DISSIPATIVE SUPERLUMINOUS BRILLOUIN SOLITONS IN AN OPTICAL-FIBER RING CAVITY

Carlos Montes, Eric Picholle, Jean Botineau
Olivier Legrand, and Claude Leycuras

Laboratoire de Physique de la Matière Condensée, (U.R.A. - C.N.R.S. N° 190)
Université de Nice-Sophia Antipolis, Parc Valrose, F-06034 Nice CEDEX, France.

Abstract : Generation of large-scale spatio-temporal coherent structures caused by stimulated Brillouin backscattering of a narrow-band laser wave in a large-gain one-dimensional nonlinear medium is studied by comparing the numerical simulations and the analytical asymptotics of the three-wave resonant model to actual experiments in a single-mode optical-fiber. This comparison recently allowed us to predict [1] and to perform the first experimental observation [2] of the "superluminous" Brillouin soliton backward propagating with respect to the cw pump in an optical-fiber ring-cavity.

I Introduction

Stimulated Brillouin scattering (SBS) is the dominant stimulated scattering process in many optical media, and particularly in optical fibers, pumped by single-frequency lasers [3] [4]. Considered as a detrimental effect for optical communications and above all for laser-plasma fusion experiments, where it is responsible of reflecting a large fraction of the laser energy, an important aim has been to limit its efficiency [5]. However, it allows also to achieve efficient pulse amplification and compression since 1968 [6].

The spatio-temporal SBS kinetics is well described by a three-wave nonlinear resonant interaction [7]. Our purpose has been to better understand the time-dependent SBS generation of large-scale coherent structures by comparing the numerical simulations and the analytical asymptotics of the three-wave coherent model to actual experiments in an optical-fiber.

In the one-dimensional problem, the forward-propagating pump wave (at frequency $\omega_p = k_p c/n_0$) couples with the thermal phonon fluctuations of the material medium (at frequency $\omega_a \simeq 2c_a\omega_p n_0/c$, where c_a is the acoustic velocity and n_0 the unperturbed refractive index) and stimulates a counterpropagating Stokes wave (at frequency $\omega_S = \omega_p - \omega_a$) downshifted by the acoustic frequency (ω_a). The acoustic wave is in turn amplified by electrostriction. The resonant condition for the three-wave coherent interaction ($\omega_p = \omega_S + \omega_a$) provides maximum power transfer when the wave-vector mismatch is zero ($\mathbf{k}_p = \mathbf{k}_S + \mathbf{k}_a$; $\Rightarrow k_a = k_p + k_S \simeq 2k_p$). Thus, assuming slowly varying envelopes for the waves, neglecting optical dispersion, and respectively denoting the complex amplitude E_p for the pump wave, E_S for the counterpropagating Stokes wave and E_a for the acoustic wave, the three-wave equations read as follows in a coherent dimensionless form [8] [9] :

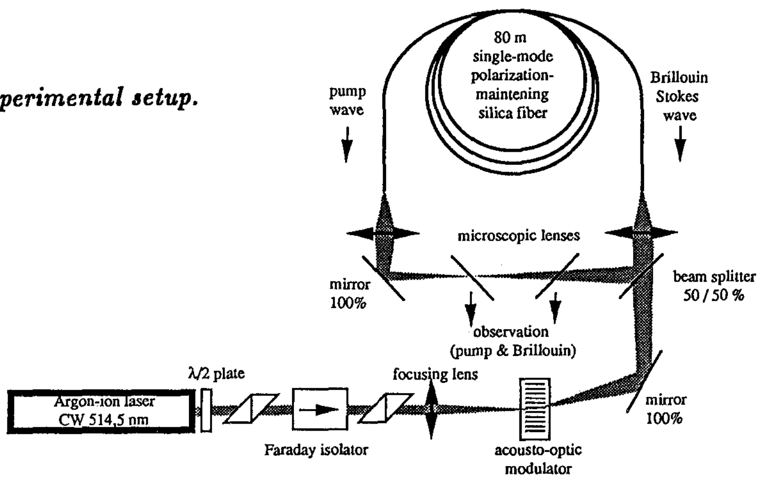
$$\begin{aligned}(\partial_t + \partial_x + \mu_e)E_p &= -E_S E_a \\ (\partial_t - \partial_x + \mu_e)E_S &= E_p E_a^* \\ (\partial_t + \mu_a)E_a &= E_p E_S^*,\end{aligned} \tag{1}$$

where the acoustic velocity has been neglected, due to its smallness relative to light velocity (normalized to unity), and where μ_e (μ_a) is the damping coefficient for the optical (acoustic) wave. Moreover, additive terms in Eqs.(1) allow our numerical simulations to account for phase modulation due to optical Kerr contribution [9].

We have carried out numerical and asymptotical studies of the coherent space-time dependent kinetics, with actual experiments in an optical-fiber ring-cavity, which lead us to consider two main SBS time dependent regimes : (i) Stokes pulse amplification and compression under nonstationary conditions [8] [9], and (ii) generation of a backscattered "superluminous" Stokes soliton accompanied by self-induced transparency for the pump [1] [2]. Let us mention two previous papers [10] [11], closely related to the theoretical and numerical models, though some confusion appears between regimes (i) and (ii).

The compression regime (i) takes place when two separated pump and Stokes pulses interact in a counter-propagating motion. The Stokes pulse behaves as a shock-wave, the amplification of which depleting the pump wave. The interaction amplifies the acoustic wave by electrostriction starting from the Stokes leading edge and propagating with the Stokes pulse in its backward traveling motion. The Stokes amplification strongly depends on the shape of its leading edge, but for long enough times it takes the shape of a "π-pulse" self-similar profile [12] whose leading maximum amplitude grows linearly in time while its width decreases as the inverse of time [8]. It is difficult to identify this asymptotic stage in an actual experiment, since we must control the initial Stokes pulse. A stable train of compressed Stokes pulses (to $\sim$ 10 ns) has been obtained in a stimulated Brillouin fiber ring laser (of length L = 83 m), similar to that shown in figure 1, but where the acousto-optic modulator (AOM) was put inside the ring-cavity,

Fig. 1 *Experimental setup.*



and the Ar-ion cw pump beam, at $\lambda_p = 514.5$ nm, was coupled into the fiber through the first Bragg order of the AOM [9]. By periodically interrupting the pump beam and the Stokes leading edge at the round-trip flight time $\Delta t_L = Ln_0/c \simeq 420$ ns, we periodically repeated the compression experience. The experimental time scale for the generated Stokes pulse width ($\simeq 10$ ns) is comparable to the spontaneous acoustic decay time in silica ($\simeq 7$ ns deduced from a spectral Brillouin bandwidth of $\Delta\nu = 150$ MHz at a pump wavelength $\lambda_p = 514.5$ nm) [13] and therefore absolutely calls for the use of the coherent description given by Eqs.(1). These studies have allowed us to better precise the experimental configuration for obtaining the second regime, which is the object of this paper.

II. Dissipative Brillouin soliton

Let us look for soliton solutions of Eqs.(1) which involve the coherent dynamics of a self-similar backward-propagating three-wave (pump-Stokes-acoustic) structure. For the three-wave interaction (1) in the nondissipative case (i.e. $\mu_e = \mu_a = 0$), specific traveling wave solutions have been studied in the forward-scattering case [14] - [16]. Corresponding solutions in the SBS case display new interesting features that persist in the dissipative case [10]. These solutions are special cases of a broader class of solitons previously described through inverse scattering transform [17] [18].

However, in the case of silica optical fibers, it is not possible to neglect the dampings ($\mu_i \neq 0$). Analytic traveling-wave solutions are still available if the pump attenuation is neglected, which is locally legitimate as long as $\mu_e/\mu_a \ll 1$; it is indeed the case in our experiment. Starting from Eqs.(1) we first perform the change of frame moving in the backscattered direction $x \longrightarrow x + vt,\ t \longrightarrow t$ which yields :

$$\begin{aligned} \left[\partial_t + (1+v)\partial_x\right]E_p &= -E_S E_a - \mu_e E_p \\ \left[\partial_t + (v-1)\partial_x\right]E_S &= E_p E_a^* - \mu_e E_S \\ \left[\partial_t + v\partial_x\right]E_a &= E_p E_S^* - \mu_a E_a. \end{aligned} \tag{2}$$

Then, by defining the A_i's fields as

$$A_1 = |1+v|^{1/2}E_p \ ; \quad A_2 = |v-1|^{1/2}E_S \ ; \quad A_3 = |v|^{1/2}E_a \ ; \tag{3}$$

and looking for stationary solutions in the new frame, we have :

$$\begin{aligned} \partial_X A_1 &= -s_1 A_2 A_3 - s_1\rho_1 A_1 \\ \partial_X A_2 &= s_2 A_1 A_3^* - s_2\rho_2 A_2 \\ \partial_X A_3 &= s_3 A_1 A_2^* - s_3\rho_3 A_3 \end{aligned} \tag{4}$$

where $X = x/|(v-1)(1+v)v|^{1/2}$; $\rho_1 = \mu_e|v-1|^{1/2}|v|^{1/2}/|1+v|^{1/2}$; $\rho_2 = \mu_e|1+v|^{1/2}|v|^{1/2}/|v-1|^{1/2}$; $\rho_3 = \mu_a|1+v|^{1/2}|v-1|^{1/2}/|v|^{1/2}$; and $s_1 = sgn(1+v)$, $s_2 = sgn(v-1)$, $s_3 = sgn(v)$. Taking into account stability arguments [18], we shall be concerned only with the solutions for which $v > 1$, i.e. a structure moving at a velocity larger than velocity of light (the $v < -1$ case only intervenes in a transient stage of the nonlinear interaction and is asymptotically unstable). Since $\mu_e \ll \mu_a$ and $v > 1$ we also have $\rho_1 \ll \rho_3$ but ρ_2 and ρ_3 are of the same order, i.e. $\mu_e \sim |v-1|\mu_a$. Thus, for $\rho_1 = 0$ and $s_1 = s_2 = s_3 = 1$, we obtain the following solution

$$A_1 = -A \tanh AX + A_0$$
$$A_2 = A_3 = A \text{ sech } AX \tag{5}$$

with $A_0 = \rho_2 = \rho_3$. Therefore

$$\frac{\mu_e}{\mu_a} = |\frac{v-1}{v}| \quad \Rightarrow \quad v = \frac{1}{1-\mu_e/\mu_a}. \tag{6}$$

From (3) we obtain the expressions for the initial amplitudes of the waves

$$E_p = P_0 - P \tanh\big[(x+vt)/\Delta\big],$$
$$E_S = S \text{ sech}\big[(x+vt)/\Delta\big], \tag{7}$$
$$E_a = SP\big[2/(S^2+P^2)\big]^{1/2} \text{ sech}\big[(x+vt)/\Delta\big],$$

with the extra relationships

$$P_0 = (\mu_e\mu_a)^{1/2}, \quad S/P = (2\mu_a/\mu_e - 1)^{1/2}, \tag{8}$$

the width of this backward complex structure being given by

$$\Delta = \frac{(\mu_e/\mu_a)^{1/2}}{P(1-\mu_e/\mu_a)}. \tag{9}$$

There is only one free parameter left which will be taken as the pump amplitude at the far left end of the structure shown in figure 2, namely $P_0 + P = E_p(x + vt \ll -\Delta) = 1$.

This superluminous self-similar Stokes pulse does not contradict by any means the special theory of relativity. Its motion can be viewed as the result of an amplification of the leading edges of the Stokes and acoustic pulses while, at the same time, their rears are depleted, the pump wave being partially restored after the interaction, in some cases with an opposite phase. This behavior is quite reminiscent of the self-induced transparency effect encountered in the coherent pulse propagation in a two-level medium [19]. No transportation of information can be obtained via this process which can only occur if a sufficiently extended background of Stokes light is available so that the superluminous self-similar amplification takes place. If at a given time (t=200 in fig.2) we perturb the soliton profile by "engraving" a signal, this one moves at the velocity of light (i.e. remains at rest in the backward moving frame of fig.2) and the Brillouin soliton runs away at the "superluminous" velocity. Figure 2 also shows the elasticity of the soliton structure. An additive interesting feature obtained by the numerical treatment of Eqs. (1) is that, when it is built, the Stokes pulse is no longer perturbed by pump phase fluctuations, which are compensated by the acoustic wave (cf. fig. 3).

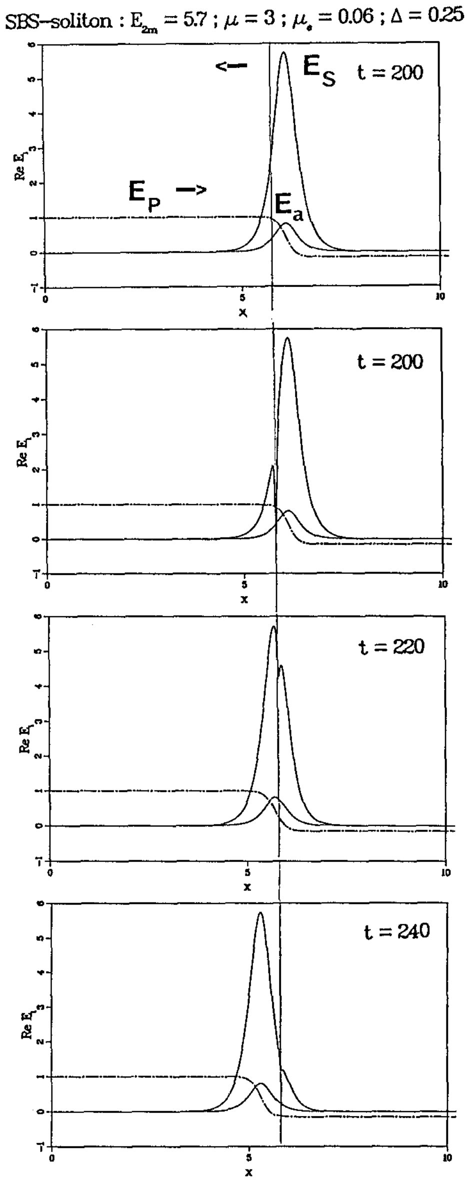


Fig. 2 *Dissipative superluminous Brillouin soliton given by formulas (7), for $\mu_a = 3$, $\mu_e = 6 \times 10^{-2}$ in the frame moving backward at the velocity of light (Stokes frame). Note that the pump is partially restored after the interaction. At time t=200 a signal is "engraved" in the soliton profile. The signal remains at rest in this backscattered frame and the Brillouin soliton runs backward away at the "superluminous" velocity. It cannot transport information.*

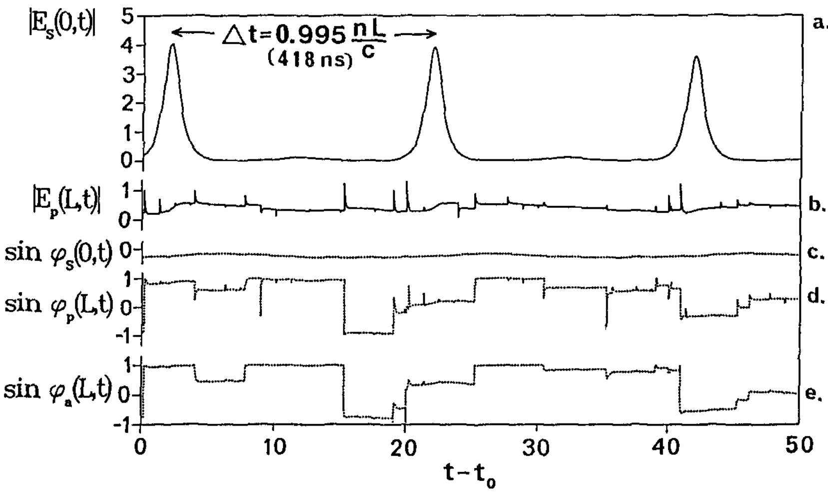


Fig. 3 *Numerical computation of Eqs.(1) for the ring cavity with* $\mu_a = 10$, $\mu_e = 10^{-2}$ *and* $L/\Lambda = 20$, *where* Λ *is the SBS characteristic length (Ref. [9]). Curve* (**a**), *train of backscattered solitons [output Stokes component* $E_S(0,t)$*]; curve* (**b**), *transmitted pump amplitude* $E_p(L,t)$ *showing spikes generated by random phase shifts; curve* (**c**), *Stokes phase, almost unsensitive to the pump phase jumps; curve* (**d**), *strongly modulated pump phase; curve* (**d**), *acoustic phase, following the pump phase jumps.*

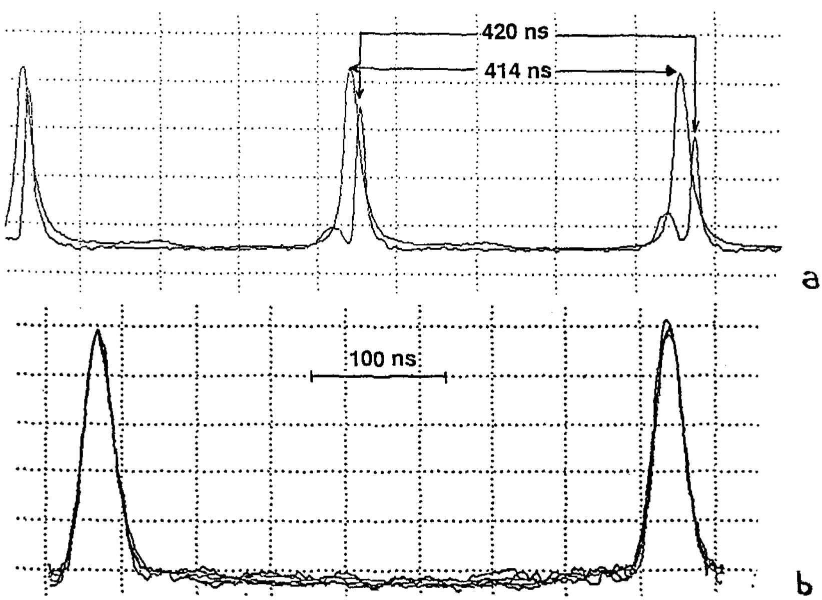


Fig. 4 (**a**) *Experimental superluminous Brillouin Stokes solitons superimposed to the evolutive pulses spaced by* Δt_L.
(**b**) *Experimental stability test of the Brillouin fiber-ring laser: six superimposed uncorrelated couples of pulses.*

III. Numerical and experimental results

In order to observe the spatial extended soliton envelopes in actual experimental situations, we must choose the interaction length as long as possible by allowing the Stokes envelope to be always in interaction with the pump in order to avoid the initial stage of the shock-wave regime (i), i.e. whithout cutting the Stokes leading edge. This has been achieved in the optical-fiber ring-cavity shown in figure 1, where the 50/50% beam splitter ensures the coupling of the cw pump wave into the cavity and the continuously recoupling of the backscattered Stokes wave. This configuration allows the Stokes wave to spread along the whole ring cavity.

However, we must also prevent the asymptotic steady state regime of the so-called *Brillouin mirror* : In the case of an ideal monochromatic cw pump wave, the numerical treatment of the coherent three-wave equations (1), with the feedback boundary conditions due to the cavity configuration, shows that, after some pulsed transients, the parametric SBS instability is saturated by the large depletion of the pump inside the medium, due to the accumulation of the phonons at the fiber input end, and gives rise to its reflection into a backscattered continuous Stokes wave [9]. This ideal case is indeed obtained in the experiment for high enough pump powers (beyond 250 mW coupled into the fiber), but for small cw input powers (below 100 mW) the experiments show pulsed nonstationary regimes exhibiting soliton-like profiles for the Stokes envelopes and self-induced transparency for the pump. These regimes are unstable and present sometimes several Stokes pulses in a round trip. But the stabilization of the pulsed regime, even at high intensity level, may be obtained by periodically interrupting the pump wave with the AOM during a time longer than the spontaneous acoustic damping time μ_a^{-1}. The modulation is achieved outside the ring cavity, so that the Stokes pulses may develop their wings through the whole ring cavity. For a good efficiency, the modulation frequency must roughly correspond to the round-trip flight time in the cavity. In this configuration the Brillouin mirror regime is totally prohibited and a stable soliton-like regime may be reached. We have obtained by this way sequences of well-shaped Stokes pulses for a few milliseconds [2]. In figure 4 are superimposed two couples of experimental Stokes pulses recorded at the output of the cavity during two consecutive sequences of qualitatively different dynamical behaviors. The first one is associated with the building phase [compressional regime (i)] of Stokes fine evolutive structures separated by 420ns, which is, within a precision of 2 ns, the round-trip flight time $\Delta t_L = Ln_0/c$ in the cavity at the group velocity of light in the fiber. The second one corresponds to a very stable sequence of hyperbolic-secant-like fitted pulses which repeat themselves every 414ns. The stability of these soliton-like pulses is quite remarkable as can be verified on figure 4(b), which shows six superimposed uncorrelated couples of Stokes pulses. We also verify that a non-zero level of pump intensity is transmitted at the other end of the fiber. Superluminous velocity and partial self-induced transparency are certainly the most undisputable proofs of the physical relevance of the dissipative Brillouin solitons. It is interesting to note that, by using such an externally modulated pump wave, a numerical treatment of Eqs.(1) shows the birth of the solitonlike profile by starting from completely stochastic amplitude and phase acoustic fluctuations.

It should nevertheless be kept in mind that the asymptotic stage is not reached. What we observe is an incompleted growing quasisoliton. In fact, in the experiment, a quasisoliton pulse is amplified and accelerated at each round trip in the cavity, but is also prevented from completing its growth every time it reaches the end of the fiber before being recoupled at the other end with a lower amplitude due to recoupling losses pertaining to the ring setup.

In conclusion, dissipative superluminous quasisolitons, observed for the first time in a Brillouin optical-fiber ring laser [2], can account for nonstationary dynamic behaviors of the backscattered Stokes wave and self-induced transparency on the transmitted pump wave. The coherent Brillouin soliton-like behaviour contributes to stabilize the Stokes output. It may be responsible of the high spectral coherence for the backscattered Stokes wave (much greater than the pump coherence) observed in recent experiments [20] [21].

[1] O. Legrand and C. Montes, J. Phys. Colloq. France, **50**, C3-147 (1989).
[2] E. Picholle, C. Montes, C. Leycuras, O. Legrand, and J. Botineau, Phys. Rev. Lett. **66**, 1454 (1991).
[3] E.P. Ippen and R.H. Stolen, Appl. Phys. Lett. **21**, 539 (1972).
[4] D. Cotter, J. Opt. Commun. **4**, 10 (1983).
[5] C. Montes, Phys. Rev. A **31**, 2366 (1985) and references therein.
[6] M. Maier, Phys. Rev. **166**, 113 (1968).
[7] W. Kaiser and M. Maier, in *Laser Handbook 2*, T.F. Arecchi and F.O. Schulz-Dubois, eds. (North-Holland, Amsterdam, 1972) p. 1077.
[8] J. Coste and C. Montes, Phys. Rev. A **34**, 3940 (1986)
[9] J. Botineau, C. Leycuras, C. Montes, and E. Picholle, J. Opt. Soc. Am. B **6**, 300 (1989).
[10] S.F. Morozov, L.V. Piskunova, M.M. Sushchik, and G.I. Freidman, Sov. J. Quant. Electron. **8**, 576 (1978).
[11] V.A. Gorbunov, Sov. J. Quant. Electron. **14**, 1066 (1984).
[12] G.L. Lamb, Jr., Phys. Lett. **29 A**, 507 (1969); Rev. Mod. Phys. **43**, 99 (1971).
[13] J. Pelous and R. Vacher, Solid State Commun. **16**, 279 (1975).
[14] J.A. Armstrong, S.J. Sudhanshu, and N.S. Shiren, J. Quant. Elect. **QE-6**, 591 (1974).
[15] K. Nozaki and T. Taniuti, J. Phys. Soc. Jpn. **34**, 796 (1973); Y. Oshawa and K. Nozaki, J. Phys. Soc. Jpn. **36**, 591 (1974).
[16] C.J. McKinstrie, Phys. Fluids **31**, 288 (1988)
[17] D.J. Kaup, Stud. Appl. Math. **55**, 9 (1976).
[18] S.C. Chiu, J. Math. Phys. **19**, 168 (1978).
[19] S.L. McCall and E.L. Hahn, Phys. Rev. **183**, 457 (1969).
[20] M. Douay, P. Bernage, and P. Niay, Opt. Comm. **81**, 231 (1991).
[21] S.P. Smith, F. Zarinetchi, and S.E. Ezekiel, Opt. Lett; **16**, 393 (1991).