The crystals D_2SbCl_6 consist of individual octahedra, the centers of which are spaced 7.5 - 10.66 Å apart, so that two-electron delocalization is perfectly realistic, as is the case in Josephson tunnels [6]. The method of nuclear γ resonance (NGR) at liquid-nitrogen temperature has made it possible to establish the presence of antimony ions with different degrees of oxidation (Sb^{3+} and Sb^{5+}), thus indicating the effect of localization of electron pairs as a result of the cooperative process. We note that in those cases when the phase-transition temperature exceeds the critical temperature of the cooperative tunnel two-electron transition, the latter temperature can be determined by the NGR method from the coalescence of the two absorption maxima into one (in the absence of quadrupole splitting).

In subsequent papers we shall present estimates of T_k for concrete systems, for the case of strong electron-phonon interaction.

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SUPPRESSION OF STIMULATED RAMAN SCATTERING IN A DISPERSIVE MEDIUM WITH A NONLINEAR REFRACTIVE INDEX

S.A. Akhmanov, M.A. Bol'shov, K.N. Drabovich, and A.P. Sukhorukov
Physics Department of the Moscow State University

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1. We report here the results of a theoretical and experimental investigation of the features of SRS of intense and picosecond and nanosecond light pulses in self-focusing liquids. The anomalous broadening of the pulse spectra, due to the nonlinearity of the refractive index, is accompanied by an appreciable decrease in the SRS intensity, and in many cases by its complete suppression. It has been established that this effect is due to simultaneous influence of rapid phase modulation of the pump, resulting from the nonlinearity of the medium, and the dispersion of the medium. The corresponding generalization of the nonstationary SRS theory makes it possible to obtain quantitative relations for this regime, which heretofore has not been investigated.

The decrease of the SRS intensity in liquids in self-focusing beams was observed in [1 - 3], but the reason of this effect was not established. It was noted in the theoretical papers [4, 5] that this can be caused by simultaneous influence of the dispersion of the medium and phase modulation of the pump, but no concrete results are given in the cited papers.

2. Experiments on SRS in self-focusing media were performed by us with the second harmonic of a neodymium laser operating in the single-mode regime (pulse duration $\tau_p \sim 10^{-8}$ sec) and in the mode-locking regime ($\tau_p \sim 3 \times 10^{-12}$ sec). Anomalous broadening of the spectrum, reaching 1000 cm^{-1} in nanosecond pulses and more than 1000 cm^{-1} in picosecond pulses were observed in carbon disulfide in a collimated beam (see Fig. 1). At a cell length ℓ = 10 cm, a broadening larger than 500 cm^{-1} leads to a practically complete suppression of the first Stokes component of the SRS. To determine the cause of the

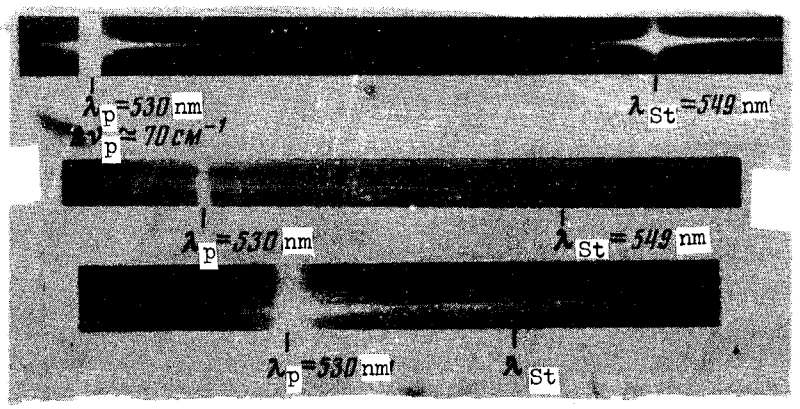

Fig. 1

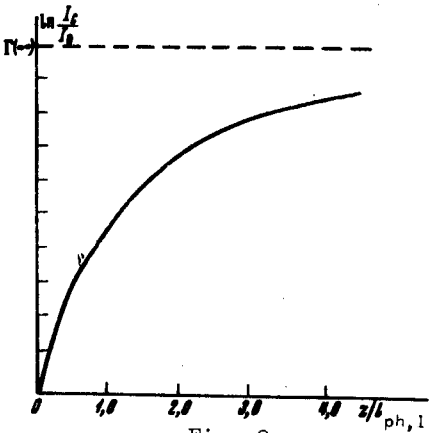

Fig. 2

Fig. 1. Spectrograms of laser pulses passing through a cell with CS_2; ℓ = 10 cm, λ_p = 530 nm. a) Spectrum of nanosecond pulse. The broadening reaches 70 cm^{-1}. One can see also the broadened band of the first Stokes component. b) Spectrum of nanosecond pulse. Broadening $\sim$1000 cm^{-1}. The Stokes component of the SRS is suppressed. c) Spectrum of picosecond pulse. The broadening exceeds 1000 cm^{-1}. The Stokes component is suppressed.

Fig. 2. Amplification of the SRS Stokes component in the field of a PM pump in a dispersive medium ($\nu \neq 0$); when $z > \ell_{ph,1}$ the intensity of the Stokes component is practically independent of z (saturation).

broadening, we measured the correlation functions of the intensity (the tracks of two-photon luminescence - TPL) of the radiation passing through the medium. It turned out that the form of the envelope remains practically unchanged. For a laser operating in the mode-locking regime, in particular, the width of the TPL spot was the same for the laser radiation and for the radiation passing through the medium. This means that the broadening of the spectrum of picosecond pulses in our experiments was uniquely connected with the phase self-modulation of the pulse in the nonlinear medium; such concepts are universally accepted in the case of nanosecond pulses.

3. Under conditions when the finite time of the transverse relaxation T_2 and the dispersion of the medium are both significant, the equations describing the SRS in a given pump field $A_p = A_p(t - z/u_p)$ are

$$\frac{\partial A_S}{\partial z} = -i\gamma_S A_p(\eta - \nu z)\sigma; \quad \frac{\partial \sigma}{\partial \eta} + T_2^{-1}\sigma = i\gamma_\sigma A_p^*(\eta - \nu z)A_S + N(\eta, z). \tag{1}$$

Here A_S is the complex amplitude of the first Stokes component, σ is the off-diagonal element of the density matrix, $\eta = t - z/u_S$, $\nu = 1/u_p - 1/u_S$, where u_p and u_S are the pump and Stokes group velocities, and N is the random force describing the spontaneous transitions in the medium.

Since allowance for the form of the amplitude modulation of the pump is of no importance in what follows (suppression of SRS takes place when the width of the pump spectrum is $\Delta\omega_p >> \tau_p^{-1}$), we shall assume the pump to be a rectangular pulse of duration τ_p with a phase that depends on the time, $A_p = A_0 \exp i\phi(t - z/u_p)$, and the effective spectrum of the PM pump is in this case

$\Delta\omega_e \simeq (\partial^2\phi/\partial t^2)^{1/2}$.

Under conditions of negligibly small dispersion of the medium ($\nu \simeq 0$), the PM pumping has practically no influence on the gain [4 - 6], and depending on the pulse duration, two SRS regimes are observed ($|A_0|^2 \sim \exp \Gamma$): when $\tau_p >> T_2$ we have $\Gamma_0 = gA_0^2 z$, and for $\tau_p < T_2$ we have $\Gamma_0 \simeq 2(2gA_0^2 z\tau_p/T_2)^{1/2}$.

4. In the case of SRS radiation of a PM pump in a dispersive medium ($\nu \neq 0$), corrections must be introduced in the gain. Let us consider first the case of relatively long pulses ($\tau_p > T_2$, $\Delta\omega_p >> \tau_p^{-1}$). We then have for the Stokes-component growth increment

$$\Gamma_1 = gA_0^2 \ell_{ph,1} \operatorname{arc\,tg} \frac{z}{\ell_{ph,1}}; \quad \ell_{ph,1} = \frac{2}{|\nu| T_2 \Delta\omega_e^2}. \tag{2}$$

It follows from (2) that when $z << \ell_{ph,1}$ the growth increment is determined by the static formula $\Gamma_0 = gA_0^2 z$. At large lengths $z >> \ell_{ph,1}$, the exponential growth of the intensity of the Stokes component stops (Fig. 2):

$$\Gamma(\infty) = \frac{\pi}{2} gA_0^2 \ell_{ph,1} \tag{3}$$

(the weak growth of the order of $|A_S|^2 \sim z$ remains, of course, in force).

Thus, to estimate the influence of the PM pumping on the gain it is necessary to determine the length $\ell_{ph,1}$. In CS_2 (for the parameters see [7]) we have $\ell_{ph,1} \leq 10$ cm if $\Delta\nu_p \geq 1$ cm^{-1}.

5. For pulses that are short compared with T_2, the same regularities apply. The limiting value of the growth factor $\Gamma(\infty)$ is determined by the same formula (3) as before. However, unlike the case of long pulses, in this case the length at which PM appears ($\Gamma \simeq 2(2gA_0^2 z\tau_p/T_2)^{1/2}$ goes over into (3)) is

$$\ell_{ph,2} = \frac{2gA_0^2}{\nu^2 T_2 \tau_p \Delta\omega_e^4} \tag{4}$$

and depends on the intensity and on the pulse duration.

6. Formulas (2) - (4) can be used for estimates in the case when the PM pumping in the nonlinear medium remains unchanged. In the experiment, such a situation is realized when SRS from not too intense picosecond pulses is observed. The PM which is invariably present in picosecond pulses lowers the gain, in accordance with the previous results; this, in our opinion, is the reason for the strong non-uniformity of the SRS along a train of picosecond pulses, observed by a number of authors (see [7]), the difficulty of obtaining SRS in the field of picosecond pulses at $\lambda = 1.06\ \mu$ in a number of media, etc.

7. If the index of the phase modulation of the wave changes in the medium (as is the case in a self-focusing medium with a refractive index $n = n_0 + n_2|E|^2$; $\partial^2\phi/\partial T^2 \simeq 2kn_2(\partial^2|A_p|^2/\partial t^2)z$), the quantitative results are altered. The characteristic length of gain saturation occurring in this case, expressed in terms of the experimentally observed spectral width $\Delta\omega_\ell$ on emerging from a sample of length ℓ, is

$$\ell_{ph,3} \approx \frac{2}{\Delta\omega_\ell}\sqrt{\frac{\ell}{|\nu|\,T_2}}. \qquad (5)$$

For CS_2 we have $\ell_{ph,3} \approx 0.3$ cm at $\Delta\nu \approx 50$ cm^{-1} and $\ell_{ph,3} \approx 0.015$ cm at $\Delta\nu \approx 1000$ cm^{-1}; this means practically complete suppression of SRS in a real experiment ($\Gamma(\infty) \approx gA_0^2\ell_{ph,3}$, cf. (3)).

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INFLUENCE OF THE MAGNETIC FIELD ON ELECTRONIC DECELERATION OF DISLOCATIONS

V.Ya. Kravchenko
Institute of Solid State Physics, USSR Academy of Sciences
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In the last few years there appeared a number of papers devoted to an experimental study of the role of electrons in processes of low-temperature plastic deformation of metals (see, e.g., [1, 2]). Usually in the experiment the electronic contribution to the deceleration of dislocations is observed on going over from the normal state to the superconducting one and vice-versa, which is realized in practice by varying the magnetic field H at a temperature $T < T_c$. Although under the experimental conditions the deformed metal in the normal state is in a magnetic field, the experimental data on the value of the electronic deceleration of dislocations agree with theoretical estimates obtained for H = 0 [3, 4], which predict the absence of a dependence on T. This apparently is connected with the fact that in fact one uses weak fields (the strong-field criterion is $r \ll \ell$, where r is the Larmor radius and ℓ is the mean free path). As will be shown below, in a strong field H there should take place a temperature and a field dependence of the electronic deceleration.

We start with the system of equations of motion of the medium with moving dislocations, the kinetic equation for the electron distribution function, and the Maxwell equations; this system is analogous to the system obtained by Kontorovich [5]:

$$\rho\ddot{u}_l = \frac{\partial}{\partial x_l}[\sigma_{ll} - \langle\chi\Lambda_{ll}\rangle] + \frac{1}{c}[\mathbf{j}\times\mathbf{H}]_l, \qquad (1)$$

$$\left(\frac{\partial}{\partial t} + v_l\frac{\partial}{\partial x_l} + \Omega\frac{\partial}{\partial\phi} + \frac{1}{\tau}\right)\chi = e\mathbf{v}\hat{\mathbf{E}} - \Lambda_{ll}\dot{w}_{ll}, \qquad (2)$$