Experiments with optical solitons

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This paper reviews the current literature on the formation and interaction of solitons in optical fibers. A number of original results of numerical modeling of the phenomena indicated are presented.

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CONTENTS

1. INTRODUCTION

The investigation of solitons and their interactions is one of the main problems of the nonlinear theory of waves.[1] Considerable progress has been achieved in the last ten years in this area due to the comprehensive application of exact, approximate, and numerical methods. At the present time, soliton solutions have been found for several tens of equations of interest in different areas of physics, from optics to biophysics.[1),2] Unfortunately, the starting equations for which solutions have been obtained in the form of stationary nonlinear waves are often model equations and describe only approximately wave processes in real media. Nevertheless, demonstration experiments were performed in electromagnetic lines, consisting of discrete linear and nonlinear elements,[3] in plasma, superconductors, and on the surface of water.[2)] The experimental realization of soliton regimes in optics gives rise to serious difficulties, since the medium and the initial conditions must satisfy quite stringent requirements.

It has been known for a long time that in the region of anomalous dispersion wave packet spreading can be balanced by self-compression, caused by nonlinearity, and this leads to the formation of envelope solitons.[4] The spatial analog of this phenomenon is the wave guide regime for propagation of light beams in self-focusing media.[5] An important difference between these two processes is that the stationary (two-dimensional) waveguide in a medium with a cubic nonlinearity is unstable, while the stationary (one-dimensional) pulse is a stable formation. Experimental observation of stationary light pulses in unbounded media is complicated, in particular, by the presence of spatial divergence or convergence, absorption, and accompanying nonlinear effects. It was proposed as far back as 1973 that optical fibers be used in experiments with time envelope solitons.[6] In an optical wave guide, the spatial profile of the intensity is determined by competition between two linear processes: diffraction and refraction and is practically unaffected by the nonlinearity. However, such experiments became a reality only recently[7] due to the realization of two prerequisites: 1) creation of single-model fibers with minimum losses in the region of anomalous dispersion; 2) perfection of sources of picosecond pulses, tunable in the near IR range. The next two sections will be concerned with a short discussion of these prerequisites.

In speaking of the significance of experiments with optical solitons, we note that they undoubtedly stimulate the interest of researchers in optical fibers as being unique, according to their possibilities, nonlinear media. Of interest are possible applications in the area of optical information systems and in problems of forming ultrashort light pulses with a given shape of the envelope for different physical experiments.

2. MODERN OPTICAL FIBERS AND THEIR NONLINEAR PROPERTIES

The current trend in fiber optics is a transition away from the visible to the near IR range.[8] This is due, in part, to the following: it is in this region ($\lambda \approx 1.55\ \mu$m) that fibers have minimum losses ~0.2 dB/km; in the near IR range, the dispersion passes through zero ($\lambda_0 \approx 1.3\ \mu$m; for $\lambda > \lambda_0$, $\omega''_{kk} > 0$, anomalous dispersion; for $\lambda < \lambda_0$, $\omega''_{kk} < 0$, normal dispersion). Figure 1 shows the dependence of the losses in dB/km for the ultimate fiber.[9] The increase in losses with decreasing λ is caused by Rayleigh scattering. Miya *et al.* associate

1) The selection of articles in this issue of the journal is entirely devoted to various aspects of the theory of solitons.

2) These problems are discussed in Solitons in Action edited by K. Lonngren and A. Scott, Acad. Pr., N.Y., 1978 (Mir, Moscow, 1981).

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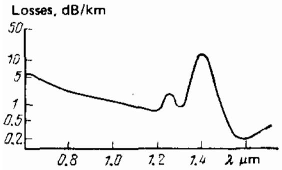

FIG. 1. The wavelength dependence of losses in a fiber.[9]

the absorption peaks visible in the figure with vibrational overtones of the hydroxyl group OH.

The dispersion properties of fibers have been studied, in particular, in Ref. 10. They are characterized by the parameter $D = 1/L\ \partial T/\partial\lambda$, which is the rate of change of the group delay time T with wavelength. A typical value of D is 20 ps/(nm · km) for $\lambda = 1.55\ \mu$m. We emphasize that the dispersion characteristics of a wave guide are determined by its structure and by the profile of the index of refraction; they are not identical with the dispersion properties of the starting material. The relationship between these properties is discussed in Ref. 11. It is shown that anomalous dispersion can be obtained in a fiber within the region of normal dispersion of the substance. Three-layer optical wave guides,[12] which have a number of advantages over two-layer wave guides, are very promising here. Japanese investigators report production of a fiber, for which the region of zero dispersion coincides with minimum losses.[13] Based on these achievements, communication lines with a rate of transmission of information up to $8 \cdot 10^8$ bits/s over a distance of up to 20 km have already been produced.[14]

Intensive experimental and theoretical investigation of the nonlinear properties of fibers began in the 1970's in connection with the fact that they limit the power and rate of transmission of signals along optical communication lines. Thus, in a fiber made of fused quartz, nonlinear effects are already observed in powers ~1 W, although in SiO_2, the nonlinear increment to the index of refraction is two orders of magnitude smaller than in carbon disulfide CS_2. The reason for this lies in the enormous interaction lengths ~1 km that can be realized in fibers with low losses.

At the present time, the following phenomena, stemming from the cubic susceptibility, are being studied experimentally: phase self-modulation and self-action of temporal envelopes, stimulated Raman scattering, four-photon mixing processes, stimulated Mandel'shtam-Brillouin scattering.[15] The first of the effects indicated is related to the dependence of the index of refraction n on the amplitude of the electric field of the light wave

$$n = n_0 + \delta n, \qquad \delta n = \frac{1}{2} n_2 |E|^2.$$

In typical materials for optical wave guides, n_2 arises due to the nonlinearity of electronic polarizability and reaches a steady state within a time $\sim 10^{-15}$ s; its characteristic order of magnitude is $n_2 \sim 10^{-13}$ CGS units.

The presence of δn can lead to self-focusing of a beam, if its power exceeds a critical level, but in single-mode fibers the changes in the spatial intensity profile are negligibly small (the characteristic power for self-focusing in a fiber with a diameter of several microns is of the order of 10^6–10^7 W). However, the increase in phase that appears $\delta\varphi = \delta n \omega z/c$ and the corresponding frequency shift $\delta\omega \propto kz\, \partial |E|^2/\partial t$, which time-dependent, has an appreciable effect on the shape of the pulse envelope and its spectrum for powers ~1 W. We note that for $n_2 > 0$, the frequency shift is negative on the front of the pulse and positive on the trailing edge. After the pulse ($\tau_0 = 200$ ps, $\lambda = 0.514\ \mu$m) traverses a distance of 255 meters along a fiber with a core diameter of 3.3 μm, the width of its spectrum doubles when the input power is 180 mW.[16] Propagation of a pulse with a broadened spectrum in a dispersive fiber greatly affects the pulse envelope. A detailed discussion of these effects in the region of normal dispersion is contained in Ref. 17, wherein the appearance of envelope shocks is demonstrated by numerical simulation for media with a finite relaxation time of the nonlinearity. Phase self-modulation greatly limits the transmission band of optical information systems. At the same time, the good correspondence between the theoretically predicted and experimentally measured shape of the spectrum permits using this phenomenon to determine n_2.[16] Picosecond pulses, undergoing phase self-modulation, are used as a wide band source in spectroscopy.

The characteristics of the occurrence of other nonlinear processes in optical fibers, stemming from the cubic susceptibility, are discussed in detail in Ref. 19. In this connection, we mention only Ref. 20, wherein stimulated Raman scattering and four-photon frequency mixing in a single-mode fiber, pumped in the region of zero dispersion, were studied experimentally.

Summarizing, we can say that from the point of view of nonlinear optics of media, optical fibers are very interesting due to their small losses, stable geometry, and low threshold power for typical nonlinear effects. We note that it is possible to isolate the self-action of temporal envelopes in "pure" form by varying the parameters of the initial pulse and by choosing appropriate fibers.

3. SOURCES OF PICOSECOND PULSES, TUNABLE IN THE NEAR IR RANGE

Sources of picosecond pulses for experiments with optical solitons must provide smooth tuning over the range 1.2–1.6 μm and must generate stable pulses with a wide spectrum, limited by the inverse pulse duration. In this connection, two classes of sources are of fundamental interest: parametric generators of picosecond pulses[21] and lasers utilizing color centers in alkali-halide crystals.[22] A typical example of the latter group is a laser utilizing F_2^+ centers in KF, described in Ref. 23, which is capable of generating pulses with duration 3–5 ps in a tuning range from 1.24 to 1.45 μm. The source operated at a temperature of 70 K and, in addition, an electron beam was used to color the crystal. Synchronous pumping was carried out with a Nd:YAG laser ($\lambda = 1.064\ \mu$m). An analogous arrangement using F_2^+ centers in NaCl (tuning range 1.35–1.75 μm), is mentioned in Ref. 7. In order to improve the spectral characteristics, the cavity dimensions were changed in order to obtain one output pulse for each pumping pulse. In addition, a birefringent element (a sapphire plate with a thickness of 4 mm) was placed in the cavity, which permitted obtaining pulses satisfying the relation $\tau_0 \Delta f = 0.18$.

Lasers capable of operating at room temperature are of great practical significance. Picosecond generation with forced mode-locking using a LiF crystal is reported in Refs. 24 and 25. Under the action of the pumping pulses ($\lambda = 0.53$ μm), two-step ionization occurs, while inversion occurs after absorption of energy by the F_2^+ centers formed. Lasing is achieved in the range $\lambda \sim 0.84-1.1$ μm.

Parametric generators of picosecond pulses,[26] which can be tuned over a wide range from 0.5 to 10 μm, form another class of promising sources. They operate on perfected solid state lasers, capable of generating powerful well-reproducible pulses with duration from 30 to 3 ps.[27] In order to improve the spectral quality of the pulses, cavity schemes with synchronous pumping are used. Recently, the possibility of using phase conjugation to obtain pulses with a spectral width limited by the inverse duration was demonstrated experimentally.[28]

Semiconducting lasers with mode-locking, capable of generating pulses with duration ~20 ps, could become very convenient and compact sources, if satisfactory spectral characteristics can be achieved.[29]

4. WAVE PICTURE OF PULSE PROPAGATION IN OPTICAL FIBERS

The theoretical analysis of the propagation of light pulses in optical fibers is usually carried out under the following assumptions[6]: 1) the electric field in the light wave is sought in the form

$$\overline{E}(x,\ y\ z\ t) = \frac{1}{2}\bar{e}\chi(x,\ y)\{\psi(z,\ t)\exp[i(\omega t - kz)] + \text{c.c.}\},$$

where $\chi(x,y)$ describes the spatial distribution of the field in the linear approximation, while the slowly varying amplitude of the temporal envelope $\psi(z,t)$ is affected by the nonlinearity; 2) second-order infinitesimals are taken into account in the expansion of $k(\omega)$ in powers of $(\omega - \omega_0)$.

In this case, $\psi(z,t)$ satisfies the nonlinear Schroedinger equation

$$i\frac{\partial\psi}{\partial z} = \pm\frac{1}{2}\left|\frac{\partial^2 k}{\partial\omega^2}\right|\frac{\partial^2\psi}{\partial\tau^2} + \alpha k\frac{n_2}{n_0}|\psi|^2\psi - i\delta_0\psi; \qquad (1)$$

where $\tau = (t - (z/v_g))$ is the instantaneous time, $v_g = \partial\omega/\partial k$, the pulse sign corresponds to anomalous dispersion, the minus sign corresponds to normal dispersion, and δ_0 is the absorption coefficient. The appearance of the factor α is related to the spatial distribution of the field. In particular, for a fiber with a step-like profile of the index of refraction $\alpha = \langle\chi^4\rangle/\langle\chi^2\rangle$ for the HE_{11} mode $\chi = J_0(\mu, \rho)$; averaging is carried out over the radial distribution and μ, is the first root of the zeroth-order Bessel function. The effect of the distribution of the index of refraction over the cross section of the fiber on the dispersion properties of the fiber is analyzed in detail in Ref. 11.

We note that Eq. (1) was obtained assuming that the optical fiber does not change the linear polarization of the input radiation (such fibers have already been developed[30]). If the polarization is not conserved, then n_2 in Eq. (1) is replaced by the polarization-averaged quantity $(5/6)\, n_2$.

In order to expose the characteristic scales of the problem, it is convenient to transform to dimensionless variables, normalizing the instantaneous time to the initial duration of the pulse $\tau' = \tau/\tau_0$, the distance z to the dispersion length $z_d = \tau_0^2/|k''_{\omega\omega}|$, and the field amplitude to the characteristic initial value $|\psi_0|$. Then, (1) assumes the form[3)]

$$i\frac{\partial\psi}{\partial z} = \pm\frac{1}{2}\frac{\partial^2\psi}{\partial\tau^2} + R|\psi|^2\psi - i\delta\psi. \qquad (2)$$

The nonlinearity parameter R has the following structure: $R = z_d/z_{nl}$, where the nonlinear length is inversely proportional to the intensity at the input to the medium $z_{nl} = cn_0^2/(8\pi k n_2 I_0)$; δ is the absorption over the dispersion length. Under real conditions, R can vary over a wide range due to a change in the input intensity I_0 as well as tuning of the laser near the wavelength corresponding to zero dispersion ($k''_{\omega\omega}$ and z_d vary). We note that for $k''_{\omega\omega} = 0$ Eq. (1) must contain higher order dispersion terms, corresponding to $k'''_{\omega\omega\omega}$ and so on.

Equation (2) (with $\delta = 0$) belongs to a class of equations that are exactly solvable by methods of the inverse problem in scattering theory.[31,32] We recall that the soliton solutions of (2) have the form (3)

$$\psi = R^{-1/2}\varkappa_j \operatorname{sech}[\varkappa_j(\tau - \tau_j + z\Omega_j)]\exp[i(\Omega_j\tau - K_j z + \varphi_j)]. \qquad (3)$$

where τ_j, φ_j, Ω_j, k_j, and $\varkappa_j^2 = \Omega_j^2 + 2K_j$ are dimensionless parameters, which determine the soliton position, phase, and velocity in a co-moving coordinate system.

For arbitrary initial conditions, the soliton component contained in them will determine the asymptotic behavior of the solution. In Ref. 32, the nonsoliton part of the solution was estimated to decrease at a rate $z^{-1/2}$. In addition, the problem of the interaction of solitons was investigated. Within the scope of Eq. (2) ($\delta = 0$), it reduces to a shift in the trajectories of their centers and phases. The additivity of the shifts as a result of the collision of a large number of solitons is also demonstrated. These results have fundamental significance, since they can be used to check approximate methods and numerical algorithms. However, application of the methods of the inverse problem in scattering theory to solving problems with arbitrary initial conditions give rise to serious difficulties. Other "exact" methods for solving nonlinear problems are presented in an accessible form in Ref. 33.

One of the approximate approaches is explained in Ref. 34. A Lagrangian, analogous to the Lagrangian for particles with a pair interaction potential, is introduced for a system of solitons with nearly equal velocities. Conditions necessary for the existence of bound states of solitons are formulated. Perturbation theory for such problems is developed in Ref. 35.

[3)]The primes on the dimensionless variables here and in what follows are dropped.

Investigation of Eq. (2) using the method of moments yields useful and graphic results.[36] For a mean-square pulse duration $\langle\tau^2\rangle = \int\tau\psi\psi^* d\tau/P_0$, it is not difficult to obtain the equation

$$\frac{P_0}{2}\frac{d^2}{dz^2}\langle\tau^2\rangle = H_0 + \frac{R}{2}\int(\psi\psi^*)^2\, d\tau, \tag{4}$$

where $P_0 = \int_{-\infty}^{+\infty}\psi\psi^* d\tau$, $H_0 = \int_{-\infty}^{+\infty}[\psi'_\tau\psi'^*_\tau - R(\psi\psi^*)^2]d\tau$ are integrals of the problem (2), which represent the total pulse energy and the Hamiltonian. It is evident that for a given initial field distribution $\psi_0(\tau, 0)$, there exists a R_0 such that for $R < R_0$, dispersive spreading dominates:

$$R_0 = \int(\psi'_\tau\psi'^*_\tau)\, d\tau\left(2\int\left(\psi\psi^*\right)^2 d\tau\right)^{-1}. \tag{5}$$

In particular, for $\psi_0 = A_0\,\mathrm{sech}\tau$, $R_0 = 0.5$ with $A_0 = 1$. In the interval $R_0 < R < 2R_0$, with z increasing from zero, the pulse begins to broaden and after several damped oscillations, its width stabilizes. In the interval $2R_0 < R < 4R_0$, as z increases, the pulse narrows, then broadens, and after a number of oscillations, its amplitude reaches a steady state value.

It is convenient to study the dynamics of this process numerically.[35] Figure 2 shows the peak amplitude of the pulse as a function of z for $A_0 = 0.8$, 1.1, 1.25, and 1.4. For $A_0 = N$ ($N = 2, 3, \ldots$) and purely real initial conditions $\psi = A_0\,\mathrm{sech}\tau$, a N-soliton bound state is formed.

The method of finite differences is used to investigate Eq. (2) numerically.[35] Separation according to physical factors,[14] when the propagation of a wave packet in a dispersive medium is described as successive transmission through layers with purely dispersive and purely nonlinear properties, is widely used. The fast Fourier transform is used at the dispersive step. The mathematical justification and estimate of the accuracy are presented in the appendix to Ref. 37. Application of the method of finite elements to problems in nonlinear optics has turned out to be very effective.[38,39]

Recently, interest has appeared in studying the self-action of pulses in multimode fibers.[40] This problem leads to a system of nonlinear Schroedinger type equations for the envelopes of pulses corresponding to different modes. The nonlinear increment to the index of refraction is proportional to the total intensity due to the orthogonality of the modes. In Ref. 40, A. Hasegawa *et al.* obtained estimates of the power for which the intermode dispersion is suppressed by the reactive (only via the index of refraction) interaction of pulses. The orders of magnitude of the quantities here are as follows: for an input power ~20 W (intensity ~10^7 W/cm^2), modes with group delay ~100 ps/km are trapped.

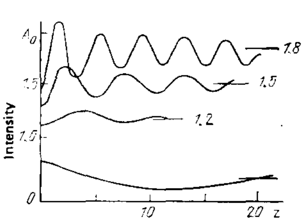

FIG. 2. The instantaneous value of the peak amplitude as a function of distance for different initial amplitudes.[35]

In the particular case of two modes, the starting system of equations can be reduced to the form[41]

$$\begin{aligned} i\frac{\partial\psi_1}{\partial z} &= \frac{1}{2}\frac{\partial^2\psi_1}{\partial\tau^2} + \alpha_{11}|\psi_1|^2\psi_1 + \alpha_{12}|\psi_2|^2\psi_1,\\ i\frac{\partial\psi_2}{\partial z} &= \frac{1}{2}\frac{\partial^2\psi_2}{\partial\tau^2} + \alpha_{22}|\psi_2|^2\psi_2 + \alpha_{21}|\psi_1|^2\psi_2, \end{aligned} \tag{6}$$

where $\alpha_{mj} = R_d R^{-1}_{nl\,m}\langle\chi^2_m\chi^2_j\rangle/\langle\chi^2_m\rangle$, averaging is carried out over the radial distribution, $R_{nl\,m}$ is the nonlinear length for the m-th mode, analogous to that introduced previously. The group delay relative to the average velocity must be taken into account in the initial conditions, which are written as follows:

$$\psi_1(\tau, 0) = \psi_{10}(\tau)\exp(i\zeta\tau),$$
$$\psi_2(\tau, 0) = \psi_{20}(\tau)\exp(-i\zeta\tau).$$

The parameter ζ is the group delay over the dispersion length, scaled to the initial duration of the pulse.

Some data from numerical calculations[41] are presented in Fig. 3. Case a) corresponds to partial suppression of intermode dispersion ($\alpha_{mj} = 1$, $\zeta = 1.63$) with total suppression of spreading of an individual pulse. Case b) illustrates total compensation of intermode dispersion for the minimum value of the detuning parameter $\zeta = 1$. The initial field distribution in both cases was purely real and had the form $\psi_{10}(\tau) = \psi_{20}(\tau) = \lambda\,\mathrm{sech}(\lambda\tau)$, $\lambda = 2$.

In concluding this section, we mention a group of papers, concerning self-action of pulses of partially co-

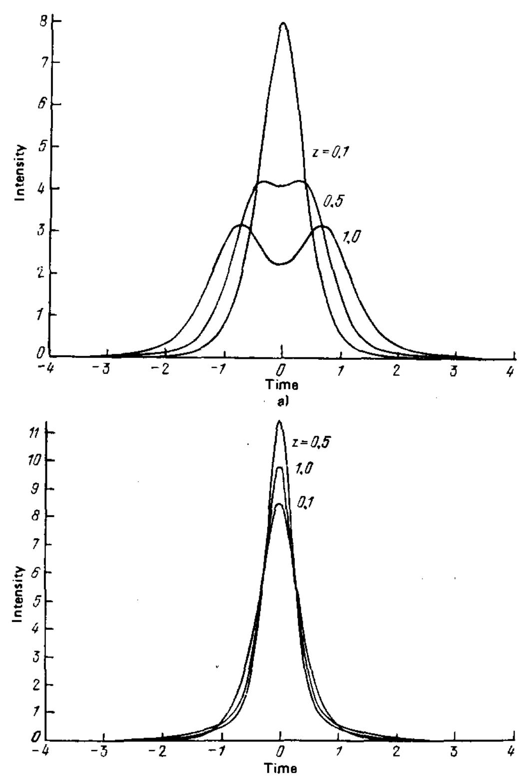

FIG. 3. a) Partial suppression of intermode dispersion with complete suppression of pulse spreading; b) complete suppression of intermode dispersion with minimum detuning.

herent radiation in fibers. Reference 42 is concerned with linear pulse propagation in a single-mode fiber. The input pulse with a Gaussian envelope represents a superposition of laser modes, taken with random phases. It is shown that fluctuations in the pulse power damp out with distance, if the intermode dispersion exceeds the dispersive spreading of the envelope. In the opposite case, fluctuations exist independently of the fiber length. The effect of intensity fluctuations on nonlinear propagation was investigated in Ref. 43. The critical power for formation of a soliton was estimated.

5. EXPERIMENTAL RESULTS. POSSIBLE APPLICATIONS

Envelope solitons in optical fibers were successfully observed experimentally in Ref. 7. Pulses from a laser based on F_2^+ centers in a NaCl crystal (tuning range 1.35–1.75 μm, working wavelength $\lambda = 1.55$ μm), entered the optical fiber. The spectrum of the signal and the autocorreation function of the intensity $B(\tau)$, which was determined by using a technique based on generating the second harmonic in ammonium dihydrophosphate, were measured at the output of the optical system. Data on the spectrum and autocorrelation function are not enough to reconstruct completely the envelope of the input pulse. The value of the product $\tau_0\Delta f \approx 0.18$ ($\tau_0 = 6$ ps) suggests that the pulses are free of phase modulation and, with respect to shape, occupy at intermediate position between pulses with an envelope of the form $\mathrm{sech}\,\tau (\tau_0\Delta f = 0.315)$ and a decaying exponential $\tau_0\Delta f = 0.11$.

A fiber made of doped fused quartz, which had a steplike profile of the index of refraction (core diameter 9.3 μm, length 700 m, losses at $\lambda = 1.55$ μm were ~1 dB/km), was used. The power of the input signal varied in the range 0.3–22.5 W.

Typical experimental profiles of the autocorrelation functions are presented in Fig. 4 for different values of the input signal power. For $P_0 = 0.3$ W, the pulse broadening corresponds to a linear regime and agrees well with the computed boradening. As P_0 increases, the pulse at the fiber output is compressed and for $P_0 = 1.2$ W, its duration equals the initial value. This indicates that dispersive spreading is completely compensated by the action of the nonlinearity and this case can be interpreted as a single soliton. Next, compression of the pulse begins (up to 2 ps with $P_0 = 5$ W). Then, the plateau in the correlation function begins to increase and two secondary maxima appear on it ($P_0 = 11.5$ W). This behavior of $B(\tau)$ can be understood by considering the temporal intensity profiles that we computed on a computer. The series of curves in Fig. 5(a) corresponds

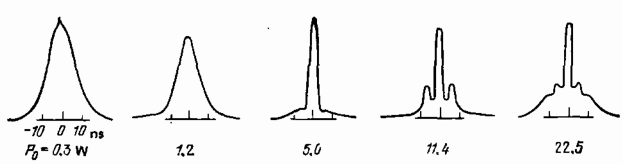

FIG. 4. The experimental profiles of the autocorrelation function of the intensity at the fiber output for different initial input signal power.[7]

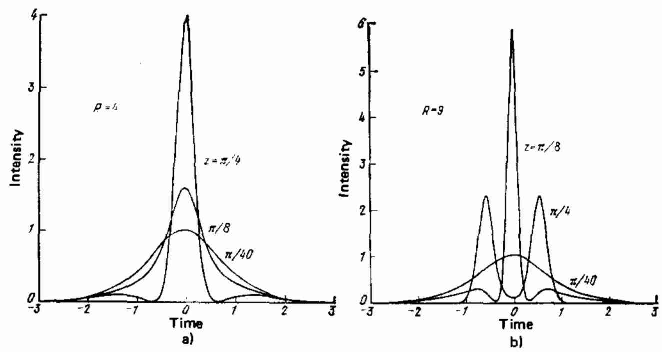

FIG. 5. a) Temporal intensity profiles for a two-soliton bound state; b) three-soliton bound state.

to a four-fold increase in power above the critical level and in Fig. 5(b) to nine-fold increase. This indicates that in the first case we are dealing with a bound state of two and in the second case of three solitons. For $z = \pi/4$ (in dimensional variables, this is 700 m), the pulse has a two-peak structure, and this is what leads to the presence of three maxima in the correlation function. Finally, $P_0 = 22.5$ W corresponds to a bound state of four solitons. The characteristic values of the power indicated above differ from the theoretical predictions on the average by 20%. We feel that this is caused by the asymmetry of the input pulse. We note that the destruction of the initial polarization in the fiber must also be taken into account.

In order to illustrate the propagation of an asymmetric pulse, we shall present data from a numerical experiment.[41] The starting distribution had the form $\psi_0 = \tau^2 \exp(-\tau^2/2)$. The intensity profiles are shown in Fig. 6 for different z. It is evident how, according to the compression resulting from the nonlinearity, the pulse is symmetrized and a nonsoliton increment is separated from it. Thus, the basic characteristics of soliton formation are insensitive to distortions of shape as compared to the classical shape.

We shall proceed to a discussion of the possible ap-

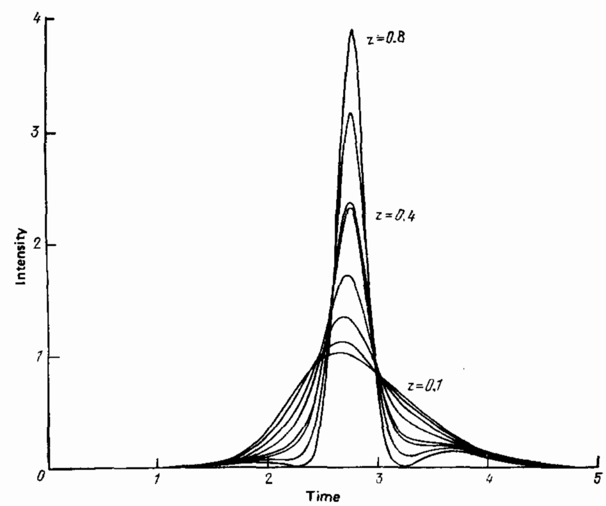

FIG. 6. Pulse symmetrization in an optical fiber.[41]

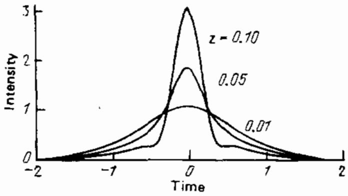

FIG. 7. Compression of a phase self-modulated pulse (linear propagation regime in a fiber).

plications of the phenomena examined above. Obtaining subpicosecond pulses is of great interest for many areas of physics. One of the possible ways to solve this problem is to compress pulses with long duration. Active control of the amplitude and phase of picosecond pulses is as yet technically not feasible. However, if the starting pulses are frequency modulated so that the frequency increases toward the end of the pulse (positive sweeping), then transmission of a pulse through a medium with anomalous dispersion causes the pulse to compress. In practice, an arrangement that provides anomalous dispersion may consist simply of a pair of diffraction gratings.[15] In order to obtain positive sweeping, a nonlinear effect, namely, phase self-modulation, is used. Thus, Laubereau *et al.*[44] experimentally compressed pulses from 20 to 2 ps.

The use of optical fibers opens up new prospects in this field. It is already evident from Fig. 5 that by increasing the parameter R (due to an increase in the input power or to tuning to adjust to the wavelength corresponding to zero dispersion), it is possible to compress the input pulse by a factor of 6–10 at a fixed distance.

Phase modulation of the input pulse gives additional possibilities. If the pulse was phase-modulated in an inertia-free strongly nonlinear medium, then it is possible to write the initial conditions at the fiber input in the form

$$\psi(\tau, 0) = \psi_0(\tau) \exp\left[-i\,|\,\psi^2\,|\frac{1}{2S}\right],$$

where the parameter S determines the rate of change of the frequency. We note that sweeping in this case is

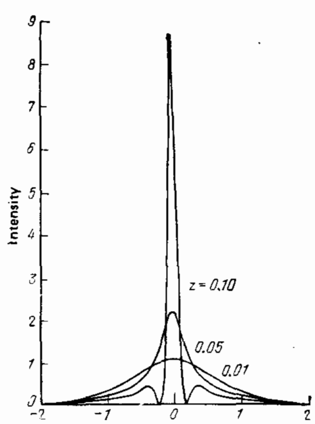

FIG. 8. Compression of a phase self-modulated pulse in the nonlinear regime.

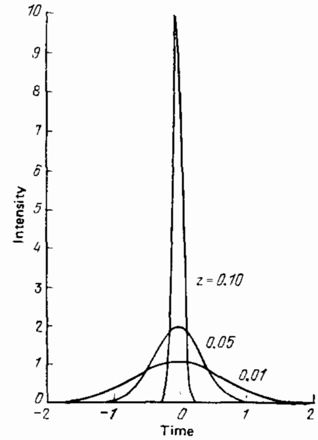

FIG. 9. Pulse compression with linear frequency sweep.

nonlinear. Figure 7 illustrates the pulse envelope along the path in a linear regime ($R=0$); the values of z are indicated in the figure in fractions of the dispersion length, $S=0.1$, and $\psi_0=\exp(-\tau^2/2)$. Figure 8 corresponds to the nonlinear regime ($R=14.3$). Aside from dispersion compression, the pulse is additionally compressed in the fiber due to self-action. Analogous profiles, obtained with linear frequency sweeping $\psi(\tau, 0) = \psi_0 \exp[-i\tau^2/(2S)]$ ($S=0.1$, $R=0$) are presented in Fig. 9. It is evident that here compression is more effective since the tails of the wave packet are pulled in toward the center more strongly.

The possiblities of optical fibers, from the point of view of the action on the shape of envelopes, are not limited to compression. The initial pulses can be symmetrized (see Fig. 6), a noisy envelope can be smoothed, and a sequence of two or more pulses following one another at short intervals can be obtained.

6. CONCLUSIONS

In conclusion, we shall list the possible lines along which the problems indicated above can be developed as well as the urgent problems. From an analysis of the published literature, it follows that experiments with optical solitons are important not only for confirming theoretical results of the nonlinear theory of waves, but they also open up new perspectives in problems of controlling the envelope and the spectrum of picosecond pulses. In this connection, the application of the methods of optimal control theory to such problems, developed for light beams in spatial problems,[45,46] is of interest. One interesting problem is the detailed examination of the nonlinear regime directly in the region of zero (in second order) dispersion. From the point of view of applications in fiber communications, experimental investigation of the possibilities of nonlinear suppression of intermode dispersion would appear to be useful. And, finally, all the directions listed above require extension to the case of partially coherent pulses.[47]

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