the collapse stage. Irradiation of such a shell with a series of pulses following a predetermined time and energy program should result in the collapse, at the center, of a series of converging dense shells and this would increase even further the final density and temperature.

Numerical calculations of the compression and thermonuclear burn-up of multilayer targets of the type described in[38,78–83] are quite difficult and time-consuming. On the other hand, it is very difficult to judge the true precision of such calculations, especially in view of the strong dependence of the thermonuclear yield on the target parameters and compression regime. Therefore, the calculated values of the energy gain should be treated with caution. The most reliable calculations suggest that an energy gain of about 100 may be expected for laser energies from 10^5 to 10^6 J.

9. EXPERIMENTAL RESULTS

We have been mainly concerned with theoretical investigations of various aspects of laser fusion. The results of such investigations demonstrate the extreme difficulty of the problem but provide some grounds for optimism. In any case, theoretical results have stimulated interest in the problem and have resulted in a considerable increase in the experimental activity. We shall discuss very briefly some experimental results obtained recently. We shall mention particularly the work of Charatis *et al.*[84] who studied compression of glass shells filled with a gaseous DT mixture. They used a two-channel neodymium-glass laser with an output energy of about 200 J. The pulse shaping system made it possible to vary the duration from 30 psec to 1 nsec. The targets were spherical glass shells with an internal diameter of 30–700 μ and walls 0.5–12 μ thick; the gas pressure inside the shells was 1–100 atm. The optical system ensured a high symmetry of target irradiation. The energy balance (reflection, x-ray emission, fast particles) was determined experimentally and x-ray photographs of the target were obtained from which the size of the target during compression could be deduced. In the best experiments the neutron yield was 5×10^6 when the energy reaching the target was 60 J and the target diameter was 50 μ, thickness 0.5 μ, and pressure in the mixture 10 atm. A compression by a factor of the order of 100 was recorded.

The results reported in[84] did not show clearly whether the observed neutrons originated from the compressed target core or whether they were formed as a result of "burning" through the glass shell. One could not be sure that a 1-percent-thick shell ($\Delta R/2R \approx 10^{-2}$) did not fracture and become mixed with the fuel in the process of hundredfold compression. However, the experiments in question were of considerable interest because they were the first to be performed on shell targets in the range of laser intensities up to 10^{16} W/cm^2.

Interesting results were also obtained in preliminary experiments on spherical glass shell targets filled with a DT gas irradiated with just one laser beam. The "disk and sphere" target configuration was designed so that the upper layer of a low-Z material was evaporated by a preliminary pulse from one side of a sphere and from a plane thin substrate to which the sphere was bonded. A plasma cloud formed around the sphere was irradiated by a second pulse and having become hot, it transferred energy by electron heat conduction to the surface of the sphere ensuring ablation of its top layer and collapse of the shell. Preliminary experiments demonstrated that the shell was compressed and neutrons were generated.[85]

The first experiments involving spherical irradiation of homogeneous CD_2 targets by neodymium-laser pulses of 1–2 nsec duration[86] also gave interesting values of the total energy input and estimates of the recoil impulse as well as of possible degree of compression of the matter at the center of the target. It was concluded in[86] that under these conditions the degree of compression should be 30, but the temperature at the center of the target at the moment of maximum compression was insufficient for the generation of thermonuclear neutrons. The neutron yield recorded in[86] was due to reactions occurring in the outer parts of the corona.

The measurements reported in[84] showed that there was strong reflection and dissipation of the laser energy and the energy losses increased with the intensity.

Somewhat different results were obtained in[84] where practically complete absorption of light was observed and attributed to the excitation of a parametric instability. Although the degree of absorption was not finally established, there was no doubt about the considerable role played by the "anomalous" processes in the range of intensities of practical interest. This was indicated by the results reported in[83,84,87], where a study was made of the spectral composition of the plasma-scattered laser radiation and components of frequencies $\omega_0/2$, $3\omega_0/2$, and $2\omega_0$ were observed (ω_0 was the frequency of the incident radiation). The shift and width of the observed lines were in good agreement with the mechanism based on the parametric growth of plasma waves. Indirect evidence of the parametric instability was provided by the amplitude modulation of the reflected laser radiation, observed in many investigations (see[84,88]). A definite experimental confirmation of the appearance of turbulent instabilities in a laser plasma was obtained in[88], where modulation of the reflected laser radiation was found to be accompanied by modulation of the x-ray emission from the target plasma.

Interesting results obtained in a study of the spectra of ions in an expanding laser plasma were reported in[89]. It was found that considerable proportion of the energy (over 50%) was carried away by very fast ions when a target was irradiated with 30 psec neodymium laser pulses of intensity exceeding 10^{16} W/cm^2. Since the energy of these ions was approximately an order of magnitude higher than the average energy of particles in a plasma, the specific recoil impulse should be considerably less and, consequently, the efficiency of conversion of the laser energy into the energy of compressed matter was much reduced. The effect was of threshold nature and was observed also at $\lambda = 10.6\ \mu$ but at a lower intensity of 10^{14} W/cm^2.

An agreed explanation of this effect is not yet available but it may be due to plasma instabilities or due to weaker heat conduction in an inhomogeneous plasma, which may result from the generation of spontaneous magnetic fields.

Careful measurements of the spontaneous magnetic fields generated in a laser plasma[73] were reported in[74–76]. It was shown that spontaneous magnetic fields in the plasma of a plane target could reach 10^6 G as a result of irradiation with neodymium or CO_2 laser light.

The experimental evidence taken as a whole demonstrates that sufficiently strong absorption does take place and targets can be compressed considerably by irradiation with pulses of intensities up to 10^{16} W/cm^2. Fuller experimental information, which will make it possible to be more definite about the validity of the theoretical predictions, may be obtained using apparatus with higher energy outputs. At present lasers suitable for fusion research and capable of an output energy of 10^3–10^4 J per pulse are being constructed in several laboratories. They will provide opportunity for obtaining a considerable thermonuclear yield and for checking the correctness of the basic principles of laser fusion and possibly even reaching the physical threshold of the fusion reaction.

Neodymium-glass lasers are the most advanced among the existing laser systems. They have a suitable wavelength, $\lambda = 1.06\ \mu$, which makes it possible to shape pulses with the required duration in the range 10^{-8}–10^{-10} sec; moreover, the radiation produced by these lasers is characterized by a high directionality and it can be focused on small targets. The first laser fusion experiments employing such equipment are planned for 1977.[90,91]

Theoretical estimates give an energy gain $G \approx 100$ for laser microexplosions when the laser pulse energy is 10^5–10^6 J. Hence, it follows that practical use of laser fusion in economic power stations would require reliable lasers capable of this energy output with an efficiency of 3–10%. The current efficiency of neodymium-glass lasers is 0.1–0.3% and it will be difficult to increase it about 1%. In view of this, it would be interesting to consider high-pressure pulse CO_2 lasers. The efficiency of these lasers may reach 3–5%. Unfortunately, the techniques for controlling the parameters of the pulses emitted by these lasers is much less developed and there is a shortage of components and optical materials suitable at $\lambda = 10.6\ \mu$. Moreover, there still remains the basic problem whether the necessary compression can be produced by long-wavelength infrared radiation because in this case the instability thresholds are much lower and the problem of heat conduction from the corona to the core is more acute. Nevertheless, large-scale systems for laser fusion with an output energy of 10^4 J per pulse are being developed.[92]

Chemical pulse lasers emitting at wavelengths of 3–4 μ and utilizing the $H_2 + F_2 = 2HF^*$ (or $D_2 + F_2 = 2DF^*$) reaction are of considerable interest: in these lasers the reaction products are in the excited state and they ensure the necessary population inversion. Recent successes in the development of high-energy and high-efficiency (up to 10% of the chemical energy) chemical laser systems make them very promising for laser fusion research.[93]

The main principles can also be tested by photodissociation iodine lasers emitting at $1.315\ \mu$.[94] Experimental work planned for the nearest future will give a definite answer to the question whether laser fusion is possible. The recent progress places the laser method in the same position as other approaches to the solution of the fusion problem, which is one of the most difficult in modern physics. However, there is no doubt that laser fusion research is of considerable interest because it provides means for laboratory investigations of matter under extreme conditions of high pressure and temperatures.

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