

Stimulated Brillouin scattering parasitics in large optical windows

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Received May 1, 1986; accepted March 18, 1987

The growth of optical radiation, scattered transverse to the pump axis by stimulated Brillouin scattering (SBS) in optical windows, is considered. Basic equations are presented, and an analytic expression that determines the parasitic buildup time is derived for a transverse SBS geometry. Losses suffered by the scattered optical radiation are included in a bulk-loss term. Calculations are performed for fused-silica windows and compared with a numerical model. This parasitic process may affect the design of laser systems that will generate multianosecond, multikilojoule, narrow-band pulses in the ultraviolet region.

Stimulated Brillouin scattering^{1,2} (SBS) is an allowed nonlinear process in all media and has previously³ been associated with damage processes in mirrors and windows. Scaling of UV excimer lasers for kilojoule, short-pulse (10 nsec), narrow-band ($<0.5 \text{ cm}^{-1}$) operation creates the potential for transverse SBS processes. The SBS-generated optical wave would be a parasitic wave that removes energy from the transmitted wave and simultaneously generates an intense acoustic wave, which could fracture large windows.

The SBS process is a three-wave mixing process involving two optical waves and an acoustic wave in the SBS medium. A compressional acoustic wave creates an index grating that scatters energy from one optical field to the other. The two optical fields form a traveling standing-wave pattern with high and low field intensity regions, which, in turn, drives the compression and rarefaction of the acoustic wave. As with all nonlinear mixing processes, energy and momentum must be conserved.

Conservation of momentum requires that the wave vectors for the three waves be coplanar; thus the phase-matching diagram of Fig. 1 can be used. Since the speed of light is many orders of magnitude larger than typical sound-wave velocities, then to a high degree of accuracy

$$\omega_a = 2\omega_p \frac{nv_a}{c} \sin(\theta/2), \quad (1a)$$

$$\omega_s = \omega_p - \omega_a, \quad (1b)$$

$$\theta_a = (\theta/2) - 90^\circ, \quad (1c)$$

where ω_p , ω_s , and ω_a are the radial frequency of the pump, Stokes, and acoustical waves, respectively; θ is the angle between the pump and Stokes wave vectors; θ_a is the angle between the pump and acoustical wave vectors; v_a is the speed of sound in the SBS medium; n is the optical index of refraction in the SBS medium; and c is speed of light.

The pump and Stokes electric fields are defined by

the slowly varying envelopes P and S through the relations

$$E_s = (\pi/c)^{1/2} S \exp[i(\omega_s t - K_s r)] + \text{c.c.}, \quad (2a)$$

$$E_p = (\pi/c)^{1/2} P \exp[i(\omega_p t - K_p r)] + \text{c.c.} \quad (2b)$$

The intensities of the Stokes and pump fields are $|S|^2/2$ and $|P|^2/2$, respectively. The equations of motion for SBS can be derived from Maxwell's equations and from the equations for continuous media, which include electrostrictive terms.² Invoking the slowly varying envelope approximation, the SBS equations are given by

$$\left(\frac{\delta}{\delta t} + \frac{c}{n} \frac{\delta}{\delta x_p} \right) P = -S\rho^* - \frac{\alpha_p}{2} P, \quad (3a)$$

$$\left(\frac{\delta}{\delta t} + \frac{c}{n} \frac{\delta}{\delta x_s} \right) S = P\rho - \frac{\alpha_s}{2} S, \quad (3b)$$

$$\left(2\tau_b \frac{\delta}{\delta t} + 1 \right) \rho = \frac{g_0 c}{4n} SP^*, \quad (3c)$$

where x_p and x_s are coordinates in the direction of propagation of the pump and Stokes waves, respectively; τ_b is the intensity decay time of the acoustic wave; g_0 is the steady-state gain coefficient for the Stokes wave; α_p (α_s) is the background loss for the pump (Stokes) intensity per unit time; and ρ is proportional to the complex amplitude of the acoustic wave. These equations are identical to those for liquids and sufficiently dense gases.¹

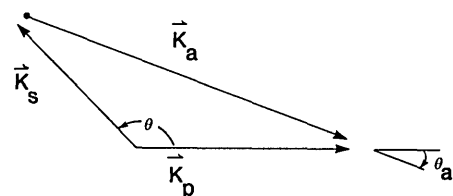


Fig. 1. Phase-matching diagram for SBS showing the pump, Stokes, and acoustic wave vectors.

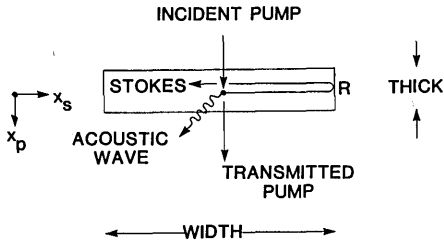


Fig. 2. Transverse SBS geometry. A uniform pump field in time and space is normally incident upon a window. A Stokes field propagating normal to the pump field builds up, reflecting off the sides of the windows with reflection coefficient R , and eventually depletes the pump field.

The angularly dependent gain coefficient in an isotropic or cubic solid for a monochromatic pump is given by^{5,6}

$$g_0 = \frac{2\pi n^7 P_{12}^2}{c\rho_0 \lambda_s^2 \nu_a \Delta\nu_b} \frac{1}{\sin \theta/2}, \quad (4)$$

where P_{12} is the elasto-optical coefficient, ρ_0 is the density of the solid, λ_s is the Stokes wavelength in vacuum, and $\Delta\nu_b$ is the FWHM linewidth of backscattered Stokes fluorescence ($\theta = 180^\circ$). For backwave SBS, $\sin \theta/2$ is 1, and the gain expression reduces to the form given by Ippen.⁶

The FWHM linewidth of the transverse Stokes wave ($\Delta\nu$) is a function of the acoustic frequency given by

$$\Delta\nu = \left(\frac{\omega_a}{\omega_{a0}} \right)^2 \Delta\nu_{b0}, \quad (5)$$

where ω_{a0} is the frequency of the acoustic wave where the backscattered linewidth $\Delta\nu_{b0}$ was measured. The pump wavelength and the scattering-angle dependence of the acoustic frequency (ω_a) are given in Eq. (1a). The decay time of the acoustic intensity is related to the line width by

$$1/\tau_b = 2\pi\Delta\nu. \quad (6)$$

Since the acoustic frequency is inversely proportional to the optical wavelength, the frequency dependence of the gain is determined solely by the wavelength variation of the optical index of refraction. The response time, however, varies as the square of the optical wavelength.

The gain for backwave SBS in fused-silica fibers has been experimentally measured⁶ at 535.5 nm to be 4.3 cm/GW. This is within experimental error of the theoretical value⁶ of 5.8 cm/GW. The index of refraction at 535 nm is 1.46, and at 248 nm it is 1.51. Including the effect of the index of refraction on linewidth, the theoretical gain for transverse SBS at 248 nm is 9.7 cm/GW. The frequency shift and response time^{7,8} (τ_b) for transverse SBS at 248 nm should be 4.9 GHz and 0.52 nsec ($\Delta\nu = 0.3$ GHz), respectively.

Consider the geometry in Fig. 2. The pump beam has a temporally flat top, uniform in the transverse dimensions, with an intensity I_{p0} . The leading edge of the pump beam is normally incident upon a window at time $t = 0$. Inside the window, a uniform Stokes wave with an initial intensity of I_{s0} is propagating at right angles to the pump wave. If the loss seen by the

Stokes wave at the window edges is distributed into an equivalent continuous loss coefficient (α_s), then the initial conditions, the boundary conditions, and the equations of motion become invariant under translation in the x_s direction. Under these conditions, the solution must also be invariant under the translation, and the partial derivatives with respect to x_s may be dropped in Eq. (3).

The equations can be further simplified by considering only the growth of the Stokes wave in the undepleted-pump regime, i.e., $I_p = I_{p0}$. In this case, Eq. 3(a) can be dropped and Eqs. 3(b) and 3(c) combined to yield

$$\left[\frac{\delta^2}{\delta t^2} + \left(\frac{1}{2\tau_b} + \frac{\alpha_s}{2} \right) \frac{\delta}{\delta t} + \left(\frac{\alpha_s}{4\tau_b} - \frac{g_0 I_{p0} c}{4\tau_b n} \right) \right] S = 0. \quad (7)$$

Assuming a solution of the form is $Ae^{\kappa t}$ and substituting into Eq. (7) results in a second-order equation for κ , whose roots are

$$\kappa_{\pm} = \pm \left[\left(\frac{1 - \alpha_s \tau_b}{4\tau_b} \right)^2 + \frac{g_0 I_{p0} c}{4\tau_b n} \right]^{1/2} - \frac{1 + \alpha_s \tau_b}{4\tau_b}. \quad (8)$$

The root with the plus corresponds to exponential growth of the Stokes electric field in the presence of a strong pump field.

When the gain in the window exceeds the loss, the transverse Stokes field will grow as a function of time. If this Stokes field has sufficient gain and sufficient interaction time, the Stokes wave will build up to a point where it can deplete significant energy from the pump radiation in the window. This causes the window to become opaque to additional pump radiation until the acoustic and optical waves decay away. Furthermore, the intense optical and acoustical fields may cause optical or mechanical damage.³ A threshold for this parasitic effect can be defined as the time required (T) for the Stokes power to reach a given fraction (f) of the pump power. Assuming an initial uniform (noise) source for the Stokes field, the solution for the Stokes intensity is exponential:

$$I_s(t) \approx I_s(0) \exp(2\kappa_+ t). \quad (9)$$

Using this form of the solution, the buildup time (T) is given by

$$T = -\frac{1}{2\kappa_+} \ln \left(f \frac{I_p t}{I_n w} \right), \quad (10)$$

where I_n is the noise intensity and $t(w)$ is the window thickness (width). Since significantly depleting the pump field corresponds to violating the assumptions that led to Eq. (8), f should be chosen to be 0.1 or less.

The major noise source for the SBS process is the thermal phonon population. However, as only the logarithm of the noise source intensity appears in Eq. (10), no significant modification of the results is expected from using the quantum noise at the Stokes frequency as the input noise source. The noise intensity was calculated as 0.0024 W/cm² for fused-silica windows pumped by KrF radiation. The acceptance angle of the noise is set by the allowed angular deviation before the SBS center frequency shifts more than the acceptance bandwidth.

Assuming an intensity reflectivity at the window

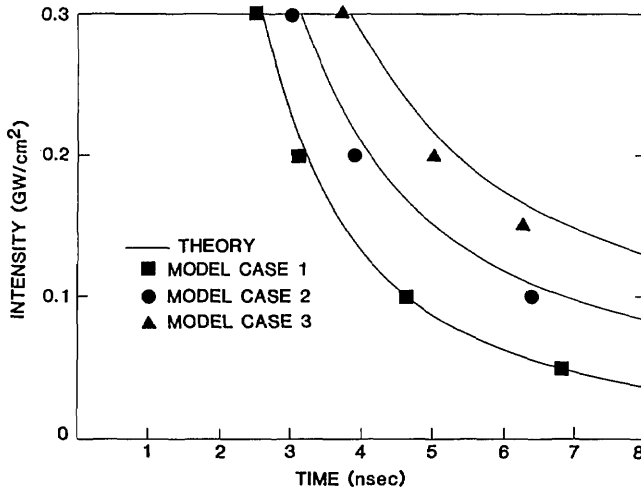


Fig. 3. Comparison of analytic results with a numerical model. Plotted is the buildup time ($f = 0.1$) to depletion threshold versus the KrF pump intensity in a uniformly filled fused-silica optical window. The first case is for a very wide window or one with perfectly reflecting edges. The second and third cases are for 40-cm- and 20-cm-wide windows, respectively, with a reflectivity of 0.001 at their edges. All curves assume a width-to-thickness ratio of 10.

edge of R , the distributed loss coefficient for the Stokes wave is given by

$$\alpha_s = -[c \ln(R)]/(nw). \quad (11)$$

This distributed loss coefficient should work well as long as the Stokes light makes a few trips across the window before threshold is reached or if the Stokes loss at the window edge is small.

A multidimensional computer model was also used to evaluate SBS parasitics in large-aperture windows. The computer model exactly solves the set of coupled partial differential equations [Eqs. (3)] as a function of time on a two-dimensional spatial mesh. There are five variables: the incident pump field, a pair of Stokes fields, and a pair of acoustic fields. The two Stokes waves propagate in opposite transverse directions to the pump wave, and there is a separate acoustic wave associated with each Stokes wave. The length-to-width ratio of the window was set at 1:10, and the edge reflectivity was varied.

The buildup time for $f = 0.1$ (10% depletion of the transmitted pump intensity) as computed with the analytic expression in Eq. (10) and with the exact computer model is plotted in Fig. 3. The three cases shown are for a very large window (or for a window with an edge reflectivity of $R = 1.0$) and for windows with widths 40 and 20 cm ($R = 0.001$). The agreement between the analytic expression and the exact solution is excellent, thereby validating the lumped-loss analytic model.

Figure 3 shows that very short pulses have high thresholds because of the finite response time of fused silica. It should be noted that this buildup time will be longer for longer-wavelength optical pump waves. Very long temporal pulses have an oscillation threshold approximation given by

$$I_p = (\alpha_s n)/(g_0 c). \quad (12)$$

For intensities below this threshold intensity, the SBS-driven Stokes wave will not build up. For pulses with pulse lengths in the 5–10-nsec regime, SBS could be a significant problem. Consider, for example, a 20 cm \times 20 cm window with a reflectivity of 0.001 at the edges. For a 5-nsec-long, narrow-band (<0.3 GHz) pulse from a KrF laser, 1 J/cm² of average pump radiation will be the upper fluence limit for avoiding this parasitic process. This corresponds to an upper limit on the total energy through the mirror of only 400 J. Exciting the transverse SBS parasitic will result in loss of transmission efficiency and probable damage to the window.

The equations presented here are based on monochromatic pumping. Including pump bandwidth will increase the threshold. The exact dependence of threshold with bandwidth depends on how the pump is modeled⁹ and is beyond the scope of this Letter. However, under the assumptions of Eq. (27) of Ref. 9, the broadband threshold intensity (I_{th}^b) is related to the calculated monochromatic threshold intensity (I_{th}^m) by

$$I_{th}^b = I_{th}^m + I_{cr}, \quad I_{cr} = (4\Delta\nu_p)/g, \quad (13)$$

where $\Delta\nu_p$ is the bandwidth of the pump, which is assumed to be much much greater than the scattering linewidth $\Delta\nu$.

In conclusion, a simple analytic expression has been developed to analyze transverse SBS geometries. This analytic form agrees well with the results of a separate exact computer model. SBS thresholds for narrow-band pump light are in the kilojoule regime for UV pulses of the order of 10 nsec. Two methods that potentially reduce this problem are to increase the laser linewidth to be much greater than the Brillouin linewidth and to use small edge-clad windows. Either of these measures will increase the threshold, but neither will eliminate the effect.

Funding for this research was provided by Los Alamos National Laboratories under contract no. 9-X65-W1478-1 and by Spectra Technology IR&D. We also acknowledge useful discussions with Norm Kurnit and Jay Ackerhalt as well as useful comments from a reviewer.

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