Thus, the results described above and of similar experiments are, in our opinion, crucial for the development of a theory of nonlinear interaction between powerful electromagnetic radiation and a plasma.

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NONSTATIONARY PARAMETRIC TURBULENCE OF A PLASMA

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We have investigated the evolution of plasma turbulence initiated by parametric instability (by the decay of the pump wave into two plasmons) and saturated by induced scattering of the plasmons from the ions. We show that parametric turbulence develops in a dense plasma during a nanosecond pulse from a powerful neodymium laser, and is essentially nonstationary as a result of the oscillations of the noise about the turbulence level.

A vital problem in the theory of parametric instability of plasma is the determination of the temporal evolution of the plasma turbulence (parametric turbulence), which is the physical basis for the mechanisms of absorption and scattering of powerful radiation, and of the heating and acceleration of the plasma particles (see, e.g., the review [1]). We report here the results of a theory of nonstationary parametric turbulence excited in a uniform plasma through decay of a pump wave (with frequency ω_0 and wave vector $\vec{k}_0$) into two plasmons [2, 3] and stabilized by induced scattering by ions[1]. Near the instability threshold, the spectral energy density of the plasma noise is given by

$$W(k, \theta, t) = \sqrt{2}(2\pi/k)^2 s_0 n_e \kappa T_e y(x, \tau)\,\delta(k - k_m),$$

where the angular distribution y(x, t) satisfies the equation

$$\frac{\partial y(x,\tau)}{\partial \tau} - \frac{1}{y(x,\tau)} = y(x,\tau)\left[a^2 - x^2 - a\int_{x_1}^{x_2} dx'\, q(x'-x)\,y(x',\tau)\right]. \quad (1)$$

Here $k_m = r_{De}^{-1}\sqrt{\omega_0 - 2\omega_{Le}}/\sqrt{3\omega_{Le}}$ is the wave number with extremal value for instability development, ω_{Le} and r_{De} are the Langmuir frequency and Debye radius of electrons with density n_e and temperature T_e (κ is Boltzmann's constant), and θ is the angle between the wave vectors $\vec{k}$ and $\vec{k}_0$. This angle is connected by the relations $x = \theta - \pi/4$ $(0 < \theta < \pi/2)$ and $x = \theta - 3\pi/4$ $(\pi/2 < \theta < \pi)$ with the variable x that determines the deviation of the propagation of the short-wave plasmons $(k >> k_0)$ from the optimal pump directions $\theta = \pi/4$ and $\theta = 3\pi/4$; $\tau = 4\tilde{\gamma}t$, where $\tilde{\gamma}$ is the plasma-wave damping decrement and takes into account the Coulomb electron-ion collisions and the Landau damping by the electrons. The quantity $a^2 = (1/2)(L - E_{min}/E_0)$ characterizes the excess over threshold, i.e., the excess of the pumping-field intensity E_0 over the minimum E_{min} necessary for instability development. The parameters α and s_0 give the relative weights of the respective contributions of the nonlinear interaction of the turbulent noise and the spontaneous noise described by the second terms in the left-hand side of Eq. (1) (c is the speed of sound, v_{Ti} and

r_{Di} are the thermal velocity and Debye radius of the ions)

$$a = (32\pi^{3/2} n_e r_{D_e}^3)^{-1} \frac{v_{T_e}^2}{c v_{T_i}} \left(\frac{\omega_0 - 2\omega_{L_e}}{2\omega_{L_e}}\right)^{1/2} \left(\frac{r_{D_i}}{r_{D_e}} + \frac{r_{D_e}}{r_{D_i}}\right)^{-2};$$

$$s_0 = (6\pi n_e r_{D_e}^3)^{-1} \frac{\tilde{y}}{\omega_{L_e}} \left(\frac{\omega_0 - 2\omega_{L_e}}{6\omega_{L_e}}\right)^{1/2},$$

and the kernel is defined by the equation ($\beta = (3/2)^{3/2} v_{Te}/c v_{Ti}$)

$$q(x) = x \int_0^{\pi/2} d\phi \cos^4\phi \left(\sin^2\phi + \frac{x^2}{2}\cos^2\phi\right)^{-1/2} \exp\left[-\frac{1}{4}\beta^2 x^2 \times \times \left(\sin^2\phi + \frac{x^2}{2}\cos^2\phi\right)^{-1}\right].$$

Equation (1) was obtained in the differential-pumping approximation, when the scale of the kernel β^{-1} is much smaller than the characteristic scale of the increment a ($a\beta \gg 1$).

Figure 1 shows a numerical solution of Eq. (1) for spectral energy density of the plasma noise $y(x, \tau)$ at the instant of time $\tau = 50$. It is seen from the figure that nonstationary noise is strongly "cut up" as a function of x, and differs qualitatively from the stationary noise $y(x, \infty)$ obtained earlier in [2]. The time dependence of either the spectral density of the noise $y(x, \tau)$ or the total noise integrated over the angles

$$\bar{y}(\tau) = \int_{x_1}^{x_2} dx\, y(x, \tau)$$

is also a nonmonotonic and oscillating function (Fig. 2). The plasma noise, as follows from (1), tends to a stationary distribution [2] in accordance with the relation ($C_1 \sim 1$; $a^2\tau \gg 1$):

$$\bar{y}(\tau) = \frac{9}{4}\frac{a^4}{\epsilon}\left\{1 + \frac{C_1}{a^2\tau}\exp\left(-\frac{9}{8}\tau\frac{\epsilon^2}{a^6}\right)\sin(0{,}1 a^3 \beta\tau)\right\} \quad (2)$$

$$\epsilon = \frac{\pi^{3/2}}{8} a \beta^{-3}$$

The characteristic damping time of the oscillations of $\bar{y}(\tau)$ is $\tau_\infty \simeq a^6/\varepsilon^2$, and is larger the higher the excess above threshold a^2. This time, $t_\infty = \tau_\infty/4\tilde{y}$, is so long that in the case of pulsed pumping the turbulence is as a rule essentially nonstationary during the lifetime of the plasma. For example, in a hydrogen plasma with density $n_e \simeq 2.5 \times 10^{20}$ cm^{-3} and an electron or ion temperature $\sim$ 1 keV, at a relative detuning of the neodymium-laser frequency $(\omega_0 - 2\omega_{Le})/(2\omega_{Le}) \simeq 0.1$ and at 10% above the minimal threshold field, the time required for the noise to reach its stationary state, $t_\infty = 10^{-3}$ sec, is longer by 5 – 6 orders of magnitude than the duration 10^{-8} – 10^{-9} sec of a typical nanosecond pulse. The noise reaches in this case the turbulent value after a time $\tau_T \simeq a^{-2}\ln(a^4/\varepsilon) \ll \tau_\infty$.

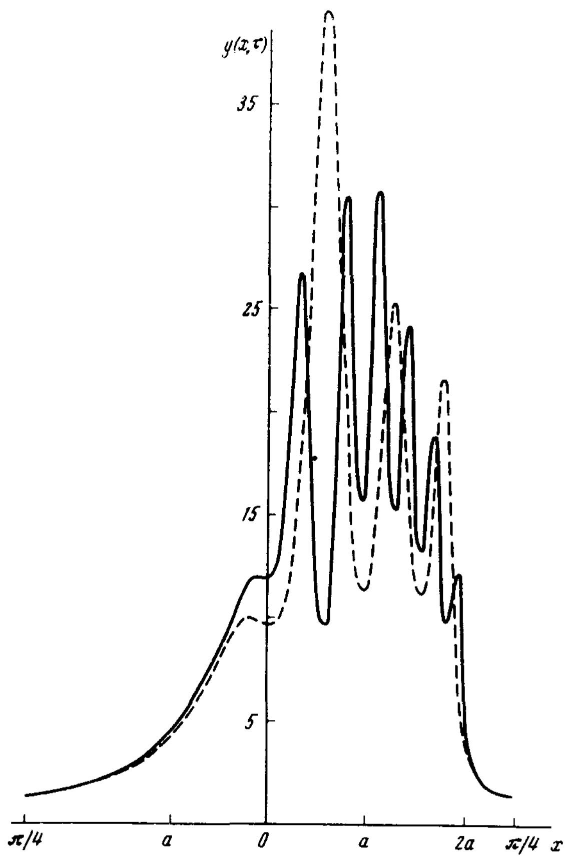

Fig. 1. Angular distribution $y(x, \tau)$ at the instant of time $\tau = 50$ for $\beta = 10$ (dashed curve) and $\beta = 20$ (c9ntinuous curve) at $a = 0.316$, $\varepsilon = 3 \times 10^{-3}$, and $x_{1,2} = \mp\pi/4$.

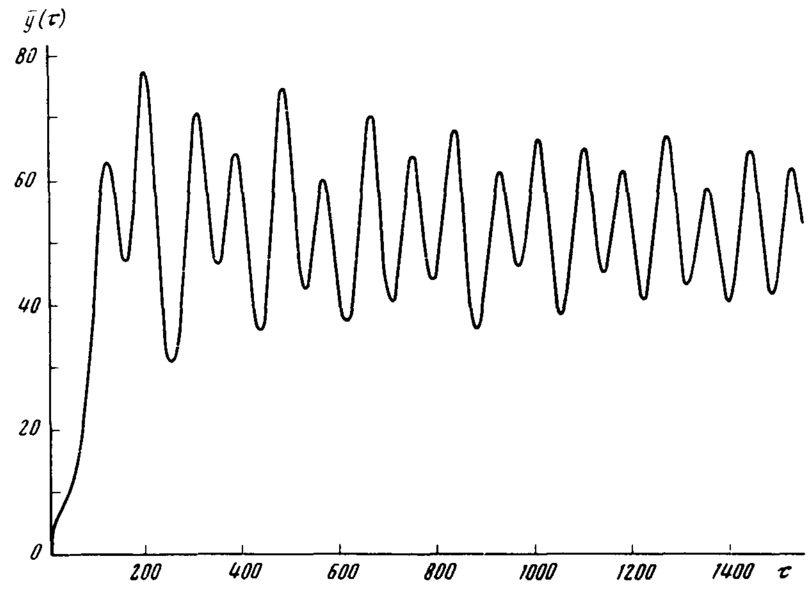

Fig. 2

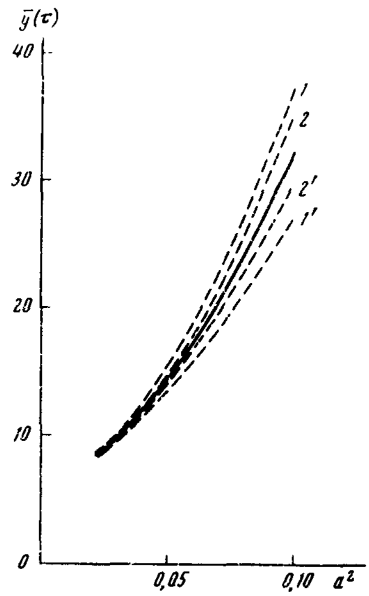

Fig. 3

Fig. 2. Total noise $\overline{y}(\tau)$ at $a = 0.2$, $\varepsilon = 10^{-4}$, and $\beta = 10$ in the angle interval $-\pi/4 \leq x \leq \pi/4$.

Fig. 3. Total noise $\overline{y}(\tau)$ as a function of the excess above threshold a^2 at $\varepsilon = 10^{-3}$, $\beta = 10$, and $x_{1,2} = \mp\pi/4$. The dashed curves 1, 1' and 2, 2' are the limits of variation of $\overline{y}(\tau)$ for the two time intervals $\tau > 50$ and $\tau > 300$, respectively. The central (solid) curve represents the stationary noise.

In the described plasma variant we have $\tau_T \simeq 10^{-10}$ sec, so that at the indicated times of laser-pulse duration the turbulence is already strongly developed. As seen from Fig. 2, at $\tau_T < \tau < \tau_\infty$ the values of the total noise $\overline{y}(\tau)$ can differ significantly from the stationary value $\overline{y}(\infty)$. The limits of variation of the total noise for two different time intervals, $\tau > 50$ and $\tau > 300$, are illustrated in Fig. 3 for different values of a^2.

The time evolution of the parametric turbulence saturated by the spectral redistribution consists of three stages. In the first, $0 \leq \tau \lesssim \gamma^{-1} \simeq a^{-2}$, the plasma noise grows exponentially. During the second stage, $\gamma^{-1} < \tau < \tau_T$, the noise saturates; then at $\tau_T < \tau < \tau_\infty$ the noise oscillates for a long time at a frequency $\sim 0.4a^3\beta\gamma$ (see (2) and Fig. 2), and assumes a stationary value at $\tau > \tau_\infty$. The third stage of noise evolution is characterized by the largest degree of choppiness of the spectrum as a function of x. The number of peaks into which the distribution $y(x, \tau)$ breaks up is determined by the relation between the increment scale a and the nuclear scale β^{-1}, and turns out to be approximately equal to $a\beta$. Figure 1 shows the increase in the number of peaks on the spectral distribution with increasing β. The distance between the peaks and the peak widths are of the order of β^{-1}.

It should be noted that the oscillating approach to the stationary value is typical of saturation mechanisms that take the waves out of the buildup region. For stabilization mechanisms such as nonlinear frequency shift, which lead to an effective increase of the instability threshold, the stationary value is assumed monotonically and much more rapidly, within a time on the order of the maximum reciprocal increment in the logarithm of the ratio of the stationary turbulent noise to the spontaneous noise.

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[1] For more details on nonlinear mechanisms see, e.g., Secs. 9 and 13 in the review [4].

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CW ARGON LASER WITH 0.5 kW OUTPUT POWER

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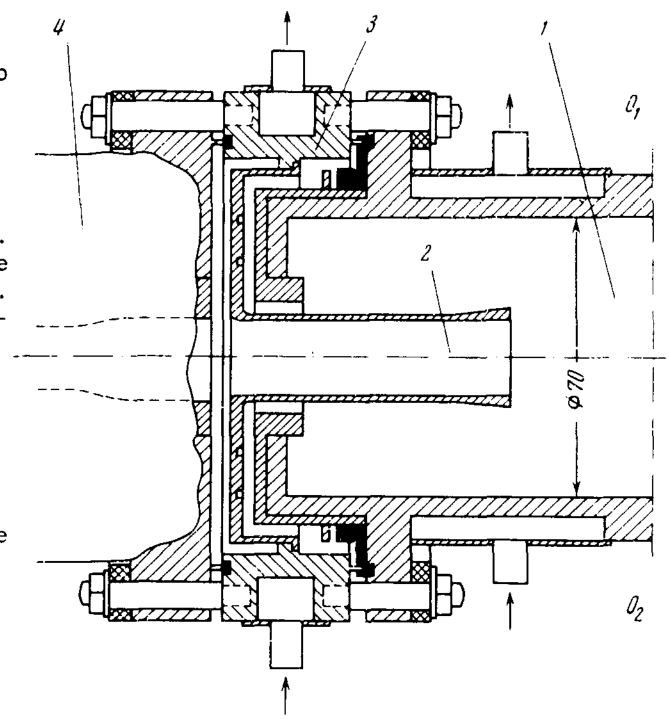
1. We report here the realization of continuous visible coherent emission of 500 W power from a laser using Ar-II ions. This is an unprecedented output power for the short-wave part of the electromagnetic spectrum. In addition to the large output power, the laser described below exceeds all the existing high-power ionic gas lasers in output power per unit discharge length, in efficiency, and in operating time. This combination of unique characteristics of the Ar-II laser was made possible by solving a number of fundamental problems of obtaining powerful stationary discharges at reduced pressure and investigating the saturation of the power of ionic lasers as a function of the pumping current.

2. Inverted population of the 4p - 4s electronic transitions of singly-ionized argon was produced in a plasma of a high-power dc discharge. The laser was constructed in the form of two large-current discharge tubes with a common cathode between them. The tubes were made up of individual aluminum sections, each coated with an Al_2O_3 film to increase the endurance to ion bombardment in the discharge and to improve the heat dissipation [1, 2, 5]. The active part of the discharge, with 1.6 cm diameter in each tube, was 1.5 m long, the individual sections were 2.5 cm long, and heat-resistant rubber was used as the vacuum sealing material. The argon was continuously fed into the cathode region and pumped out slowly at the anode ends. The construction of the dismountable cold cathode is shown in the figure. The operating principle is based on retaining the cathode spots of the arc inside the cathode cavity by means of a self-heating refractory bushing [1, 3]. The use of this principle has made it possible to construct a cathode capable of delivering a stationary current of approximately 1 koloampere in a low-pressure discharge and keep the evaporation from the cathode from affecting the plasma in the active part of the discharge. To reduce the cathode voltage drop and increase the stability of the cathode, metallic bismuth was added to the working cavity. The optical resonator was 4 m long and made up of two internal multilayer mirrors produced by cathode sputtering. The transmission of one mirror was 8%, that of the other 2%, and the curvature radius of both mirrors was $R \sim 10$ m.



Arc cathode (the construction is symmetric about the O_1O_1 axis): 1 – working cavity of copper, 2 – heating bushing of molybdenum, 3 – bushing holder, 4 – discharge tube. The teflon vacuum seal is shown black; the arrows show the direction of the cooling water.

3. The laser generation was produced under conditions of saturated discharge