Effect of Laser Mode Structure on Stimulated Brillouin Scattering

PAUL NARUM, MARK D. SKELDON, AND ROBERT W. BOYD

Abstract—**The gain and reflectivity of the stimulated Brillouin scattering (SBS) process used as a Stokes-wave generator are shown theoretically to be independent of the mode structure of the pump laser provided that the pump-laser mode spacing exceeds the Brillouin linewidth and that the laser coherence length exceeds the characteristic gain length of the SBS process. Under the same set of conditions, the gain of an SBS amplifier is found to depend on the degree of correlation between the laser and Stokes fields. However, due to nonlinear coupling, these two fields become correlated within several characteristic gain lengths, and the subsequent propagation of the two fields is governed by the same set of equations that apply for the case of a single-mode pump laser. These theoretical predictions are tested experimentally for an SBS generator using acetone, carbon disulfide, and methanol as the Brillouin-active media, and the results are in full agreement with the theoretical predictions.**

INTRODUCTION

STIMULATED Brillouin scattering (SBS) is a process that is useful for producing the phase conjugate of an arbitrary wavefront [1], for beam combining [3], and for pulse compression [2]. The theoretical description of SBS for the case of a single-mode pump laser is well known and has been reviewed by Kaiser and Maier [4]. Many applications of SBS involve the use of multimode pump lasers. For example, high-energy lasers must often be operated multilongitudinal mode in order to extract all of the energy stored in an inhomogeneously broadened gain medium [5], [6]. The extent to which the efficiency of the SBS process is degraded through the use of a multimode pump laser has not been entirely clear. In fact, in many circumstances, the efficiency of the SBS process is badly degraded through the use of a multimode pump laser [7], [8].

There has been considerable theoretical and experimental work on stimulated scattering with broad-band lasers, including both stimulated Raman [9]–[12] and stimulated Brillouin scattering [13]–[17]. D'yakov [13] has treated theoretically the case of a laser with a continuous spectrum whose width Γ_L is much larger than the Brillouin linewidth Γ in the limit of negligible pump depletion. He showed that the threshold for the SBS process is independent of the laser bandwidth as long as the laser coherence length exceeds the characteristic gain length for SBS. It was also predicted that the spectrum of the scattered light replicates that of the pump laser. These predictions were verified in part in experiments of Popovichev *et al.* [14] and Zubarev and Mikhailov [15].

In this paper, we consider SBS with a multilongitudinal-mode pump laser. We assume that the width of each mode is much smaller than the Brillouin linewidth, but allow the total laser linewidth to be arbitrarily broad. Our theoretical approach is similar to that of Zel'dovich and Shkunov [16], but is extended to include pump depletion, optical absorption, and thermal contributions to the Brillouin gain. We show theoretically and experimentally that under a set of conditions that can readily be achieved in practice, the characteristics of the SBS process are independent of the number of modes present in the pump laser. In particular, we show that the spatial evolution of the laser and the Stokes fields are described by three coupled nonlinear differential equations that describe the spatial evolution of the intensities and of the correlation function C between the two fields. Analogous equations have been shown to hold for stimulated Raman scattering in the forward direction [11]. For the case of an SBS generator, we show that these three equations reduce to the two equations that describe the SBS process with a single-mode pump laser. For the case of an SBS amplifier, the initial gain depends on the initial value C_0 of the correlation function and is always smaller than the single-mode gain. However, the nonlinear interaction drives C towards unity, and after the Stokes wave has propagated a distance approximately equal to the inverse of the product of C_0 with the gain per unit length, the equations that describe the interaction are again the same as those for the case of a single-mode pump laser. The main criteria that have to be met in order for the SBS to be independent of the laser mode structure are that the spacing between the individual laser modes must be larger than the Brillouin linewidth and that the coherence length of the laser must be larger than the characteristic gain length of the SBS process. The theoretical predictions have been tested experimentally in an SBS generator, and excellent agreement between theory and experiment has been found.

THEORY

In this section, we develop the theory that describes stimulated Brillouin scattering with a multilongitudinal-mode laser. For completeness, we allow the possibility that the scattering medium exhibits linear absorption, and

Manuscript received March 11, 1986; revised May 30, 1986. This work was supported by the New York State Center for Advanced Optical Technology.

The authors are with the Institute of Optics, University of Rochester, Rochester, NY 14627.

IEEE Log Number 8610273.

thus our treatment encompasses both SBS and its thermal analog, stimulated thermal Brillouin scattering (STBS). We assume that the laser field is linearly polarized and propagates in the positive z direction and can be represented by the plane wave expansion

$$E_L(z, t) = \tfrac{1}{2} \sum_n E_{L,n}(z)\, e^{-i\omega_n(t - n_0 z/c)} + \text{c.c.} \tag{1}$$

Here c is the speed of light in vacuum and n_0 is the refractive index, which we take to be the same for all modes. We also assume that the Brillouin Stokes field propagates in the negative z direction and can be represented as

$$E_S(z, t) = \tfrac{1}{2} \sum_m E_{S,m}(z)\, e^{-i\nu_m(t + n_0 z/c)} + \text{c.c.} \tag{2}$$

We have used the convention here and throughout the paper that ω refers to a laser frequency and ν to a Stokes frequency. The total electric field $E = E_L + E_S$ satisfies the driven wave equation

$$\frac{\partial^2 E}{\partial z^2} - \frac{\alpha n_0}{2c}\frac{\partial E}{\partial t} - \frac{n_0^2}{c^2}\frac{\partial^2 E}{\partial t^2} = \frac{4\pi}{c^2}\frac{\partial^2 P^{NL}}{\partial t^2} \tag{3}$$

where α is the linear intensity absorption coefficient. The nonlinear polarization P^{NL} appropriate to the SBS process is given by

$$P^{NL}(z, t) = \frac{1}{4\pi}\left(\frac{\partial \epsilon}{\partial \rho}\right)_T \rho(z, t)\, E(z, t) \tag{4}$$

where $\rho(z, t)$ is the deviation of the material density from its equilibrium value ρ_0 and ϵ is the dielectric constant.

SBS is the coherent scattering of light from sound waves. The sound waves are described by the wave equation for the density variation of the material. This equation is obtained from the linearized hydrodynamic equations of motion for a viscous liquid, and has the following form [18], [19]:

$$\frac{v^2}{\xi}\frac{\partial^2 \rho}{\partial z^2} + \frac{\eta}{\rho_0}\frac{\partial^2}{\partial z^2}\frac{\partial \rho}{\partial t} - \frac{\partial^2 \rho}{\partial t^2} = \frac{\gamma^e}{8\pi}\frac{\partial^2 E^2}{\partial z^2} - v^2\rho_0\frac{\beta}{\xi}\frac{\partial^2 T}{\partial z^2}. \tag{5}$$

Here v is the velocity of sound, ξ is the ratio of the specific heat at constant pressure C_P to that at constant volume C_V, $\gamma^e = \rho_0(\partial\epsilon/\partial\rho)_T$ is the electrostrictive constant, $\beta = -(\partial\rho/\partial T)_P/\rho$ is the coefficient of thermal expansion at constant pressure, and $\eta = \eta_B + 4\eta_S/3$ is the viscosity where η_B is the bulk viscosity and η_S is the shear viscosity. The first driving term in this equation is due to the pressure gradient induced by the spatially modulated total optical intensity through the process of electrostriction. Note that the laser and Stokes fields are coupled together because this driving term involves E^2, and hence leads to a response at the beat frequencies between the two waves. The second driving term is due to the inhomogeneous temperature distribution produced by linear absorption. This temperature distribution is given by the energy transport equation

$$\rho_0 C_V\frac{\partial T}{\partial t} - \lambda_T\frac{\partial^2 T}{\partial z^2} - C_V\frac{\xi - 1}{\beta}\frac{\partial \rho}{\partial t} = \frac{n_0 c\alpha}{4\pi}E^2 \tag{6}$$

where λ_T is the coefficient of heat conduction. The driving term is the power per unit volume deposited in the medium by absorption of the spatially modulated total optical intensity. The solution to (5) and (6) for the fields given by (1) and (2) is

$$\rho(z, t) = \frac{1}{16\pi}\sum_{n,m} E_{S,m}(z)\, E^*_{L,n}(z) \cdot \frac{(\gamma^e + i\gamma^a_{nm})q^2_{nm}}{(q_{nm}v)^2 - \Omega^2_{nm} + 2i\Gamma_{nm}\Omega_{nm}} \cdot e^{i(\Omega_{nm}t - q_{nm}z)} + \text{c.c.} \tag{7}$$

where $\Omega_{nm} = \omega_n - \nu_m$ is the frequency of the acoustic wave driven by the beating between laser mode n and Stokes mode m, $q_{nm} = n_0(\omega_n + \nu_m)/c$ is the propagation constant of this disturbance, $\Gamma_{nm} = (\eta/2\rho_0)q^2_{nm}$ is the acoustic damping rate or equivalently the Brillouin linewidth, and $\gamma^a_{nm} = \alpha v^2 c\beta/2n_0 C_P\Omega_{nm}$ is the absorptive coupling constant. In deriving this equation, we have made the standard assumptions that the phonons are strongly damped and that the slowly varying amplitude approximation is valid. We have also ignored terms that give rise to coupling due to stimulated Rayleigh scattering. The amplitude of the density wave induced by the beating of laser mode n and Stokes mode m is maximum when $\Omega^2_{nm} = (q_{nm}v)^2$ or, equivalently, when

$$\omega_n - \nu_m = \pm\frac{2vn_0\omega_n}{c} \tag{8}$$

where the displayed form follows from the fact that the speed of sound is much smaller than the speed of light. The width of the resonance in (7) is Γ_{nm}. If the total laser bandwidth Γ_L obeys the inequality

$$\Gamma_L \ll \Gamma_{nm}\frac{c}{2vn_0}, \tag{9}$$

the variation of the right-hand side of (8) with ω_n is small compared to the width Γ_{nm} of the resonance, and hence ω_n can be replaced on the right-hand side by the center frequency of the laser ω_0. In this case, the Brillouin frequency shift $\Omega = 2vn_0\omega_n/c$ is the same for all laser modes. We shall henceforth assume that the Stokes frequencies appearing in the expansion (2) are chosen so that for each laser mode n having frequency ω_n, there is a Stokes mode that we shall also denote with the index n and which has frequency $\nu_n = \omega_n - \Omega + \Delta\nu$. The case $\Delta\nu = 0$ corresponds to exact resonance. We ignore the possibility that some other laser mode has frequency $2\nu_n - \omega_n$, and hence can act as a Stokes mode to Stokes mode n. We further assume that the laser mode spacing is much larger than the Brillouin linewidth Γ_{nn} and that the detuning $\Delta\nu$ is much smaller than the Brillouin frequency shift Ω. Under these conditions, the only terms that contribute apprecia-

bly to the double sum over n and m appearing in (7) are those for which $n = m$, and therefore the double sum collapses to the single sum

$$\rho(z, t) = \frac{(\gamma^e + i\gamma^a)^q}{32\pi v(\Delta\nu + i\Gamma)} \sum_n E_{S,n}(z)\, E^*_{L,n}(z) \cdot e^{i((\Omega - \Delta\nu)t - q_{nn}z)} + \text{c.c.} \tag{10}$$

where the Brillouin linewidth Γ and the absorptive constant γ^a now are the same for all laser modes. Through use of (4) and this form of the density variation, the nonlinear polarization driving the laser and Stokes fields become

$$P^{NL}_L(z, t) = \frac{i\omega n_0 \gamma^e(\gamma^e - i\gamma^a)}{128\pi^2\Gamma v c \rho_0(1 + i\Delta\nu/\Gamma)} \cdot \sum_n E_{S,n}(z)\, e^{-i\omega_n(t - zn_0/c)} \times \sum_k E_{L,k}(z)\, E^*_{S,k}(z)\, e^{2in_0 z(\omega_k - \omega_n)/c} + \text{c.c.} \tag{11}$$

and

$$P^{NL}_S(z, t) = -\frac{i\omega n_0 \gamma^e(\gamma^e + i\gamma^a)}{128\pi^2\Gamma v c \rho_0(1 - i\Delta\nu/\Gamma)} \cdot \sum_n E_{L,n}(z)\, e^{-i(\omega_n - \Omega)(t - zn_0/c)} \times \sum_k E^*_{L,k}(z)\, E_{S,k}(z)\, e^{-2in_0 z(\omega_k - \omega_n)/c} + \text{c.c.} \tag{12}$$

By inserting these expressions for the nonlinear polarization into the wave equation (3) for the electric field and using the slowly varying amplitude approximation, we obtain the following expressions for the spatial evolution of the laser and Stokes waves:

$$\frac{\partial E_{L,n}(z)}{\partial z} = -\frac{\gamma^e\omega^2(\gamma^e - i\gamma^a)\, E_{S,n}(z)}{32\pi\Gamma v c^2 \rho_0(1 + i\Delta\nu/\Gamma)} \cdot \sum_k E_{L,k}(z)\, E^*_{S,k}(z)\, e^{2in_0 z(\omega_k - \omega_n)/c} - \frac{\alpha}{2} E_{L,n}(z) \tag{13}$$

and

$$\frac{\partial E_{S,n}(z)}{\partial z} = -\frac{\gamma^e\omega^2(\gamma^e + i\gamma^a)\, E_{L,n}(z)}{32\pi\Gamma v c^2 \rho_0(1 - i\Delta\nu/\Gamma)} \cdot \sum_k E^*_{L,k}(z)\, E_{S,k}(z)\, e^{-2in_0 z(\omega_k - \omega_n)/c} + \frac{\alpha}{2} E_{S,n}(z). \tag{14}$$

We note that the growth of Stokes mode n is due to the coherent scattering of laser mode n off the density variation induced by the beating of laser mode k with Stokes mode k for all mode pairs k. For the mode pair with $k = n$, the exponential phase factor vanishes, implying that this process is perfectly phase matched. For all other mode pairs, there is a wavevector mismatch given by $2n_0(\omega_k - \omega_n)/c$. In order to find the conditions under which this wavevector mismatch can be neglected, we perform a change of variables in (13) and (14): $E_{S,n}(z) \to E_{S,n}(z) \exp(i\kappa_{S,n}z)$ and $E_{L,n}(z) \to E_{L,n}(z) \exp(i\kappa_{L,n}z)$. It is then easily seen that the phase factors $\kappa_{L,n}$ and $\kappa_{S,n}$ can be chosen in such a way that the coupled amplitude equations retain their original form, but with the complex exponentials replaced by unity provided that

$$\Lambda \ll \frac{c}{n_0\Gamma_L} \tag{15}$$

where Λ is the characteristic gain length of the SBS process, that is, the distance over which the amplitude of any of the waves varies appreciably. The right-hand side of this inequality is the coherence length of the laser.

We next derive the coupled intensity equations for the Stokes and laser fields. We multiply (14) by $E^*_{S,n}$, add the equation to its complex conjugate, and sum it over all Stokes modes n to obtain

$$\frac{\partial I_S(z)}{\partial z} = -gC(z)\, I_L(z)\, I_S(z) + \alpha I_S(z). \tag{16}$$

Similarly, from (13), we obtain

$$\frac{\partial I_L(z)}{\partial z} = -gC(z)\, I_L(z)\, I_S(z) - \alpha I_L(z) \tag{17}$$

where $I_L(z) = (n_0 c/8\pi)\, \Sigma |E_{L,n}(z)|^2$ and $I_S(z) = (n_0 c/8\pi)\, \Sigma |E_{S,n}(z)|^2$ are the intensities of the Stokes and laser waves,

$$C(z) = \frac{\left|\sum E^*_{L,n} E_{S,n}\right|^2}{\sum |E_{L,n}|^2 \sum |E_{S,n}|^2} \tag{18}$$

is the normalized correlation coefficient between the laser and Stokes fields, and g is the total gain coefficient given by

$$g = g_B + g_a \tag{19a}$$

where

$$g_B = \frac{\gamma^{e2}\omega^2}{2\Gamma n_0 v c^3 \rho_0} \frac{1}{1 + (\Delta\nu/\Gamma)^2} \tag{19b}$$

is that part of the gain coefficient which is due to electrostrictive coupling and

$$g_a = -\frac{\gamma^e\omega^2}{2\Gamma n_0 v c^3 \rho_0} \frac{\gamma^a \Delta\nu/\Gamma}{1 + (\Delta\nu/\Gamma)^2} \tag{19c}$$

is that part which is due to absorptive coupling. These gain coefficients are identical to those that apply for the case of a single-mode pump laser [4].

From the definition of $C(z)$ and (13) and (14), the differential equation that governs the spatial evolution of $C(z)$ is found to be

$$\frac{\partial C(z)}{\partial z} = -g[I_S(z) + I_L(z)]\, C(z)\, [1 - C(z)]. \tag{20}$$

This equation has two spatially invariant solutions given by $C(z) = 0$ and $C(z) = 1$. However, the solution $C(z) = 0$ can be shown by means of a linear stability analysis to be unstable. For the case of a Brillouin amplifier in which both the laser and the Stokes fields are generated externally, the boundary value of $C(z)$ depends on the boundary values of these fields, and in general is less than unity. However, due to the nonlinear coupling, the relative phases of the various modes change as the beams propagate, and (20) shows that within a few gain path lengths, the correlation coefficient $C(z)$ approaches the stationary value unity. For the case of a Brillouin generator in which the Stokes field is created entirely within the scattering medium, the amplitudes and phases of the Stokes modes are determined by the characteristics of the density fluctuations that initiate the SBS process. Since all laser modes scatter from the same density fluctuations, their relative amplitudes and phases are preserved in the scattered Stokes field, and hence even initially C is equal to unity. For C equal to unity, the coupled intensity equations (16) and (17) are exactly the same as those that apply for the case of a single-mode pump laser. We thus see that under the conditions assumed in this calculation, not only is the gain the same for a single- or multimode laser, but in addition, the reflectivity of an SBS generator is the same because the equations describing the spatial evolution of the fields are the same and depend only on the total intensities. The characteristic gain length introduced in inequality (15) can then be shown from (16)–(19) to be given by $\Lambda = 2/g(I_L + I_S)$, and the condition (15) on the total laser bandwidth becomes

$$\Gamma_L \ll \frac{c}{2n_0} g(I_L + I_S). \tag{21}$$

The conditions that must be satisfied in order for the SBS process to be independent of the pump laser mode structure can be summarized as follows. We have made several approximations which are standard approximations in the theory of SBS with a single-mode pump laser and which for most physical systems do not restrict the limits of validity of our calculation. These approximations are the use of the slowly varying amplitude approximation, the assumption that the phonons are strongly damped, and the use of the equations of hydrodynamics in their linearized form. In order to extend the theory to apply to the case of a multimode pump laser, we have had to make several additional assumptions. We have assumed that the laser bandwidth satisfies the inequality (9) and have ignored the dispersion of the refractive index over this bandwidth; for most physical systems, these conditions are very easily satisfied. We have also had to assume that no two laser modes are separated by precisely twice the Brillouin frequency shift in order to dismiss the possibility that the Stokes mode associated with one laser mode could act as an anti-Stokes mode for a different laser mode. In addition, we have assumed that the system is in steady state in the sense that the mode amplitudes in (1) and (2) depend only on the spatial variable z and not on time. We have also assumed that the laser mode spacing is much larger than the Brillouin linewidth and that the laser coherence length is substantially larger than the gain length as stated in inequality (15).

The physical interpretation of these last two conditions is as follows. The condition that that laser mode spacing must be larger than the Brillouin linewidth implies that the medium cannot respond to the beat frequency between laser mode n and Stokes mode m for $n \neq m$, and hence implies that the acoustic waves are driven only by the beating between laser mode n and its associated Stokes mode n. The coupling between different mode pairs arises from the fact that they scatter from each other's acoustic waves. This scattering of one mode pair from the acoustic wave formed by another mode pair is not necessarily phase matched. The condition that the laser coherence length substantially exceed the gain length ensures that this phase mismatch can be neglected. Note that the gain length Λ and not the total interaction distance L enters into inequality (15) because the nonlinear coupling modifies the propagation vectors in such a manner as to decrease the effective phase mismatch.

Experiment

In order to test these theoretical predictions, we have measured the threshold and the reflectivity under conditions of strong pump depletion for an SBS generator using the liquids methanol, acetone, and carbon disulfide. The experimental setup is shown schematically in Fig. 1. The frequency-doubled output from a passively Q-switched Nd:YAG laser (Quantel YG471-C) is focused into a cell containing the Brillouin active liquid. The pulse width is 15 ns (FWHM), the beam divergence is less than 0.5 mrad, and the beam diameter is approximately 7 mm. The focusing lens is planoconvex, has a focal length of 3.7 cm in air, and serves as the entrance window to the SBS cell. The energies of the laser and SBS pulses are measured with photodetectors $D1$ and $D2$, respectively, and their spectra are measured by imaging their Fabry–Perot interferograms onto a vidicon. By misaligning the laser oscillator, it can be made to run on up to four modes separated by 260 MHz, and up to three modes separated by 5 GHz. The spectral width of each individual laser mode was less than 50 MHz. The pulse repetition rate of the laser is 10 pps, and for each pulse, the outputs from the photodetectors and the interferograms are digitized and stored by a computer.

Typical Fabry–Perot interferograms for the laser and SBS return from methanol are shown in Fig. 2 both for the case when the laser is running single mode and when it is running multimode with a mode spacing of 5 GHz. The relative mode amplitudes in the SBS return are seen to be the same as those of the pump laser. This result is in agreement with the theoretical predictions that the correlation coefficient $C(z)$ is equal to unity, and hence that the spectrum of the pump laser is preserved in the scattered light. In Fig. 3, the SBS reflectivity is plotted as a function of laser pulse energy for a given number of modes

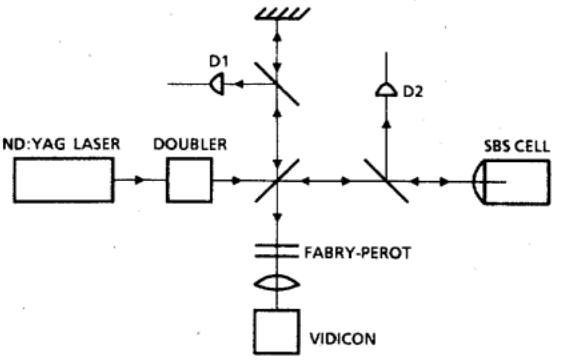

Fig. 1. Schematic of the experimental setup. The energies of the laser and SBS pulses are measured with detectors $D1$ and $D2$, and their spectra are measured by imaging their Fabry–Perot interferograms onto a vidicon. Not shown is an arrangement of lenses, beam splitters, and knife edges that images the interferogram of the laser onto one half of the vidicon and that of the laser onto the other.

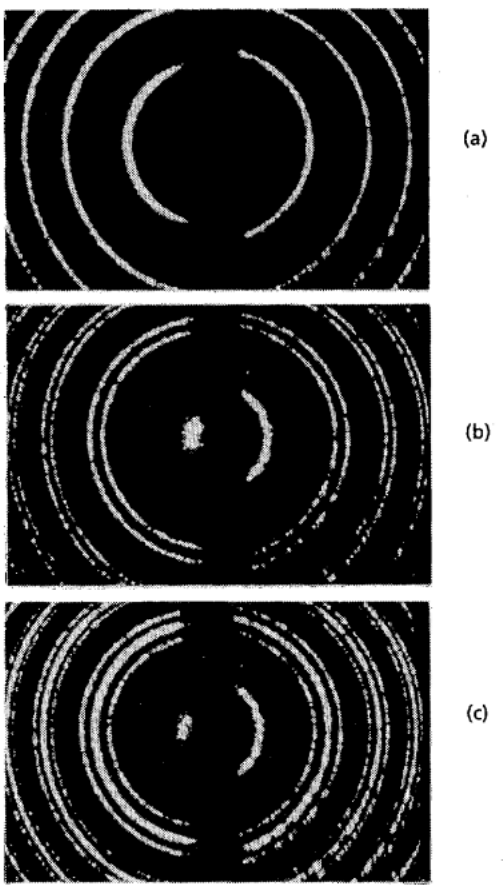

Fig. 2. Fabry–Perot interferograms for the laser and SBS return from methanol for the laser running (a) single mode, (b) on two modes separated by 5 GHz, and (c) on three modes separated by 5 GHz. In each case, the interferogram of the linear pulse is shown on the right-hand side and that of the SBS return on the left-hand side of the picture. The free spectral range of the interferometer is 30 GHz.

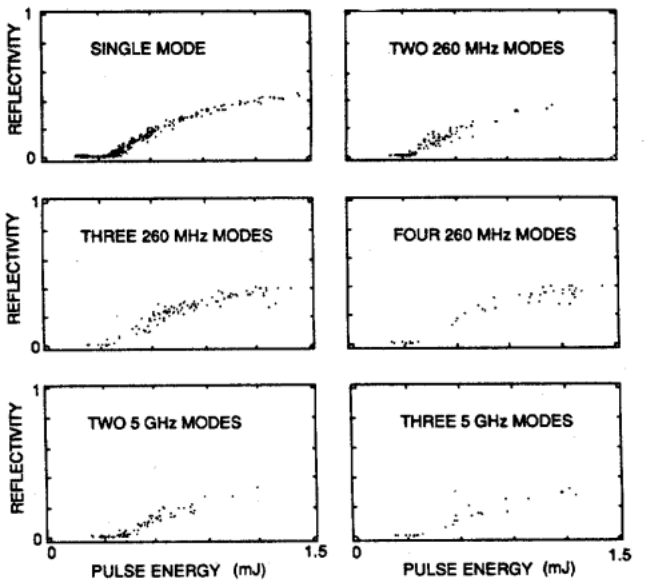

Fig. 3. SBS reflectivity of methanol plotted as a function of the laser pulse energy. Each graph corresponds to a different mode structure of the exciting laser.

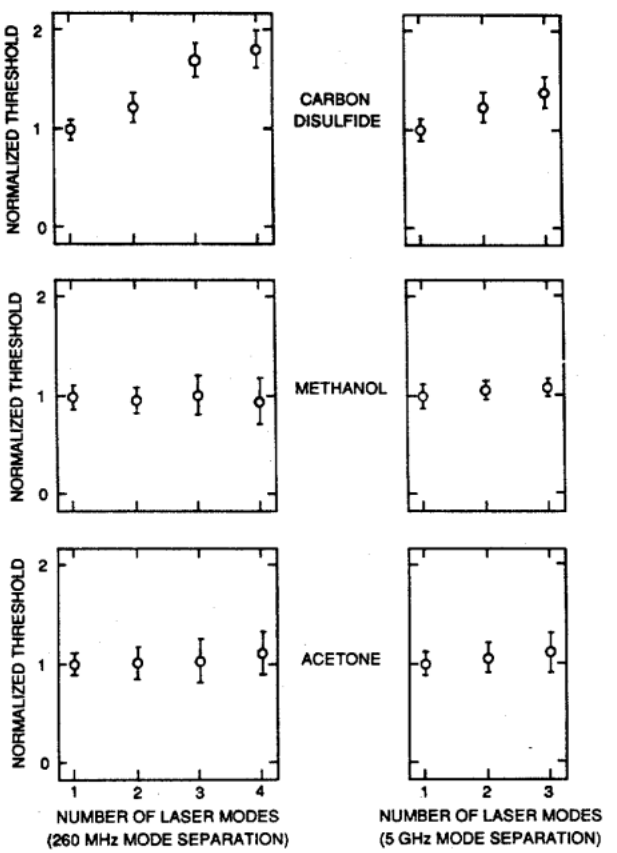

Fig. 4. Threshold pump energy for SBS normalized by its value for single-mode operation of the pump laser plotted as a function of the number of laser modes for three different liquids.

in the pump laser for methanol as the SBS-active medium; similar plots have also been obtained for acetone and carbon disulfide. Approximately 1400 laser pulses are collected and subsequently sorted by laser mode structure for each of these liquids. The threshold and the reflectivity at a laser pulse energy of 10 mJ are determined from these graphs and are plotted in Figs. 4 and 5 as a function of the number of laser modes. For each liquid, the threshold is normalized to its single-mode value. The bars indicate the pulse-to-pulse variations in the experiment. These variations are due to statistical variations in the SBS process, and are much larger than the variations resulting from measurement noise.

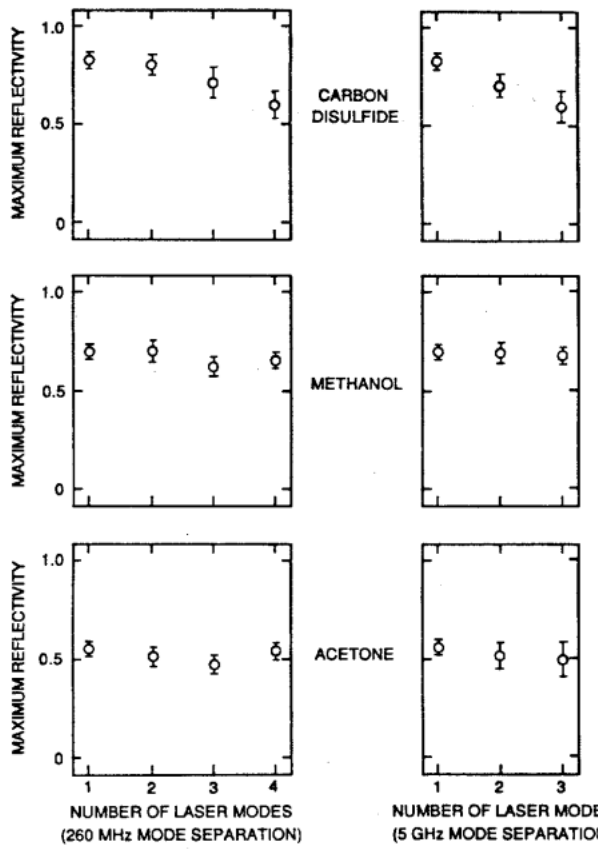

Fig. 5. SBS reflectivity at a pump-laser pulse energy of 10 mJ plotted as a function of the number of modes in the pump laser for three different liquids.

TABLE I
PHYSICAL PROPERTIES AT A WAVELENGTH OF 0.53 μm FOR THE THREE BRILLOUIN-ACTIVE LIQUIDS STUDIED EXPERIMENTALLY

	$\Omega/2\pi$ (GHz)	g_B^{max} (cm/GW)	g_a^{max}/α (cm^2/GW)	$\Gamma/2\pi$ (MHz)	g_R (cm/GW)	$\Gamma_R/2\pi$ (GHz)
Carbon Disulfide	7.7	130	162	60	31	7
Acetone	6.0	20	17	160	1.2	260
Methanol	5.6	13	10	210	0.5	270

Note: Ω is the Brillouin frequency shift, g_B^{max} is the gain coefficient for stimulated Brillouin scattering at $\Delta\nu = 0$, g_a^{max}/α is the gain coefficient for stimulated thermal Brillouin scattering at $\Delta\nu = \Gamma$ divided by the linear absorption coefficient α, Γ is the Brillouin linewidth, g_R is the gain coefficient for stimulated Raman scattering, and Γ_R is the Raman linewidth.

In order to compare these experimental results to the theoretical prediction, we first check the validity of the assumptions made in the theoretical section. Table I presents relevant physical parameters for the liquids used in the experiment [4], [20]. For all three liquids, the measured absorption coefficient at 0.53 μm is less than 0.01 cm^{-1}. Table I can then be used to infer that in each case, $g_B >> g_a$ or, equivalently, that the electrostrictive coupling is much greater than the absorptive coupling. Hence, according to (19), the total gain is maximum when the detuning $\Delta\nu$ from the Brillouin resonance is equal to zero. It is easily seen from Table I and the facts that the maximum laser bandwidth was 10 GHz and the confocal parameter was approximately 0.5 mm that the requirements of inequalities (9) and (15) are satisfied. The requirement that the laser mode spacing be greater than the Brillouin linewidth is seen to be well satisfied for the case of a 5 GHz mode spacing. For the case of a 260 MHz mode spacing, this requirement is close to not being satisfied, in particular for methanol where the Brillouin linewidth $\Gamma_B/2\pi$ is 210 MHz. However, our experimental results show no dependence on the mode spacing. Finally, our theory assumes that the SBS process is in steady state. This assumption is questionable for carbon disulfide where the phonon lifetime $\frac{1}{2}\Gamma_B$ of 1.4 ns is comparable to the laser pulse length of 15 ns.

As shown above, our theory predicts that there should be no dependence of the SBS gain or reflectivity on the number of laser modes. From Figs. 4 and 5, we see that there is no dependence on the pump-laser mode structure in either the threshold or the reflectivity for the case of ethanol or methanol as the SBS active medium. However, for carbon disulfide, the threshold increases and the reflectivity decreases as the number of modes in the pump laser increases. We believe that this dependence is due to competition between SBS and other nonlinear optical processes such as stimulated Raman scattering (SRS) and self focusing. As seen from Table I, the ratio of the gain for SRS to that of SBS is considerably larger for carbon disulfide than for the other two liquids. In addition, the SBS process is in the transient regime for much of the duration of the laser pulse, and hence the SBS gain is less than its steady-state value. Moreover, the SRS process has an extremely short response time of 11 ps, and hence can respond to the peaks of the intensity fluctuations that occur due to mode beating when the laser is operated multimode. Since the nonlinearity giving rise to self focusing is characterized by a short response time, competition due to this process is also sensitive to the mode structure of the laser.

CONCLUSIONS

We have shown theoretically that under conditions that are often met in practice, the characteristics of the SBS process are independent of the number of longitudinal modes present in the pump laser. We have performed an experiment that verifies these predictions for the threshold and reflectivity of SBS generation in the liquids acetone and methanol for which the assumed conditions are met. For the case of carbon disulfide, which does not fulfill these requirements, we find a dependence on the mode structure of the pump laser.

ACKNOWLEDGMENT

We gratefully acknowledge useful discussions with G. Dube, K. Lee, and M. G. Raymer.

REFERENCES

[1] B. Ya. Zel'dovich, V. I. Popovichev, V. V. Ragul'skii, and F. S. Faizullov, "Connection between the wave fronts of the reflected and exciting light in stimulated Mandel'shtam-Brillouin scattering," *Sov. Phys. JETP Lett.*, vol. 15, pp. 109–113, 1972.

[2] D. T. Hon, "Pulse-compression by stimulated Brillouin scattering," *Opt. Lett.*, vol. 5, pp. 516–518, 1980.
[3] M. Valley, R. Aprahamian, G. Linford, and G. Lombardi, "Beam-combining by stimulated Brillouin scattering," in *CLEO '85, Dig. Tech. Papers*, 1985, p. 85.
[4] W. Kaiser and M. Maier, "Stimulated Rayleigh, Brillouin, and Raman spectroscopy," in *Laser Handbook, Vol. 2*, F. T. Arrechi and E. O. Schulz-Dubois, Eds. Amsterdam: North-Holland, 1972. ch. E2, pp. 1078–1150.
[5] W. Koechner, *Solid State Laser Engineering*. Berlin: Springer-Verlag, 1976.
[6] D. C. Brown, *High Peak Power Nd:glass Laser Systems*. Berlin: Springer-Verlag, 1981.
[7] R. Giles, R. Fedosejevs, and A. A. Offenberger, "Reduced Brillouin scattering from multiline CO_2 laser interaction with a plasma," *Phys. Rev. A*, vol. 26, pp. 1113–1115, 1982.
[8] C. Montes, "Reflection limitation by driven stimulated Brillouin rescattering and finite-bandwidth spectral interaction," *Phys. Rev. A*, vol. 31, pp. 2366–2374, 1985.
[9] G. P. Dzhotyan, Yu. E. D'yakov, I. G. Zubarev, A. B. Mironov, and S. I. Mikhaliov, "Amplification during stimulated Raman scattering in a nonmonochromatic pump field," *Sov. Phys. JETP*, vol. 46, pp. 431–435, 1977.
[10] M. G. Raymer, J. Mostowski, and J. L. Carlsten, "Theory of stimulated Raman scattering with broad-band lasers," *Phys. Rev. A*, vol. 19, pp. 2304–2316, 1979.
[11] J. Eggleston and R. L. Byer, "Steady-state stimulated Raman scattering by a multimode laser," *IEEE J. Quantum Electron.*, vol. QE-16, pp. 850–853, 1980.
[12] A. Flusberg, D. Korff, and C. Duzy, "The effect of weak dispersion on stimulated Raman scattering," *IEEE J. Quantum Electron.*, vol. QE-21, pp. 232–236, 1985.
[13] Yu. E. D'yakov, "Excitation of stimulated light scattering by broad-spectrum pumping," *JETP Lett.*, vol. 11, pp. 243–246, 1970.
[14] V. I. Popovichev, V. V. Ragul'skii, and F. S. Faizullov, "Stimulated Mandel'shtam-Brillouin scattering excited by radiation with a broad spectrum," *JETP Lett.*, vol. 19, pp. 196–198, 1974.
[15] I. G. Zubarev and S. I. Mikhailov, "Stimulated scattering of light in a noisy pumping field with a spectrum wider than the frequency shift of the Stokes component," *Sov. J. Quantum Electron.*, vol. 4, pp. 683–684, 1974.
[16] B. Ya. Zel'dovich, and V. V. Shkunov, "Influence of the group velocity mismatch on reproduction of the pump spectrum under stimulated scattering conditions," *Sov. J. Quantum Electron.*, vol. 8, pp. 1505–1506, 1978.
[17] M. D. Skeldon, P. Narum, and R. W. Boyd, "Stimulated Brillouin scattering with a multimode laser," in *SPIE Conf. Proc.*, vol. 613, to be published.
[18] F. V. Hunt, "Propagation of sound in fluids," in *AIP Handbook*, D. E. Gray, Ed. New York: McGraw-Hill, 1957, pp. 3-37–3-57.
[19] R. D. Mountain, "Spectral distribution of scattered light in a simple fluid," *Rev. Mod. Phys.*, vol. 38, pp. 205–214, 1966.
[20] A. Penzkofer, A. Lauberau, and W. Kaiser, "High intensity Raman interactions," *Prog. Quantum Electron.*, vol. 6, pp. 55–140, 1979.

Paul Narum was born in Toten, Norway, on March 9, 1951. He received the Sivilingnioer degree in electrical engineering from The Norwegian Institute of Technology in 1976.

From 1976 to 1979 and from 1981 to 1985 he was at the Norwegian Defence Research Establishment developing military E-O systems. In 1980 he was a Visiting Research Associate at The Institute of Optics, University of Rochester, Rochester, NY, and since 1985 he has been a Research Associate at The Institute of Optics. His present research interests include phase conjugation by stimulated scattering and four-wave mixing, quantum optics, and new materials for nonlinear optics.

Mark D. Skeldon was born in Buffalo, NY, on May 31, 1949. After serving in the United States Navy he received the B.S. degree in mathematics in 1975, the M.S. degree in physics in 1977, both from Old Dominion University, Norfolk, VA, and the M.S. degree in optics in 1983, from the University of Rochester, Rochester, NY.

He became active in laser research at the Night Vision Laboratory, where he worked until 1982. Presently, he is pursuing the Ph.D. degree in optics at the University of Rochester. His research interests include nonlinear optics, phase conjugation, and stimulated scattering.

Robert W. Boyd, for a photograph and biography, see p. 1101 of the July 1986 issue of this JOURNAL.