We point out that, at identical pressures and pumps, the output energy of the anti-Stokes component is smaller for longer chambers. This fact, as well as the presence of the "dip," contradicts the conclusions of the theory that does not take into account the spatially-limited phase capture, and is the direct consequence of the latter (see item b above).

Thus, the statements made at the beginning of the article concerning the mechanism and singularities of the axial ASRS can be regarded as valid. We note in conclusion that spatially-limited phase locking should play an important role also in other parametric interactions between fields, in which real transitions in matter take part.

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GENERATION REGIMES OF SOLID-STATE RING LASER

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The overwhelming majority of investigations of lasers with ring resonators were performed with a gaseous active medium having an inhomogeneously broadened luminescence line. Yet undisputed interest attaches to the study of the dynamics of generation of ring lasers based on active crystals with homogeneously broadened luminescence line. In such lasers, there is a strong competition between the opposing waves [1, 2]. We have investigated theoretically and experimentally the generation regimes in a solid-state ring laser.

1. The dynamics of ring-laser generation is described on the basis of the following system of equations for the complex amplitudes of the opposing waves $\tilde{E}_{1,2}$ and the inversion density N:

$$\dot{\tilde{E}}_{1,2} = -\frac{\omega}{2Q}\tilde{E}_{1,2} + \frac{i\tilde{m}_{1,2}}{2}\tilde{E}_{2,1} \mp i\frac{\Omega}{2}\tilde{E}_{1,2} + \frac{\sigma}{2T}\left(\tilde{E}_{1,2}\int_0^{\ell} N dx + \tilde{E}_{2,1}\int_0^{\ell} N e^{\pm i2kx}\,dx\right),$$

$$\dot{N} = W - \frac{N}{T_1} - \frac{a}{T_1}N\left|\tilde{E}_1 e^{-ikx} + \tilde{E}_2 e^{ikx}\right|^2. \qquad (1)$$

Here $\tilde{m}_{1,2} = m_{1,2}\exp(\pm i\theta_{1,2})$ are the complex coefficients of coupling via back scattering ($m_{1,2}$ and $\theta_{1,2}$ are the moduli and phases of the coupling coefficients), ω/Q is the resonator bandwidth, $T = L/c$ is the time necessary for the light to travel around the resonator, σ is the transition cross section, W is the pumping rate, T_1 is the longitudinal relaxation time, ℓ is the length of the active element, $k = 2\pi/\lambda$ is the wave number, $\Omega = \omega_1 - \omega_2$ is the difference between the natural frequencies of the resonator for the opposing waves, and $a = \sigma c T_1/8\pi\hbar\omega$. It is assumed that the generation regime is single-mode, with a frequency close to the center of the luminescence line.

2. In a ring laser it is possible to have mutual synchronization of the opposing waves (the amplitudes of the waves $E_{1,2}$ and the phase difference $\Phi = \phi_1 - \phi_2$ are constant in time). The synchronization regime can take place when the detuning of the resonator natural frequencies lies inside the synchronization band $|\Omega| < \Omega_0$. Inside the synchronization region, the intensities and phase differences of the opposing waves are functions of the detuning Ω. If the moduli of the coupling coefficients are equal ($m_1 = m_2 = m$), the condition for the stability of the regime in which the opposing waves are synchronized is

$$m\left|\sin\frac{\theta_1-\theta_2}{2}\right| > \frac{\omega}{Q}\left[1-\frac{\sqrt{1+8(1+\eta)}-1}{2(1+\eta)}\right]. \qquad (2)$$

It follows therefore that in the case of equal phases of the coupling coefficients ($\theta_1 = \theta_2$) the standing-wave regime is unstable at all m. In the case $\theta_1 \neq \theta_2$, the standing-wave regime can become stable if the coupling is strong enough. It is easily seen that at large values of η the standing-wave regime can become unstable.

The synchronization bandwidth Ω_0 is given by

$$\Omega_0 = \sqrt{m^2\sin^2\frac{\theta_1-\theta_2}{2} - \frac{1}{4}\left(\frac{\omega}{Q}\eta - m\left|\sin\frac{\theta_1-\theta_2}{2}\right|\right)^2}. \qquad (3)$$

We see therefore that for a specified modulus of the coupling coefficient Ω_0 is largest at $|\theta_1 - \theta_2| = \pi$. The width of the synchronization band depends on the pump, and decreases with increasing η.

3. In a ring laser it is also possible to have a stationary generation regime with essentially unequal wave amplitudes (the regime of unidirectional generation). This regime was investigated by us under the assumption that $\Omega = 0$ and the coupling coefficients are complex-conjugate ($m_1 = m_2 = m$, $\theta_1 = \theta_2$). The regime of unidirectional generation can exist under the condition $m \ll \omega\eta/Q(1 + \eta)$ and stable if $m < (1/2)\sqrt{(\omega/Q)\eta(1 + \eta)/(2 + \eta)T_1}$.

It follows from the results that in the case when the coupling coefficients are close to complex-conjugate, and $m > (1/2)\sqrt{(\omega/Q)\eta(1 + \eta)/(2 + \eta)T_1}$ the regime of unidirectional generation and the regime of

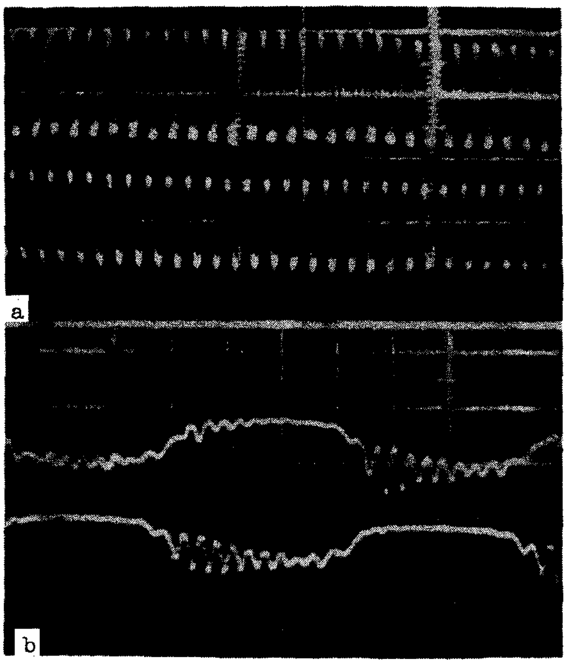

Fig. 1

opposing-wave synchronization become unstable, i.e., generation with stationary amplitudes and a phase difference between the opposing waves is impossible. In this case an automodulation regime should set in (the amplitudes and phase differences vary periodically with time).

4. The generation regimes of a YAG:Nd^{3+} ring laser were investigated experimentally. The laser operated continuously at a wavelength λ = 1.06 μ. The end faces of the active element were made non-reflecting (the residual reflection coefficient from the end face did not exceed 0.4%).

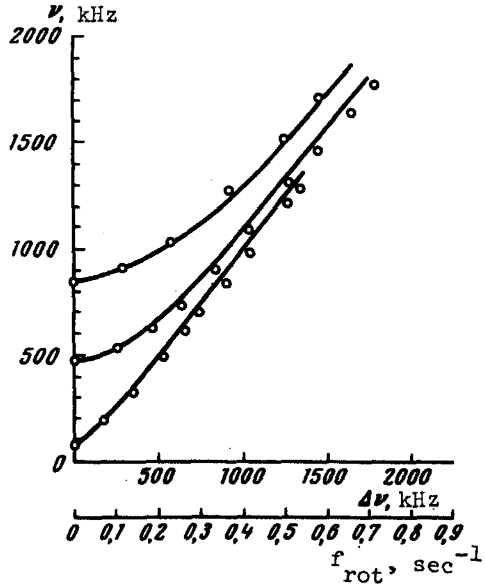

Fig. 2

It was shown that the generation regimes depend significantly on the tuning of the ring resonator and on the speed of rotation (the laser was mounted on a rotating platform). When the laser was at rest, depending on the tuning of the resonator (e.g., when the position of one of the mirrors was altered), two regimes were observed: Synchronization of two opposing waves and automodulation of the intensities of the opposing waves. Inside the synchronization region, the difference between the intensities of the opposing waves was a function of the speed of rotation and increased when the synchronization region was approached. In the laser at rest, the intensity difference depended on the tuning of the resonator. The tuning affected also the extent of the synchronization region, and at definite tunings it was impossible to go out of the synchronization region (the maximum difference of the natural frequencies of the rotating laser reached 1 MHz). It was observed that the regime of synchronization of the opposing waves can become unstable when the pump excess over threshold exceed a certain value that depends on the resonator tuning. All these results are in qualitative agreement with the theoretical analysis presented above.

Typical oscillograms of the automodulation regimes are shown in Fig. 1 (200 μsec/div). It follows from Fig. 1 that the automodulation of the intensities of the opposing waves can be periodic (Fig. 1a) or can have a more complicated character (Fig. 1b - the automodulation oscillations are modulated at a lower frequency). The frequencies of the automodulation oscillations depend significantly on the resonator tuning and range from several kHz to 1 MHz. The self-oscillations of the intensities of the opposing waves are close to being in phase opposition. The dependence of the frequency of the automodulation oscillations ω_m (at a fixed tuning) on the speed of rotation of the ring laser is described sufficiently well by the formula $\omega_m(\Omega) = [\omega_m^2(0) + \Omega^2]^{1/2}$. The experimental data and the values of ω_m calculated from this formula are shown in Fig. 2.

The automodulation frequency $\omega_m(0)$ in the laser at rest is determined by the value of the coupling [2]. In the case of complex-conjugate coupling coefficients we have $\omega_m(0) = m$. Representing the coupling coefficient in the form $m = \sqrt{R/T}$ (R is the intensity ratio of the back-scattered and incident waves) and assuming $T = 0.3 \times 10^{-8}$ sec and $\omega_m/2\pi \simeq 10^6$ Hz we obtain $R \simeq 4 \times 10^{-4}$.

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TRANSFORMATION OF HIGH-FREQUENCY WAVES AND VANISHING OF LOW-FREQUENCY INSTABILITIES IN A RADIALLY-INHOMOGENEOUS BEAM-PLASMA DISCHARGE

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The influence of transformation of high-frequency waves on the operating regime of a plasma-beam discharge is analyzed. Experiments show that by varying the radial distribution of the plasma density it is possible to ensure effective radiation of the oscillations near the upper hybrid resonance and by the same token prevent transfer of the energy into the low-frequency oscillations. The latter quench the low-frequency instabilities in a plasma-beam discharge.

In experimental investigations of the linear transformation in an inhomogeneous plasma, principal attention is usually paid to plasma heating upon transformation of the transverse waves into longitudinal ones, or to radiation of the energy from the system in the oppo ite case [1 - 3]. Yet the wave transformation can exert a more diverse action on the system and, in particular, influence its stability (the possibility of using wave transformation to increase the stability of the system was first pointed out in [4]). The influence of the wave transformation on the system stability can be conveniently traced with a qualitative example, in which there are two "modes" K_1 and K_2 localized in the same region, the common region of localization of the two modes being separated by opacity barriers from the vacuum-connected transparency region of one of the modes (then part of the energy can radiate out from the system to the vacuum). Assume that the mode K_2 can be excited in the nonequilibrium system. In a homogeneous medium, the modes K_1 and K_2 are not connected with each other, and therefore the mode K_1 does not influence the instability development. In an inhomogeneous medium there can exist points where the wave vectors of both modes coincide (regions of mode "intersection") and in this case the transformation of one mode into another is frequently close to 100% in the "intersection" region [4]. Then only a "superposition mode" $K_1 + K_2$ can be excited in the medium, although far from the "intersection" region it is convenient, as before, to designate the "components" of the new mode by K_1 and K_2 respectively (see, e.g., [1, 4]). Let L be the distance between the "intersection" points of the modes, and let the dimension L of the instability region be smaller than L. In this case the condition of instability following excitation of the "superposition" mode can be written in the form

$$\alpha \exp\left(\frac{\gamma_2' L'}{V_{g2}} - \frac{\gamma_2 L}{V_{g2}} - \frac{\gamma_1 L}{V_{g1}}\right) > 1 \qquad (1)$$

where γ_2' is the growth increment of mode K_2 in the region L'; γ_2, γ_1, V_{g2}, and V_{g1} are respectively the increments and group velocities of modes K_2 and K_1; α^{-1} is the factor by which the amplitude of the excited wave is reduced because of radiation of energy to the outside ($\alpha < 1$; γ_2', γ_1, $\gamma_2 < \omega$).

The situation described above is typical, e.g., of the case of interest to us, the interaction of an electron beam with high-frequency oscillations of a radially-inhomogeneous plasma (K_1 is a mode of the Bernstein type, propagating at an angle to the magnetic field H_0, K_2 is a "cold" mode with a singularity in