# Raman interaction in the field of opposing light waves

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It is shown theoretically and experimentally that in four-wave Raman interaction in the field of opposing pump beams there exists a regime of unstable reversal of the wave front of a weak signal wave, such that the reflection coefficient of the wave increases exponentially in time. In the absence of a signal wave this instability leads to spontaneous generation of radiation in the reversing mirror. In the presence of an additional reflecting surface, spontaneous generation is possible also in the region of stability of the reversing mirror, provided that the self-excitation condition is satisfied for the optical resonator made up by this mirror and the reflecting surface. Self-excitation of radiation between a reversing mirror and a reflecting surface was observed in experiment, with the lasing threshold dependent on the tilt of the reflecting surface and on the optical losses.

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## 1. INTRODUCTION

The problem of four-wave Raman interaction (FWRI) of light wave in the field of opposing pump beams has recently attracted great interest. The reason is that it become possible to produce as a result of such an interaction a reversal of the wave front (WFR) of optical radiation, with a large reflection coefficient $R$.[1]

We explain here the physical mechanism that make possible in FWRI a large reflection coefficient of a weak signal wave subjected to WFR, and show that these mechanisms are due to instability effects. These effects manifest themselves in the fact that at sufficiently high pump-wave intensities the signal-wave reflection coefficient begins to increase rapidly with time. Instability effects as applied to degenerate four-wave interaction of light waves were discussed in Refs. 2–4, and the rate of the exponential growth of the reflection coefficient with time was determined in Ref. 4. That reference, however, dealt with wave interaction in a non-delaying medium, and the instability investigated there was connected with the final time of passage of the opposing light wave through a layer of a nonlinear medium.

We shall dwell below on the case of nondegenerate interaction that has a number of distinguishing features. Principal attention, both in theory and experiment, will be paid to instability effects due to the inertness of the medium as a result of the finite relaxation time of the refractive-index perturbations.

In the absence of a light wave subject to reversal, the instability effects lead to generation of radiation in the inverting mirror. In some cases, however, lasing can occur also in the region where the FWR mirror is stable. This occurs if the FWR mirror is optically coupled with some reflecting surface in conjunction with which it forms an optical resonator.[5] If the WFR mirror has a sufficiently large reflection coefficient, such a resonator becomes self-excited, and the transverse structure of the light beam generated in it can "attune itself" to the reflecting surface in a way that the diffraction losses of the radiation in the resonator are a minimum. This self-excitation in the resonator can be regarded, on the one hand, as parasitic, since it affects adversely the properties of the FWR mirror. On the other hand, the investigated self-excitation effects in the resonator are of independent interest for the development of methods of "self-attuning" (adaptive) generation of radiation in extended resonators, the development of interferometers for the measurement of fluctuations in the lengths of such resonators, measurement of small optical-radiation frequency shifts, etc.

We obtain the temporal growth rates of the field inside the resonator made up of the FWR mirror and the reflecting surface, and show that the summary intensity of the opposing pump beams needed for its self-excitation has a periodic dependence on the resonator length and on the difference between the frequencies of the waves generated in it, and is in the general case a nonmonotonic function of the optical losses.

Self-excitation of optical radiation was observed experimentally in a resonator formed by a reflecting surface and the FWR mirror, and consisting of a cell with acetone—a medium active to stimulated Brillouin scattering (SBS). This cell was pumped by two opposing beams in the field of which there was realized the FWRI that lead to the WFR of the radiation generated in the resonator. The divergence of this radiation was 2–3 times larger than the diffraction divergence. The lasing threshold depended on the optical losses in the resonator and on the ratio of the frequencies of the opposing pump beams.

## 2. THEORY OF INSTABILITY EFFECTS IN A WFR MIRROR AND OF THE SELF-EXCITATION OF RADIATION BETWEEN THIS MIRROR AND A REFLECTING SURFACE

To attain maximum possible values of $R$ in parametric mixing of the light wave to be reversed and the more powerful pump waves, a number of conditions must be satisfied. One is that the realization of large reflection coefficients must not lead to a growth of the waves primed by parasitic components due to the inaccuracy of the WFR of the pump beams relative to each other, diffraction of the pump waves by the edges of the cell, scattering of these waves by optical inhomogeneities of the medium, and others. It follows therefore that the frequency of the reversed wave, and hence the

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frequency of the wave subject to inversion, should differ from the frequencies of the high-power pumps.

Another condition stems from the fact that in an optical experiment there is always a relatively weak backscattering of the reversed wave from some reflecting surface. This scattering returns radiation at the frequency of the inverted to the WFR mirror, causing secondary reflection. If the reflection coefficient $R$ of the WFR mirror is large enough, this process can lead to spontaneous generation of optical radiation between the reversing mirror and the reflecting surface. To exclude this generation it is necessary that the secondary reflection of the already reversed wave from the FWR mirror be negligible. This condition is compatible with the requirement that the reflection coefficients be large if the frequencies of the wave to be reversed and the reversed wave are different.

Both conditions on the frequencies of the interacting wave can be satisfied by the WFR process in FWRI in which the perturbations of the refractive index are excited as a result of interference between light waves having frequencies that differ by an amount close to the natural frequency $\Omega_0$ of these perturbations. The most typical example of such an interaction is FWR in FWRI with participation of hypersound.[6] Here, in contrast to the ordinary SBS, the medium is illuminated from opposite sides by two opposing pump beams with complex amplitudes $A_0^+$ and $A_0^-$ with frequencies $\omega_0^+$ and $\omega_0^-$. The frequency $\omega_0^-$ of the pump $A_0^-$ is detuned by an amount $\Omega_0$ from the frequency $\omega_1^+$ of the signal wave $A_1^+$, which is fed into the same medium at an angle close to 180° relative to the wave $A_0^-$. The interaction of the waves $A_0^-$ and $A_1^+$ excites a hypersound wave, and scattering of the pump $A_0^+$ by this wave leads to emission of a reversed wave of frequency $\omega_1^- = \omega_0^+ + \omega_0^- - \omega_1^+$ (Fig. 1). We examine now the formation and growth of this wave with time in the absence and in the presence of a reflecting surface outside the volume of the nonlinear medium.

## 1. Development of unstable perturbations in FWRI of opposing light waves in a WFR mirror

In a given field of opposing pump beams $E_0^\pm = A_0^\pm \exp(i\omega_0^\pm t \mp k_0^\pm z) +$ c.c., the resonant FWRI between the signal light wave

$$\mathbf{E}_1^+ = \mathbf{A}_1^+ \exp(i\omega_1^+ t - ik_1^+ z) + \text{c.c.},$$

and the reversed wave

$$\mathbf{E}_1^- = \mathbf{A}_1^- \exp(i\omega_1^- t + ik_1^- z) + \text{c.c.},$$

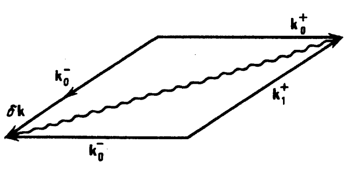


FIG. 1. Wave-vector diagram illustrating the WFR effect in FWRI. The wave line indicates the hypersound wave vector (its direction is in the positive $z$ direction at $\omega_1^+ > \omega_0^-$ and in the negative at $\omega_1^+ < \omega_0^-$).

with participation of a phonon wave

$$Q = q \exp[i(\omega_1^\pm - \omega_0^-)t - i(k_0^- + k_1^+)z] + \text{c.c.}$$

having a frequency $\Omega_0 = |\omega_0^- - \omega_1^+|$ is described in the quasioptical approximation by a system of three equations:

$$\frac{1}{v_1^+}\frac{\partial \mathbf{A}_1^+}{\partial t} + \frac{\partial \mathbf{A}_1^+}{\partial z} + \frac{i}{2k_1^+}\Delta_\perp \mathbf{A}_1^+ = -i\frac{k_1^+(\partial\varepsilon/\partial Q)}{2\varepsilon}\mathbf{A}_0 q, \tag{1}$$

$$\frac{1}{v_1^-}\frac{\partial \mathbf{A}_1^-}{\partial t} - \frac{\partial \mathbf{A}_1^-}{\partial z} + \frac{i}{2k_1^-}\Delta_\perp \mathbf{A}_1^- = -i\frac{k_1^-(\partial\varepsilon/\partial Q)}{2\varepsilon}\mathbf{A}_0 q^* \exp(i\delta kz), \tag{2}$$

$$\frac{\partial q}{\partial t} + \frac{1}{T_2}q = -i\beta(\mathbf{A}_1^+\mathbf{A}_0^{-*} + \mathbf{A}_0^+\mathbf{A}_1^{-*})\exp(-i\delta kz). \tag{3}$$

Here

$$k_1^\pm = \omega^\pm/v_1^\pm, \quad \omega_1^- = \omega_0^+ - \omega_0^- - \omega_1^+, \quad \delta k = k_0^+ + k_1^- - k_0^- - k_1^+,$$

$T_2$ is the relaxation time of the phonon wave, $\beta$ is the nonlinearity coefficient, and $v_1^\pm$ are the wave velocities. In the derivation of (1)–(3) it was assumed that the pump waves do not interact with each other ($||\omega_0^- - \omega_0^+| - \Omega_0| \gg T_2^{-1}$), and it was also assumed that $\omega_1^+ - \omega_0^- = \Omega_0 > 0$; on the other hand if $\omega_1^+ - \omega_0^- = -\Omega_0 < 0$, it is necessary to make in (1)–(3) the substitutions $q \to q^*$ and $\beta \to -\beta$.

Assume that the amplitudes $\mathbf{A}_0^+$ and $\mathbf{A}_0^-$ are independent of the coordinates, and put in addition $k_1^+ = k_1^-$ in the left-hand sides of (1) and (2), i.e., we neglect the diffraction mismatch of the spatial distribution of the fields $\mathbf{A}_1^+$ and $\mathbf{A}_1^-$ of the signal and reversed waves within the volume of the nonlinear medium. This is justified if the characteristic length $z_n$ of the nonlinear interaction is small compared with $z_{\text{scat}} = \rho_1^2 |1/k_1^+ - 1/k_1^-|^{-1}$ where $\rho_1$ is the scale of the transverse modulation of the waves $\mathbf{A}_1^+$ and $\mathbf{A}_1^-$. Then, when a spatially coherent signal wave

$$\mathbf{A}_1^+(0, r_\perp, t) = \mathbf{c}_1^+(0, t)\psi_1^+(0, \mathbf{r}_\perp)$$

is incident on the nonlinear medium, it is convenient to seek the solution of Eqs. (1)–(3) in the form

$$\begin{gathered} \mathbf{A}_1^+ = \mathbf{c}_1^+(z, t)\psi_1^+(z, \mathbf{r}_\perp), \\ \mathbf{A}_1^- = \mathbf{c}_1^-(z, t)\psi_1^{+*}(z, \mathbf{r}_\perp), \quad q = \bar{q}(z, t)\psi_1^*(z, \mathbf{r}_\perp), \end{gathered} \tag{4}$$

where $\psi_1^+$ satisfies Eq. (1) with zero right-hand side. Substitution of (4) in (1)–(3) yields

$$\frac{1}{v_1^+}\frac{\partial \mathbf{c}_1^+}{\partial t} + \frac{\partial \mathbf{c}_1^+}{\partial z} = -i\frac{k_1^+(\partial\varepsilon/\partial Q)}{2\varepsilon}\mathbf{A}_0^-\bar{q}, \tag{5}$$

$$\frac{1}{v_1^-}\frac{\partial \mathbf{c}^-}{\partial t} - \frac{\partial \mathbf{c}_1^-}{\partial z} = -i\frac{k_1^-(\partial\varepsilon/\partial Q)}{2\varepsilon}\mathbf{A}_0^+\bar{q}^* \exp(-i\delta kz), \tag{6}$$

$$\frac{\partial \bar{q}}{\partial t} + \frac{1}{T_2}\bar{q} = -i\beta(\mathbf{A}_0^{-*}\mathbf{c}_1^+ + \mathbf{A}_0^+\mathbf{c}_1^{-*}\exp(-i\delta kz)). \tag{7}$$

If the time of travel of the light through the nonlinear medium is short compared with the pump-pulse duration, as well as compared with the relaxation time $T_2$, we can neglect in (5)–(7) the term

$$(v_1^\pm)^{-1}\partial \mathbf{c}_1^\pm/\partial t.$$

A beneficial characteristic of the system described by Eqs. (5)–(7) is its response to a short signal-wave pulse $c_1^+(0, t) = \delta(t)$. The Laplace transform of this response

$$L(p) = \int_0^\infty dt\, c_1^-(0, t)e^{-pt} \quad \text{at} \quad c_1^-(L, t) = 0, \quad \bar{q}(z, 0) = 0$$

is equal to (at $\omega_1^+ > \omega_0^-$ and $\delta k = 0$)

$$L(p)=\frac{A_0^+A_0^-(\exp[(M^++M^-)/2(1+pT_2)]-1)}{I_0^-\exp[(M^++M^-)/2(1+pT_2)]+I_0^+}, \qquad (8)$$

where $M^\pm = gI_0^\pm L$, $g = 2\pi k(\partial\varepsilon/\partial Q)\beta/\varepsilon^2 v$ is the growth rate of the stimulated scattering per unit length and per unit intensity, $I_0^\pm$ are the pump intensities; it is assumed here for simplicity that $v_1^+ = v_1^- = v$ and $k_1^+ = k_1^- = k$. At $\omega_1^+ > \omega_0^-$ we have $g > 0$. If, however, $\omega_1^+ < \omega_0^-$, Eq. (8) remains in force, but in this case $g < 0$, while $A_0^+A_0^-$ must be replaced by $(A_0^+A_0^-)^*$.

Taking the inverse Laplace transform of (8) we obtain the response to a delta function and this, by virtue of the linearity of Eqs. (5)–(7) enables us to determine the response to an arbitrary function. In particular, the unit function $c_1^+(0,t) = 1(t)$ has a response that coincides with the reflection coefficient $R(t)$, namely

$$R(t)=\frac{A_0^-}{A_0^{+*}}\left\{\exp\left(-\frac{t}{T_2}\right)\sum_{n=0}^{\infty}\left(-\frac{I_0^-}{I_0^+}\right)^n[f(n+1,t)-f(n,t)]\right.$$
$$\left.+\frac{1}{T_2}\int_0^t dt'\exp\left(-\frac{t'}{T_2}\right)\sum_{n=0}^{\infty}\left(-\frac{I_0^-}{I_0^+}\right)^n[f(n+1,t')-f(n,t')]\right\}, \qquad (9)$$

where

$$f(n,t)=I_0(2n(M^++M^-)t/T_2)^{1/2},$$

and $I_0(x)$ is a modified Bessel function. Equation (9) is convenient for the calculation of $R(t)$ at $I_0^- < I_0^+$, since all the terms that enter in it decrease at sufficiently large $n$. On the other hand if $I_0^- > I_0^+$, it is more convenient to use a different form of the same formula, obtained by making the substitutions $A_0^- \to A_0^+$ and

$$f(n,t)\to -J_0(2n(M^++M^-)t/T_2)^{1/2},$$

where $J_0(x)$ is a Bessel function. It follows also from (9) that

$$R(t,g,I_0^+,I_0^-)=-R(t,-g,I_0^-,I_0^+).$$

At large values of the argument, the modified Bessel function takes the form $I_0(x) \approx e^x/(2\pi x)^{1/2}$ ($x \gg 1$). Therefore at sufficiently large values of $t$ the sum in (9) contains terms proportional to the factor

$$\exp\left[\left(\frac{2(M^++M^-)nt}{T_2}\right)^{1/2}-n\ln\frac{I_0^+}{I_0^-}+in(\pi+2\pi k)\right]$$

($k = 0, \pm1, \ldots$), which reaches at

$$n=\frac{(M^++M^-)t}{2T_2[\ln(I_0^+/I_0^-)-i\pi(2k+1)]^2}$$

a maximum equal to

$$\exp\left\{\frac{(M^++M^-)[\ln^2(I_0^+/I_0^-)+i\pi(2k+1)]t}{2T_2[\ln^2(I_0^+/I_0^-)+\pi^2(2k+1)^2]}\right\}.$$

It follows therefore that the total gain is determined by the growth rate, whose real and imaginary parts $p_k'$ and $p_k''$ are given by

$$T_2p_k'=\frac{{}^1/_2(M^++M^-)\ln(M^+/M^-)}{(\ln(M^+/M^-))^2+\pi^2(2k+1)^2}-1, \qquad (10)$$

$$T_2p_k''=\frac{(2k+1)\pi{}^1/_2(M^++M^-)}{(\ln(M^+/M^-))^2+\pi^2(2k+1)^2}. \qquad (11)$$

Equations (10) and (11) can be obtained also by another method, which reduces to setting the denominator of (8) equal to zero and finding in this manner the poles of the function $L(p)$.

At $\mathrm{Re}\,P_k > 0$ the value of $R(t)$ increases exponentially with time and the WFR process is unstable. This means that when a nonlinear medium is irradiated by two opposing pump beams, spontaneous generation of noise radiation is possible. The minimum value of the parameter $M^+ + M^-$ at which such generation is possible is reached at $k = 0$ and 1, and is equal to $4\pi$; in this case $|\ln(I_0^+/I_0^-)| = \pi$. If, for example, $\omega_1^+ > \omega_0^-$, at a constant ratio $I_0^+/I_0^- = e^\pi$ the dependence of the maximum growth rate $p_{0,-1}$ on the excess over the instability threshold $\rho = (M^+ + M^-)/4\pi$ is written in the form

$$p'_{0,-1}=(\rho-1)T_2, \quad p''_{0,-1}=\pm\rho T_2^{-1}. \qquad (12)$$

The nonvanishing of $p_0''$ and $p_{-1}''$ means that after the threshold is exceeded the reversed wave has two growing components that are displaced by $\rho T_2^{-1}$ in both directions relative to the frequency $\omega_1^- = \omega_0^+ + \Omega_0$.

## 2. Conditions for self-excitation of optical radiation

If $\mathrm{Re}\,p_k < 0$, the WFR process is stable. At sufficiently large values of the growth rate $M^+ + M^-$, however, the reflection coefficient $R$ can reach very high values even in the stability region. Under these conditions, a substantial role can be assumed by effects connected with the backscattering of the reversed wave from some reflecting surface. If, for example, the pump frequencies are equal ($\omega_0^+ = \omega_0^-$), the backscattered wave of frequency $\omega_1^- = \omega_0^+ + \omega_0^- - \omega_1^+$, returning to the cell, is transformed with WFR into a wave of frequency $\omega_1^+$. The latter in turn, when scattered by the indicated surface in the backward direction, again returns to the cell, and the entire re-reflection cycle is repeated. It is clear that this can lead to self-excitation of the optical radiation between the reversing mirror and the reflecting surface.

We now obtain the conditions for this self-excitation at $\omega_0^+ = \omega_0^-$. Assume that at a distance $\mathscr{L}$ from the reversing mirror there is located a reflecting surface with a reflection coefficient $r$. It is easily seen that the optical resonator made up of the WFR mirror and this reflecting surface becomes self-excited if the real part of the root of the dispersion equation is positive (the pump waves $A_0^+$ and $A_0^-$ are assumed plane):

$$|r|^2\exp(i\psi-p\tau)L(p)L'(p)-1=0, \qquad (13)$$

where $\tau = 4\mathscr{L}/v$ and $\psi = 2(\omega_1^+ - \omega_1^-)\mathscr{L}/v$ are the travel time of the light and the phase difference over the entire re-reflection cycle. The functions $L(p)$ and $L'(p)$ describe WFR of waves of frequency $\omega_1^+ > \omega_0^-$ and $\omega_1^- < \omega_0^+$, respectively. Accordingly, the expression for $L(p)$ is determined by Eq. (8) with $g = |g|$, while $L'(p)$ is determined by the same Eq. (8) but with $g$ replaced by $-g$ and $A_0^+A_0^-$ by $(A_0^+A_0^-)^*$.

We now obtain the solution of Eq. (13) in the limiting case when the time $\tau$ is short compared with the characteristic times of the change of the generated wave. Putting in this approximation $\exp(-p\tau) = 1$, we obtain

$$p_k'T_2=\frac{(M^++M^-)\ln(x_1^2+x_2^2)}{\ln^2(x_1^2+x_2^2)+4[\mathrm{arctg}(-x_2/x_1)+\pi\alpha_k]^2}-1, \qquad (14)$$

$$p_k''T_2 = \frac{2[\mathrm{arctg}(-x_2/x_1)+\pi\alpha_k](M^++M^-)}{\ln^2(x_1^2+x_2^2)+4[\mathrm{arctg}(-x_2/x_1)+\pi\alpha_k]^2}. \qquad (15)$$

Here $x_1$ and $x_2$ are the real and imaginary parts of the complex number $x$, which is the solution of the algebraic equation

$$x^2+\left[\frac{a^2+1-2a|r|^2e^{i\psi}}{a(1+|r|^2e^{i\psi})}\right]x+1=0 \qquad (16)$$

and depends on the parameters $a = I_0^-/I_0^+$, $\psi$, $|r|^2$. The value of $\alpha_k$ indicates the number of the longitudinal mode; $\alpha_k = 2k$ if $x_1 > 0$ and $\alpha_k = 2k+1$ if $x_1 < 0$.

At $|r|^2 = 0$ Eqs. (14) and (15) go over into (10) and (11). It follows from (14) that the self-excitation condition is written in the form

$$M^++M^- \geqslant M_k^{cr} = \frac{\ln^2(x_1^2+x_2^2)+4[\mathrm{arctg}(-x_2/x_1)+\pi\alpha_k]^2}{\ln(x_1^2+x_2^2)}. \qquad (17)$$

The function $M_k^{cr}(a, |r|^2, \psi)$, defined by Eq. (17), is rather cumbersome. We shall therefore analyze only several particular cases.

We put first $a = |r|^2 = 1$. In this case $x_1 = 0$, $x_2 = \tan(\psi/4 \pm \pi/4)$, and

$$M_k^{cr} = \frac{\pi^2(2k+1)^2+(2\ln|\mathrm{tg}(\psi/4\pm\pi/4)|)^2}{|2\ln|\mathrm{tg}(\psi/4\pm\pi/4)||}. \qquad (18)$$

The minimum of $M_k^{cr} = 2\pi|2k+1|$ is reached at

$$|\mathrm{tg}(\psi/4\pm\pi/4)| = \exp(\pi/2).$$

At $\psi = \pi n$ the values of $M_k^{cr}$ go off to infinity (Fig. 2). We see thus that a change of the length $\mathscr{L}$ leads to variations of the phase difference $\psi$, causing thereby oscillations of the "critical" growth rate $M_k^{cr}$. By measuring the period of the oscillations it is easy to determine the value of $\Omega_0 = (\omega_1^+ - \omega_1^-)/2$. This effect is a consequence of the fact that in the approximation considered to travel time of the light $\tau$ is short compared with the relaxation time $T_2$. In this case the efficiency of hypersound excitation depends essentially on the relative phases of the two generated waves with frequencies $\omega_1^+$ and $\omega_1^-$.

We consider now self-excitation under conditions when $|r|^2 = a$. (This situation is realized in experiments similar to that described below, where the reflecting surface for the generated waves $A_1^+$ and $A_1^-$ is the same flat mirror that transforms the pump $A_0^+$ into

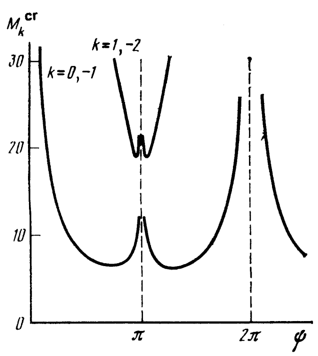


FIG. 2. $M_k^{cr}$ vs the phase difference $\psi$ at $a = |r|^2 = 1$.

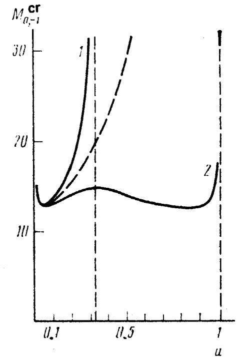


FIG. 3. $M_{0,-1}^{cr}$ vs $a$ at $|r|^2 = a$ [curve 1—$\psi = 2\pi n$, 2—$\psi = (2n+1)\pi$]. The dashed lines indicate the plot of $M_{0,-1}^{cr}(a)$ at $|r| = 0$.

a backward pump $A_0^-$ of intensity $I_0^- = aI_0^+$.) We confine ourselves for simplicity to a discussion of the functions $M_{0,-1}^{cr}(a)$ only for two fixed values of the phase difference: $\psi = 2\pi n$ and $\psi = (2n+1)\pi, n = 0, \pm 1 \ldots$. In both indicated cases $x_2 = 0$. The functions $M_{0,-1}^{cr}(a)$ are shown in Fig. 3 [curves 1 and 2 are for $\psi = 2\pi n$ and $\psi = (2n+1)\pi$, respectively]. For comparison, the same figure shows a plot of $M_{0,-1}^{cr}(a)$ at $|r| = 0$ (dashed curve). It follows from Fig. 2 that at $a > 1/3$ no self-excitation is reached at all at a phase difference $\psi = 2\pi n$.

It is also of interest to consider the dependence of $M_k^{cr}$ on $|r|^2$ at fixed values of $a$ and $\psi$. In particular, it is easy to show that at $a = 1$ and $\psi = 2\pi n$ there is no self-excitation, but at $\psi = (2n+1)\pi$ the function $M_{0,-1}^{cr}(|r|)$ is nonmonotonic (Fig. 4). This behavior is again due to the influence of the phasing of the generated wave on the hypersound excitation.

The fact that the total growth rate $M_0^+ + M_0^-$ exceeds $M_k^{cr}$ means that optical radiation is generated between the reversing mirror and the reflecting surface. Furthermore, as seen from (14), each of the longitudinal modes of this radiation has its own growth rate, so that ultimately mode selection takes place and some of the modes predominate.

We shall not consider here the selection of transverse modes. We note only that modes with minimum diffraction losses, which "attune" themselves to the re-

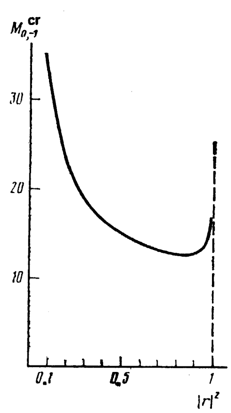


FIG. 4. $M_{0,-1}^{cr}$ vs $|r|^2$ at $a = 1$ and $\psi = (2n+1)\pi$.

flecting surface, will have the largest growth rate when the phase difference is optimal and the dependence of $p_k'$ on $|r|$ is monotonic. At an arbitrary phase difference $\psi$, however, this is not always the case, as is exemplified by the situation illustrated in Fig. 4. In such cases, in principle, the self-excited mode can be not the one with the minimum loss but with a certain optimal diffraction loss.

To prevent the discussed self-excitation of optical radiation between the reversing mirror and the reflecting surface, the wave interaction must be effected in such a way that the reflected reversed wave is no longer subject to the WFR process. It is easily verified that this calls for the use of pumps with different frequencies. If

$$|\omega_0^+-\omega_0^-|\gg T_2^{-1} \text{ and } ||\omega_0^+-\omega_0^-|-2\Omega_0|\gg T_2^{-1},$$

then the frequency of the reflected backward-reversed wave $\omega_1^-$ differs from the frequency $\omega_0^-$ of the pump that interferes with it by an amount larger (or smaller) than the natural oscillation frequency $\Omega_0$. In this case there is no effective hypersound excitation and the secondary reflection of the radiation from the reversing mirror is negligible. If, however, backscattering produces on the reflecting surface a mixing of the radiation frequency, then self-excitation can occur in principle also if the pump frequencies are different. (The required frequency shift can be easily obtained, for example, if the light scattered by the reflecting surface is again aimed on the same surface with its frequency changed by another reversing mirror.)

## 3. EXPERIMENTAL INVESTIGATION OF THE EFFECTS OF INSTABILITY AND SELF-EXCITATION

### 1. Unstable WFR regime

The WFR was investigated in a cell with acetone ($L$ = 3 cm) pumped by two opposing single-mode beams from a neodymium laser with pulse duration 30 nsec and beam diameter 1 mm. [One of the beams ($A_0^-$) was produced by reflecting the other ($A_0^+$) from a plane glass plate placed across the light beam at a distance $\mathscr{L}$= 10 cm behind the cell with the nonlinear medium.] The weak signal wave $A_1^+$ was the radiation from the same laser, shifted in frequency into the anti-Stokes region by an amount equal to the SBS shift, and attenuated (in intensity) by a factor $10^4$–$10^6$ relative to the high-power pump wave $A_0^+$. The weak wave entered the cell at an angle ~1° to the pump wave $A_0^+$. The presence of an unstable WFR regime was detected by an abrupt (threshold-like) increase of the measured reflection coefficient $R$ and of the reversed-wave power $p_1^-$ as the pump intensity $I_0^+$ was increased (the value of $R$ reached ~$10^5$ and more, Ref. 1). Figure 5 shows for two values of the ratio $I_0^+/I_1^+(0)$ and for $a = 2.5\times10^{-3}$ the threshold pump intensities $I_{0\,\mathrm{thr}}$ normalized with respect to $1/gL$. (The threshold was determined from the condition that the reversed-wave intensity $I_1^-$ reach 2% of the pump intensity $I_0^+$.) The same figure shows the threshold for the observation of the noise radiation in the absence of a signal wave (the value of $\ln[I_0^+/I_1^+]$ is assumed in this case to be 25). The solid line of Fig. 5

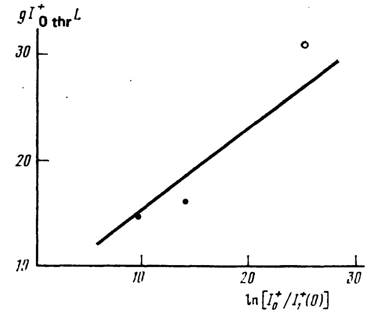


FIG. 5. Threshold dependence of $I^+_{0,\mathrm{thr}}$ vs $\ln[I_0^+/I_1^+(0)]$ (•—experiment in the presence of a signal wave, ○—experiment in the absence of a signal wave, solid line—theoretical calculation).

indicates (in the same dimensionless units) the approximate theoretical dependence of the threshold intensity $I_{0\,\mathrm{thr}}$ on $\ln[I_0^+/I_1^+(0)]$, calculated using Eqs. (9) and (10). The results confirm the existence of the unstable WFR regime.

To obtain high-quality WFR in this regime, the signal-wave power must not be too small; it is necessary at least that it exceed substantially the power of the priming noise source, as determined by the spontaneous scattering. It is clear also that the signal-wave pulse must enter the nonlinear medium either ahead of or simultaneously with the pump pulse, otherwise the growth of the unstable perturbations, which are produced by the spontaneous scattering, will decrease the intensity of the pump field in the cell even prior to the arrival of the signal wave. The accuracy of the WFR was qualitatively investigated precisely under these conditions: the pulses of the two pumps and of the signal entered the cell simultaneously, the intensity of the signal wave was not less than $10^{-6}$ of the intensity of the pump $A_0^+$. It was found when the divergence of the signal wave was equal to its diffraction value (1.5 $\times 10^{-4}$ rad), the divergence of the reversed wave was not more than 2.4 times larger, even when the pump intensity was comparable with the threshold $I^+_{0\,\mathrm{thr}}$ of the observation of the noise radiation in the absence of a signal.

### 2. Self-excitation of radiation between the WFR mirror and the reflecting surface

To investigate the self-excitation of optical radiation between the reversing mirror and the reflecting surface, we used the fact that the glass plate that reflected the pump beam $A_0^+$ was also the reflecting surface for the wave that could be generated in the resonator made up of the reversing mirror and this plate.

To observe the self-excitation effects, the WFR mirror was irradiated by only two pump waves, and no signal wave was applied. Under these conditions we registered, from the direction of the stronger pump, the angular spectrum of the radiation propagating along the pump beams. When the surface reflecting the opposing pump was covered, we observed the usual SBS angular distribution picture corresponding to incoherent

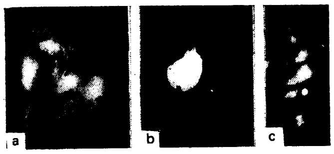

FIG. 6. Transverse distribution of angular spectrum of emission from cell with acetone in the field of two opposing pumps (a—weak pump obstructed, b—frequencies of both pumps equal, c—frequencies of the two pumps different).

scattering the field of a single-mode beam [Fig. 6(a)].[7] [In the experiment we photographed on the same frame of the film, besides the SBS radiation, also the angular spectrum of the high-power single-mode pump beam (bright spot at the center of the photograph in Fig. 6)].

When the surface that reflected the opposing pump was uncovered, however, the transverse distribution of the observed angular spectrum changed [Fig. 6(b)]. The relative intensities of the higher modes decreased, and the angular spectrum revealed most clearly the weakly diverging component with angular-spectrum width $\sim 4\times 10^{-4}$ rad. (The opposing pump beam was so weak that it produced practically no imprint on the film.) The generation threshold of this weakly diverging radiation was lower than the SBS threshold. For example, at $a = |r|^2 = 2\times 10^{-2}$ it was half the SBS threshold in the same cell with acetone. The generated radiation pulse duration was ~15 nsec. When the reflectance of the surface was decreased to $4\times 10^{-3}$ the generation stopped and the picture of the angular spectrum hardly differed from that of ordinary SBS.

The result indicated that self-excited optical radiation is produced in the resonator made up of the WFR mirror and the reflecting surface. This affects adversely the quality of the WFR. First, the radiation generated in the resonator is superimposed on the inverted wave and distorts its properties. Second, the generation of the indicated radiation weakens the intensities of the pump waves, thereby decreasing the coefficient of reflection of the useful signal with WFR. At the same time, the observed self-excitation of optical radiation is of independent interest for the investigation of the feasibility of generation of a "self-attuning" light beam in a resonator with a WFR mirror.

To suppress this self-excitation we must decrease the coefficient of reflection of the opposing pump from the glass plate (in our experiment, as already indicated, the self-excitation vanished at a coefficient smaller than $4\times 10^{-3}$). In a number of cases, however, it is more convenient to use pump waves with frequencies that differ by much more than the half-width of the spontaneous Brillouin scattering line $\delta\omega = T_2^{-1}$ in the nonlinear medium (acetone). To verify this premise, the reflecting glass plate was replaced by a cell with compressed xenon ($p = 38$ atm) on which SBS by the strong pump $A_0^+$ produced a weak opposing pump $A_0^-$. Its frequency shift relative to the high-power pump was 0.01 cm$^{-1}$, or ten times the value $\delta\omega/2\pi c \approx 10^{-3}$ cm$^{-1}$ typical of acetone. Under these conditions the observed (from the side of the strong pump beam) picture of the angular spectrum of the radiation from the cell with the acetone differed greatly from the case of equal pump frequencies and recalled the angular spectrum of ordinary SBS [Fig. 6(c)]. Thus, a frequency difference between the opposing pump beams eliminates the self-excitation between the WFR mirror and the reflecting surface.

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