PHYSICAL REVIEW A VOLUME 42, NUMBER 5 1 SEPTEMBER 1990

Soliton decay and gain focusing in stimulated Raman scattering

Joseph W. Haus and Michael Scalora
Physics Department, Rensselaer Polytechnic Institute, Troy, New York 12180-3590
(Received 28 June 1990)

Solitons are generated in stimulated Raman scattering using a π phase shift of a weak Stokes input field; diffraction effects are retained in the calculation. For strong diffractive coupling, we observe an on-axis increase in both the laser and the Stokes intensities and a corresponding narrow transverse profile for both fields; these results are elucidated and explained.

Solitons have been observed in stimulated Raman scattering (SRS) when a π phase shift was introduced in the incident Stokes light.[1] The phase shift in the Stokes pulse leads to a temporary repletion of the pump beam and a consequent loss in the Stokes energy. This phenomenon was theoretically predicted for plane-wave propagation of the electromagnetic fields.[2] The phase shift (of order π) can also be induced spontaneously from quantum noise,[3] leading to the same phenomenon. This is a special example of the decay of an unstable state of which the transient buildup of laser oscillation and superfluorescence are other important examples in quantum optics.

The media in which the Raman process occurs are diverse and applications can be found in the generation of coherent light over a wide range of frequencies in the far-infrared region;[4] in the cleanup of fluctuations in light,[5] and in the compression of excimer laser pulses.[6] These are only a few areas of interest. Moreover, SRS has similarities with schemes used for developing x-ray lasers, whose radiation would be initiated from quantum noise.[7]

In this Rapid Communication we will study the propagation effects on soliton behavior with the aid of the semiclassical model for SRS. We opt to generate the soliton by inducing a π phase shift in the Stokes field (deterministic model), and use full three-dimensional propagation effects in order to study the role that diffractive coupling plays in such systems.

Under general conditions, assuming the system remains mostly in its ground state, we write the equations of motion for the following complex, slowly varying fields: the pump field E_L, the Stokes field E_S, and the polarization Q. The length of the medium is scaled to a unitless number and the coupling coefficients have been removed by scaling both the space and time coordinates and the polarization. The equations are then written as[3]

$$\frac{\partial}{\partial z}E_L+\frac{1}{v}\frac{\partial}{\partial t}E_L-\frac{i}{F}\nabla_\perp^2 E_L=-QE_S\,, \tag{1}$$

$$\frac{\partial}{\partial z}E_S+\frac{1}{v}\frac{\partial}{\partial t}E_S-\frac{i}{F}\nabla_\perp^2 E_S=Q^*E_L\,, \tag{2}$$

$$\frac{\partial}{\partial t}Q=-\gamma Q+E_L E_S^*\,, \tag{3}$$

where v is the group velocity, z is the longitudinal coordinate, $\nabla_\perp^2$ is the Laplacian that operates on the transverse coordinates (x,y); the strength of this diffractive coupling is measured by the Fresnel number $F=2kd^2/L$, which is assumed to be equal for both the Stokes and laser fields, d is a typical width of the laser field profile, k is the wave number, and L is a characteristic length of the medium. γ is a phenomenological damping factor for the polarization Q (we choose $\gamma=100$ in this Rapid Communication, which corresponds to a steady-state gain-length product of 50 in our numerical calculations below).

In order to study soliton behavior under the full propagation effects, we solve the set of Eqs. (1)-(3) using the method of characteristics, combined with the split operator method[8] in order to account for diffractive coupling. The details of our numerical method will be given elsewhere.[9] We assume that propagation takes place in the z direction and we include the effects of both transverse coordinates by using a two-dimensional fast Fourier transform (FFT). The initial and boundary conditions were given for all the fields; at the input boundary the laser field was taken with a Gaussian profile

$$E_L(0,t)=50\exp[-(x^2+y^2)]\,; \tag{4}$$

the Stokes field is $E_S(0,t)=E_L(0,t)/25$ for times less than 0.01 and the negative of this value for greater times. The polarization Q is initially assumed to be zero. When the polarization advances by the time increment Δt, it follows that the fields must have advanced a distance $\Delta z=v\Delta t$, and the effective propagation distance is then $\Delta\xi=v\Delta t\sqrt{2}$. We have modified a second-order approximation split-step operator method combined with a predictor-corrector method to efficiently solve these equations. The velocity is taken to be $v=40$ in our scaled units, but we have used values from 0.1 to 200 and found no qualitative changes in our results. The gain is reversed between the laser and Stokes energy at $t=0.01$ when the input Stokes field has its phase shifted by π. The laser intensity grows and repletes itself in a short pulse characteristic of solitons in nonlinear media.

The full three-dimensional propagation effects are accounted for by using a two-dimensional FFT and it was tested for convergence, giving excellent agreement with known results. As expected, we found that in this case, where no fluctuations play any role, radial symmetry is preserved. We used a 32×32 grid for the FFT transform, which was large enough to give accurate results without spurious interference from the periodic boundary conditions. We tested a 32 point and 64 point grid with one-dimensional FFT with no observable changes in the results. Furthermore, there is no qualitative difference between the one- and two-dimensional simulations.[9]

In Fig. 1 we show the on-axis laser intensity for a Fresnel number $F=15$. When the medium is short, no depletion is observed. For small distances the soliton is sharp, but it eventually broadens and dissipates due to diffraction, as seen at $z=1.2$. The soliton peak is actually larger than the input laser intensity at $z\cong0.6$. From Eq. (4), $|E_L|^2_{\max}=2500$ at the input ($z=0$), and we find the on-axis field at $z\cong0.6$ $|E_L|^2_{\max}\cong2800$. In Fig. 2 we plot the corresponding Stokes intensity. For short distances, there is a characteristic dip in the Stokes intensity when the laser intensity is repleted. However, at larger propagation distances, and especially at $z=1.2$, the Stokes intensity increases as the laser intensity is increased. The effect is not expected to be small, and for $F=15$ the Stokes intensity increases by more than a factor of 2 at $z=1.2$ ($|E_S|^2_{\max}\cong5600$). This is a surprising result in view of the fact that for infinite Fresnel number, the sum of the Stokes and laser intensity are constant and we have introduced diffraction into the equations of motion which would normally further diminish the on-axis intensity.

These figures can be understood as gain focusing of the Stokes and laser intensities. Our gain-focusing phenomenon is quite distinct from that discussed in earlier papers.[10,11] Those papers treated linearized equations where there is a balance between gain and diffraction. Our version of gain focusing occurs off-axis in a regime that can only be described through a full nonlinear analysis. In Figs. 3 and 4, we show the laser and Stokes fields as a function of time with only one transverse coordinate displayed at $z=1.2$. For short propagation times, the laser is depleted and conical emission of the field is observed. The reappearance of the laser field due to the π phase shift is strong and nearly complete on the axis; the transverse profile of the laser field is about a factor of four narrower than the input field. Satellite intensities are also predicted from the simulation.

In Fig. 4, the Stokes intensity initially grows, saturates and then spreads as it propagates. No conical emission is observed for this field. As the laser soliton reappears, the Stokes intensity collapses to the axis and itself becomes much narrower. Once the laser has again depleted, the Stokes intensity again spreads.

Diffraction in this system plays an, at first, unexpected

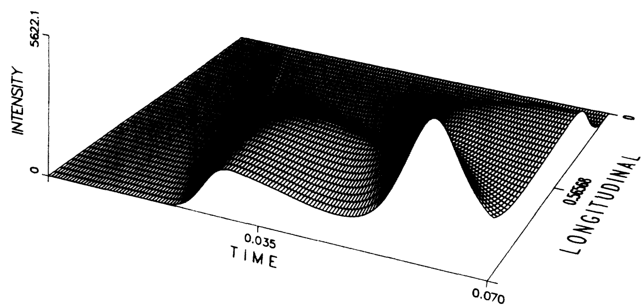

FIG. 2. On-axis Stokes intensity vs time and longitudinal (z) distance.

role, but the consequences can be understood with simple physical arguments. As the laser pump propagates inside the medium, the higher intensities of the input Gaussian shape are strongly absorbed, thus attaining the general shape of an annulus, and with only free space propagation a Poisson spot should appear on the axis. The annulus shape of the laser intensity lies in a region where the Stokes intensity is initially small. The laser propagates in a cone toward the center and outer portion of the beam.

Since the Stokes field is spreading by diffraction to the outer region, the laser propagates in its presence and is amplified with a gain that is comparable to the gain which the on-axis laser field experiences. The part of the laser field diffracting into the hole is amplified by the phase change in the Stokes field and this converging component is responsible for the much narrowed laser pulse. Concurrently the energy is being removed from the wings of the Stokes and focused onto the axis; the Stokes energy is in turn also redirected by the strong focusing field of the laser. Therefore, the presence of the side wings in the laser intensity will promote energy flow to the center of the beam. As this takes place, the Stokes gain, which is intensity dependent, pushes the center intensity up, thus making transverse coupling even more important. The effect, as outlined above, is then essentially initiated by a Poisson-spot circumstance.[12]

The integrated intensity over the transverse coordinates, shows a result that is expected from the plane-wave theory. Namely, as the integrated laser intensity is repleted, the integrated Stokes intensity decreases. The energy

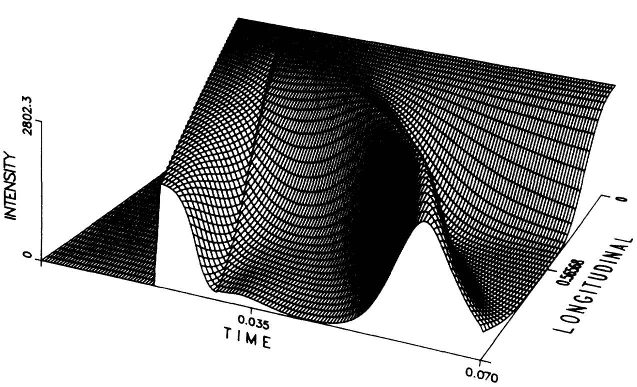

FIG. 1. On-axis laser intensity vs time and longitudinal (z) distance in the medium. The scaling of the coordinates and intensity makes them unitless.

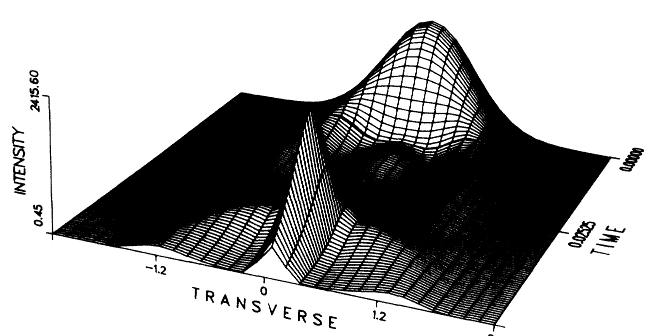

FIG. 3. The transverse profile of the laser field at $z=1$ vs time. $F=15$.

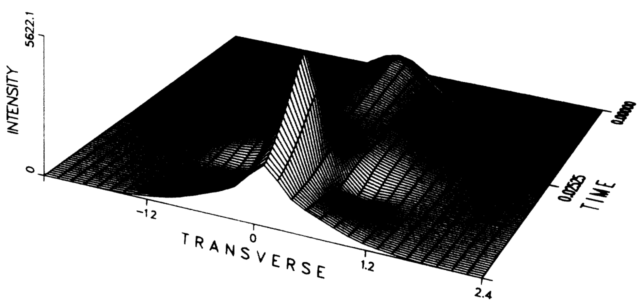

FIG. 4. The transverse profile of the Stokes intensity vs time.

is, of course, conserved, but a surprising redistribution occurs along the transverse direction due to diffraction. The close agreement between experiment and the plane-wave theory on the generation of spontaneous solitons[3] is due to the fact that the noise is important during the initiation process and only the integrated intensities were studied. It will be interesting to study the correlations in the transverse intensity profile which cannot be predicted by plane-wave theory.

We believe that our new predictions could be observed in carefully designed experiments. First, the conical emission in the laser, the narrowing and satellite peaks should be considered signatures of the on-axis effect. Second, intensity must be time resolved and our results apply to the near field. Finally, experiments that integrate over slices of the profile are not likely to observe the effect and therefore, either an array of detectors or a detector which resolves the on-axis intensity will be required.

In conclusion, we find that diffractive coupling also causes the solitons to dissipate, once they are created. Diffraction is responsible for an energy redistribution which causes an unexpected sharp gain in the on-axis Stokes intensity at the same time that the laser intensity is increased. The transverse profile of the soliton and the Stokes emission is narrowed by gain and diffraction coupling. We find that the process is initiated by conditions that strongly favor the onset of diffractive coupling, namely intensity discontinuity in the beam due to strong absorption in its center.

This research was partially supported by NSF Grant No. ECS-8813028. A grant of computer time at the Cornell National Supercomputer Facility and at Rensselaer Polytechnic Institute is gratefully acknowledged. Thanks are also due to M. Kuperschmidt and M. Miller for computer system help.

[1]K. Druhl, R. G. Wenzel, and J. L. Carlsten, Phys. Rev. Lett. **51**, 1171 (1983); R. G. Wenzel, J. L. Carlsten, and K. J. Druhl, J. Stat. Phys. **39**, 621 (1985); D. C. MacPherson, R. C. Swanson, and J. L. Carlsten, Phys. Rev. A **39**, 6078 (1989).

[2]F. Y. Chu and A. C. Scott, Phys. Rev. A **12**, 2060 (1975); D. J. Kaup, Physica **6D**, 142 (1983); H. Steudel, *ibid.* **6D**, 155 (1983).

[3]M. G. Raymer and J. Mostowski, Phys. Rev. A **24**, 1980 (1981); C. M. Bowden and J. C. Englund, Opt. Commun. **67**, 71 (1988); M. G. Raymer and I. A. Walmsley, Prog. Opt. (to be published).

[4]P. Rabinowitz, B. N. Perry, and N. Levinos, IEEE J. Quant. Electron. **QE-22**, 797 (1986).

[5]N. F. Andreev, V. L. Bespalov, A. M. Kiselev, and A. Pasmanik, Sov. J. Quantum Electron. **9**, 585 (1979).

[6]J. Murray, J. Goldhar, D. Eimerl, and A. Szoke, IEEE J. Quantum Electron. **QE-15**, 342 (1979).

[7]G. Hazak and A. Bar-Shalom, Phys. Rev. A **38**, 1300 (1988).

[8]M. Scalora and J. W. Haus, J. Opt. Soc. Am. B **6**, 1714 (1989).

[9]M. Scalora and J. W. Haus (unpublished); M. Scalora, Ph.D. thesis, Rensselaer Polytechnic Institute, 1990 (unpublished).

[10]H. Kogelnik, Appl. Opt. **4**, 1562 (1965); D. Kotter, D. C. Hanna, and R. Wyatt, Appl. Phys. **8**, 333 (1975).

[11]K. J. Druhl, S. Shakir, and M. Yousaf, Opt. Lett. **11**, 446 (1986).

[12]M. Born and E. Wolf, *Principles of Optics,* 6th ed. (Pergamon, London, 1980), p. 375.