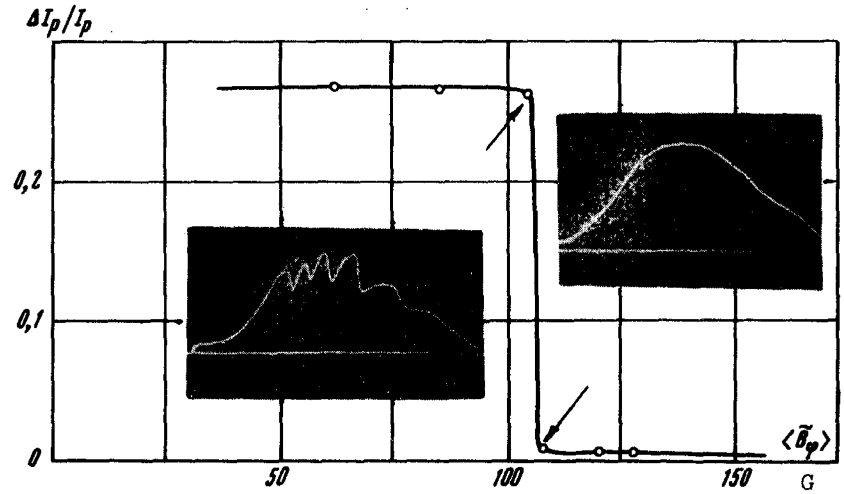
Fig. 3. Relative current-oscillation amplitude $\Delta I_p/I_p$ vs. the stabilizing HF field intensity. I_p = 3 kA, B_z = 6 kG.



The effective stability margin q_{eff} ($q_{eff} = 2\pi/(i_0 + i)$, where i_0 is the twist angle of the force line of the stellarator magnetic field on the filament boundary and i is the twist angle produced by the quasiconstant current) in the stabilized regime is $q_{eff} \approx 1$. In the stabilized regime, the pressure of the HF field, as well as the pressure of the field current, was much lower than the plasma pressure, i.e., $\tilde{\beta}_\phi \approx \bar{\beta}_\phi >> 1$, thus indicating effective heating of the plasma by the current flowing through it.

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SPATIALLY-BOUNDED PHASE CAPTURE AND AXIAL ANTI-STOKES RADIATION IN SRS IN GASES

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It is shown theoretically and experimentally that the onset in all the singularities of the behavior of the anti-Stokes component of stimulated Raman scattering propagating along the axis of the pump beam are due to phase locking of the interacting waves, which occurs in a bounded segment of the interaction path.

1. We present here the results of experimental and theoretical investigations of the anti-Stokes component of SRS (ASRS), propagating along the axis of the pump beam. These investigations were undertaken for the purpose of explaining the mechanism whereby axial ASRS is produced.

The first to observe axial ASRS was Garmire [1]. Observation of such radiation was subsequently reported in a number of publications. The following hypotheses were advanced concerning the mechanism whereby the axial ASRS is produced: 1) it is due to self-focusing [1]; 2) it constitutes Raman scattering by excited molecules [2]; 3) it is connected with lasing in the resonator made up of the parallel end faces of the cell [1, 3].

In [3, 4], devoted to the investigation of ASRS in liquids, and compressed H_2, respectively, it was shown that none of these mechanisms take place. The authors of [3] have concluded that axial and conical ASRS differ in nature. This conclusion was apparently the result of the disparity between the experimental facts established in [3] and the existing theory of ASRS [5, 6]. Further investigations did not help determine the nature of the axial ASRS.

2. It will be shown below that to explain axial ASRS it is not necessary to involve mechanisms other than the occurrence of conical ASRS. All the singularities of ASRS can be attributed to phase capture of the interacting waves, which takes place in a limited section along the beam axis. The spatial limitation of the phase-locking region is due to the reaction of the scattered-light components on the pump intensity.

The equations that take such an influence into account can be obtained by the method used in [7]. If $\vec{k}_p$, $\vec{k}_s$, and $\vec{k}_a$ are collinear ($E_j = \mathcal{E}_j \exp[i(\omega_j t - k_j z)] + c.c.$; $j = p$, s, and a pertain respectively to the pump, Stokes, and anti-Stokes components), these equations take the form[1)]

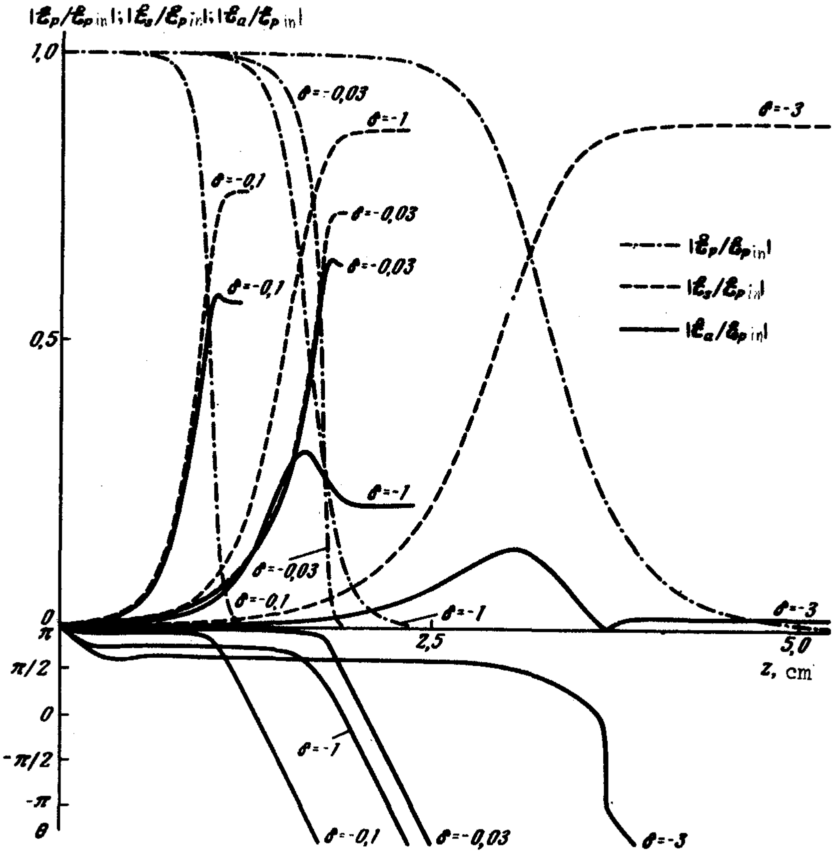

Fig. 1. Distribution of $\mathcal{E}_a$, $\mathcal{E}_s$, $\mathcal{E}_p$, and θ along the propagation direction (H_2 pressure 75 atm). The point z = 0 corresponds to the chamber cross section in which S^2, rising from the noise level at the input, reaches $10^{-6}\,\mathcal{P}_{in}^2$.

[1)]In accordance with the experimental conditions, we assume the pump pulse to be quasistationary, and neglect the saturation of the populations.

$$\frac{dS}{dz} = q_s \mathcal{P}^2 (S + A\cos\theta) \quad (1); \qquad \frac{dA}{dz} = -q_a \mathcal{P}^2 (A + S\cos\theta); \quad (2);$$

$$\frac{d\mathcal{P}}{dz} = q_p \mathcal{P}(A^2 - S^2) \quad (3); \qquad \frac{d\theta}{dz} = \mathcal{P}^2\left(q_a \frac{S}{A} - q_s \frac{A}{S}\right)\sin\theta + \Delta \, (4);$$

where

$$\mathcal{P}, A = |\mathcal{E}_{p,a}|; \quad S = \frac{\langle r_s^2 \rangle}{\langle r_s r_a \rangle}|\mathcal{E}_s|; \quad \theta = \arg(\mathcal{E}_p^2/\mathcal{E}_s \mathcal{E}_a) + \Delta z;$$

$$\Delta = 2k_p - k_s - k_a; \quad q_{a,s} = \frac{2\pi N \omega_{a,s} T \langle r^2_{a,s} \rangle}{\hbar^3 \epsilon_{a,s} v_{a,s}}; \quad q_p = \frac{q_a \omega_p \epsilon_a v_a}{\omega_a \epsilon_p v_p};$$

$$r^2_{a,s} = \left\langle \left[\sum_n \left(\frac{d_{1n}^{(a,p)} d_{n2}^{(p,s)}}{\omega_{n2} - \omega_{p,s}} + \frac{d_{1n}^{(p,s)} d_{n2}^{(a,p)}}{\omega_{n2} + \omega_{a,p}} \right) \right]^2 \right\rangle;$$

here $d_{mn}^{(j)}$ are the projections, on the direction of the j-th field, of the dipole moments of the transitions that determine the cross section of the SRS on the 1 - 2 transition; ω_{mn} are the frequencies of these transitions: T is the reciprocal line width of the 1 - 2 transition; N is the particle-number density; v_j is the velocity of the j-th wave in the medium; the symbol < > denotes averaging over the molecule orientations.

A qualitative investigation of Eqs. (1) - (4) shows that at the start of the interaction we have $\pi/2 < \theta < 3\pi/2$ ($\theta \simeq \pi$ in the case of high pump power); the Stokes and anti-Stokes components increase, since parametric transfer of the energy from the Stokes and pump components to the anti-Stokes component ($S + \mathcal{P} \to A$; see the second term in (2)) exceeds the pure Raman transformation of the anti-Stokes component into the pump component ($A \to \mathcal{P}$). Further, the change of the phase θ (see (4)) as a result of the increase of A/S and of the decrease of $\mathcal{P}$ (independently of the sign of Δ) decreases the rate of growth of A and increases the rate of growth of S; when θ differs from π to such an extent that the processes $S + \mathcal{P} \to A$ and $A \to \mathcal{P}$ balance each other, A reaches a maximum and then decreases. Since S continues to increase, the pump decreases to a very small value and the interaction of the waves practically ceases; a constant value of A and $\theta \simeq \Delta z$ is established. A computer solution of Eqs. (1) - (4) (Fig. 1) confirms the qualitative picture described above. We note that the qualitative behavior of the solutions remains unchanged at low pump energies, when the reaction of the anti-Stokes component on $\mathcal{P}$ and S can be disregarded. In this case the solution for A is expressed in terms of the Gauss hypergeometric

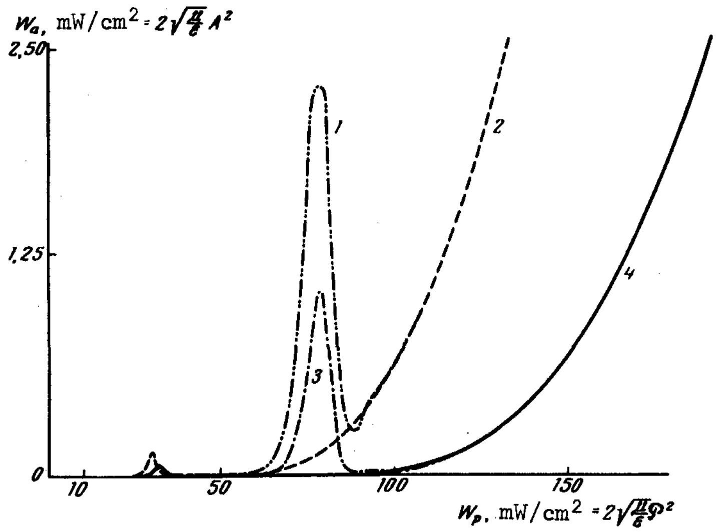

Fig. 2. Intensity of the axial ASRS at the chamber output vs. the input pump intensity: 1 - 75 atm, 5 cm; 2 - 75 atm, 12 cm ; 3 - 110 atm, 5 cm; 4 - 110 atm, 12 cm.

function ${}_2F_1[1/2 - i\delta/2;\ 1 - i\delta/2;\ 2 - i\delta/2;\ (1 + u^2 \exp 2\xi)^{-1}]$, in which $\delta = \Delta[q_s \mathcal{P}_{in}^2(1 + u^2)]^{-1}$; $\xi = q_s \mathcal{P}_{in}^2 (1 + u^2)z$; $u^2 = q_p S_{in}^2 / q_s \mathcal{P}_{in}^2$; $\mathcal{P}_{in} S_{in} = \mathcal{F}$, $S|_{z=0}$.

The asymptotic form of the solution as $z \to \infty$ is $\lim A^2 \simeq q_0^2 \pi\Delta / q_p q_s^2 \times \sinh(\pi\Delta / q_s \mathcal{P}_{in}^2) = \text{const}^{2)}$; $\theta = \Delta z$. Thus, the linear theory of [5, 6] does not hold also in the case of small conversion into the anti-Stokes component, for distances over which an appreciable fraction of the pump energy is transformed into the Stokes component, since this theory cannot account for the spatially-limited phase capture.

3. The solution for A (Fig. 1) can be used to find the relation $A^2 = f(\mathcal{P}_{in}^2)$ (see Fig. 2), for fixed lengths of the chamber and fixed gas pressures. Let us indicate the characteristic features of this relation: a) the intensity of the anti-Stokes component varies non-monotonically with increasing pump; b) an increase of the chamber length at fixed pressure should lead to a decrease of the maximum of A^2; c) the intensity of the anti-Stokes component at the output, at sufficiently powerful pump and at constant pressure, does not depend on the chamber lengths.

The purpose of the experiment was to observe these qualitative irregularities in the behavior of the axial ASRS. To this end, radiation of wave lengths 5300 Å and duration 20 nsec was focused into a chamber with compressed hydrogen in such a way that its length was always shorter than the focal region; the transverse dimension of the latter was 0.5 mm. The windows of the chamber, to avoid resonator effects, were misaligned by inclining them 3° to the pump beam axis. With increasing pump power, an anti-Stokes component (λ_a = 4350 ± 1.5 Å) was produced at the output of the chamber, first in the direction of the pump axis, and then along the generator of a cone with angle 2×10^{-2} rad.

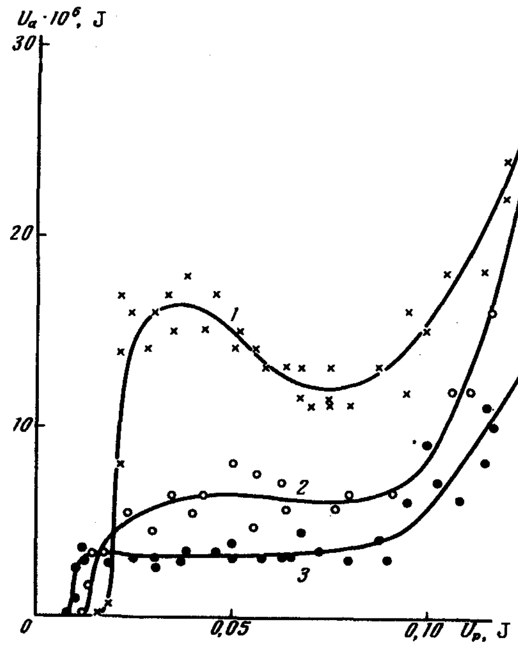

Fig. 3. Connection between the integrated energies (over the cross section and over the pulse duration) of the axial ASRS at the output and the pump at the input of the chamber: 1 - 75 atm, 5 cm; 2 - 75 atm, 12 cm; 3 - 110 atm, 12 cm.

The experimental dependence of the energy of the axial ASRS on the pump energy for two pressures and two different chamber lengths is shown in Fig. 3. As expected, the non-rectangular pump distribution in space and in time led to a smoother dependence of the integral energy of the anti-Stokes component on the pump energy than the relation between their intensities (Fig. 2). As a result, a "plateau" rather than a "dip" was observed on curves 2 and 3 of Fig. 3 (chamber length 12 cm). Calculations show that medium chamber lengths (~5 cm) are most favorable for the observation of the "dip"; such a dip is clearly pronounced on curve 1 of Fig. 3.

2) This quantity increases sharply when $\mathcal{P}_{in}^2 = \pi|\Delta|q_s^{-1}$, but the influence of A on S and $\mathcal{P}$ is essential for such values of the pump.

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.

The authors are grateful to R.V. Khokhlov for constant interest in the work and for useful discussions.

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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:

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

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