Generation of ultrashort pulses by spectral filtering during stimulated Raman scattering in an optical fiber

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The generation of an intense continuum when a quartz fiber is addressed by a Nd:YAG laser ($\lambda = 1.064\,\mu$m, pulse length 60 ps) is investigated. A package of computer programs has been developed for the computation of the nonlinear dynamics of ultrashort pulses produced during stimulated Raman scattering in optical fibers. The physical origin of the intense continuum due to the combined effect of nonlinear Raman frequency conversion, self-phase modulation, and group velocity dispersion of pulses channeled by optical fibers is examined. The dynamics of frequency scanning in pump pulses during SRS in single-mode optical fibers is investigated in detail and it is shown that, under the conditions of time-dependent saturation of stimulated Raman scattering, radiation with fundamental and Raman frequencies undergoes strong spectral broadening due to the appearance of sharp changes in amplitude and phase structure of pump pulses and SRS. It is shown that this mechanism for producing an intense continuum during SRS in optical fibers is universal, i.e., it occurs independently of the initial amplitude and phase structure of the pulses. It is shown experimentally that spectral filtering of the continuum can be used to generate high-contrast ultrasharp pulses. The technique has been used to produce 2.5-ps pulses.

Considerable attention is currently being devoted to nonlinear-optics processes that occur when powerful ultrashort pulses are channeled by optical fibers. This has been due mainly to the successful application of optical fibers to the solution of one of the most actively studied problems in modern quantum electronics, namely, generating ultrashort pulses in the picosecond and femtosecond ranges, while controlling properties such as the amplitude and phase envelopes, carrier frequency, and spectral composition.

Three lines of development can now be identified in this area.

1. The first method relies on dispersive compression of frequency-modulated pulses, using single-mode optical fibers for wide-range frequency scanning. As early as 1975, Lugovoĭ[1] proposed that the nonlinearity of the refractive index of the material of the single-mode optical fiber could be used to produce wide-range frequency scanning of laser pulses, so that the resulting frequency scanning interval could substantially exceed the initial spectral width of the pulse.

In single-mode optical fibers, the field distribution over the fiber cross section is determined in the linear approximation (i.e., up to power levels $I_{\rm cr} \sim 10^6$ W, where $I_{\rm cr}$ is the critical self-focusing power) by the competition between two linear effects, namely, diffraction and refraction, and corresponds to an approximately Gaussian radial distribution of the guiding mode, which remains unaltered along the entire length of the optical fiber. Hence, in contrast to self-phase modulation of the field in ordinary laser beams in homogeneous nonlinear media, the broadening of the spectrum in the single-mode optical fiber is not accompanied by a change in the spatial structure of the beam, or the redistribution of energy over the cross section by diffraction and the accompanying nonlinear effects, such as self-focusing, nonlinear absorption, and so on. This, together with the very considerable propagation range of low-loss optical fibers ($\lesssim 1$ dB/km), means that it is possible to produce broadening of the laser pulse spectrum that is the same over the entire cross section of the beam,[2] and use the optical fiber as a nonlinear phase modulator in dispersive pulse-compression systems. Considerable advances have already been made in this area.[3–7] For example, Fujimoto *et al.*[7] use the compression of frequency-modulated pulses in linear dispersive compression systems to produce 16-*fs* pulses, while Nicolaus and Grischkowsky[3] and Akhmanov *et al.*[6] have built a frequency-tunable generator of subpicosecond light pulses. It is important to note that, although this method of producing ultrashort pulses has been under development for more than 15 years (see, for example, Ref. 8), only since the advent of high-quality single-mode optical fibers has it been possible to reach the femtosecond range and produce up to 80-fold compression of pulses.[4] We emphasize that this method of pulse compression is intended mostly for the visible and near-infrared ranges, in which group velocity dispersion is positive (normal).

2. The second method is the "soliton compression" of pulses. It is well-known that, in the region of negative (anomalous) group velocity dispersion, the combined effect of dispersion and nonlinearity can lead to the self-compression of optical pulses and the generation of envelope solitons.[9–11] The propagation of a multi-soliton pulse along an optical fiber is accompanied by a periodic variation in the envelope, namely, self-compression and fragmentation of the pulse where, at the point of maximum self-compression, a narrow central intensity peak is produced on a broad pedestal.[12] This method was successfully used in Refs. 13 and 14 to achieve a 27- and 110-fold compression of a laser pulse in the spectral region corresponding to negative group velocity dispersion ($\lambda = 1.55\,\mu$m). It has been shown[15–17] that interesting new opportunities for controlling the nonlinear dynamics of picosecond pulses are provided by optical fibers for the middle infrared range,[15] optical fibers containing res-

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onant absorbing impurities,[16] and graded fibers.[17]

3. A new approach to the generation of ultrashort pulses in optical fibers involves the combined use of two nonlinear effects, namely, stimulated Raman scattering (SRS) and self-phase modulation.

It has been suggested[18] that multistage generation of high-order Stokes SRS components could be used in the region corresponding to negative group velocity dispersion for the nonlinear conversion of solitons in low-mode optical fibers. It is shown in Ref. 18 that stimulated Raman frequency conversion in a low-mode optical fiber can be used to transform a bound state of several solitons at the pump frequency into a high-energy single soliton pulse at the Stokes frequency when the mismatch in the pulse group velocities at the fundamental and Raman frequencies is balanced by intermode dispersion in the optical fiber.

Dianov *et al.*[19] have suggested that pulse compression can be produced in the region corresponding to positive group velocity dispersion by exploiting the "multiplication" of the initial frequency-scanning range during the generation of the higher Stokes components of SRS in an optical fiber, proposed in Ref. 20. They used the method of multistage nonlinear conversion of the time dependence of the multimode solid-state laser beam to generate ultrashort pulses at the Raman frequencies falling into the "zero" and negative group velocity dispersion regions.

The combined effect of stimulated Raman scattering and self-phase modulation during the propagation of powerful light pulses in an optical fiber can be used to generate a spectral continuum which extends several hundred reciprocal centimeters into the anti-Stokes region while, in the Stokes region, it overlaps the SRS spectrum.[21] It is shown in Ref. 21 that spectral filtration of the radiation leaving the fiber can be used to isolate single picosecond pulses at a frequency shifted relative to the pump frequency.

It is worth noting that high-intensity spectral continua have been observed and investigated during high-intensity pico- and subpicosecond excitation in different media (see, for example, the review given in Ref. 8), including optical fibers. The different mechanisms that have been proposed to explain the origin of these continua include self-phase modulation (in some cases, self-focusing), cascade ionization, and four-photon parametric processes. (We note that four-photon shifts in optical fibers have recently been under active investigation; see Ref. 22 and the references therein.)

The aim of the present research was to investigate the physical nature of the generation of the spectral continuum accompanying SRS in single-mode optical fibers and the possibility of a frequency-tuned source of ultrashort pulses exploiting the spectral filtration of radiation leaving an optical fiber.

1. EXPERIMENTAL METHOD

A detailed description of the experimental arrangement is given in Ref. 21. The beam of a continuously pumped Nd:YAG laser operated with active mode locking and Q-switching was introduced into a single-mode fused quartz optical fiber, 10 m long. The width of the pump pulses at half-height was 60 ps.

The radiation leaving the fiber was intercepted by an $LiIO_3$ crystal correlator which produced noncollinear frequency doubling for pulse-length measurement, using a zero-background autocorrelation scheme. The phase-match width of the $LiIO_3$ crystal was 6 nm. The second-harmonic radiation was then allowed to enter a grating monochromator incorporating an FEU-28 photomultiplier behind the exit slit, and the resulting electrical signal was processed by a detection module.

We were thus able to analyze the time characteristics for a given spectral selection of the radiation components leaving the optical fiber. Selection was produced both by scanning the monochromator slit (spectral slit width 5 nm) and suitably rotating the nonlinear crystal through small angles.

When the pump power P_p was raised to a value close to the SRS threshold, the pump spectrum leaving the fiber was broadened by self-phase modulation to about 10 cm^{-1}. Further increase in P_p produced the first Stokes component of SRS, which was shifted relative to the pump line by 440 cm^{-1} and had a width of 80 cm^{-1} (Fig. 1). SRS generation was accompanied in the spectrum by a broad continuum extending to both the Stokes and anti-Stokes regions.

We have estimated the critical pump power P_p entering the fiber, for which the intensities of the first Stokes component of SRS and of the pump were equal at exit from the optical fiber. The estimate was based on the theoretical results obtained in Ref. 23 for $P_{cr} gL / S_{eff} = 16$, where S_{eff} is the effective mode area. For our optical fiber with a core diameter of 7.5 μm and refractive index difference between the core and the cladding of 4×10^{-3}, this area amounted to 3.9×10^{-7} cm^2. The gain for linear polarization was $g = 0.92\times10^{-11}$ cm/W (Ref. 24). The critical power P_{cr} was 1.5–2 kW, the spread in the estimated P_{cr} being due to the fact that, in our experiment, the polarization in the optical fiber was not linear.

Figure 2 shows the autocorrelation functions for the pump radiation during effective SRS. The three peaks on the autocorrelation function correspond to the two pulses at $\lambda = 1.064\,\mu$m, produced as a result of the transfer of pump power from the central portion of the pump pulse to the pulse at the Stokes frequency (illustrated below in Figs. 5d–f). The separation between these pulses depends on the SRS

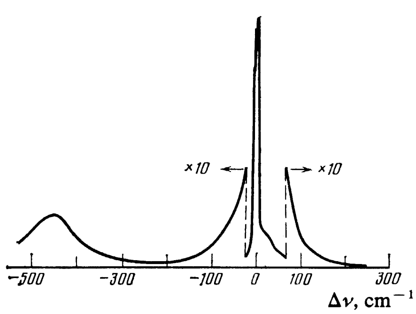

FIG. 1. Spectrum at exit from an optical fiber pumped by laser pulses with $\lambda_p = 1.064\,\mu$m and $\tau_p = 60$ ps.

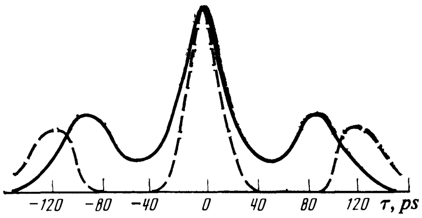

FIG. 2. Autocorrelation functions of radiation at the pump frequency during SRS in a fiber for different pump levels: equal to P_p—broken curve, greater than P_p—solid curve.

conversion efficiency, i.e., on the extent to which P_p exceeds P_{cr} (Fig. 2).

For $\Delta\nu = 440\ \text{cm}^{-1}$ (intensity maximum of the Stokes component of SRS), the width of the autocorrelation function for $P_p \simeq P_{cr}$ (Fig. 3) is 27 ps at half-height which, for a Gaussian pulse, gives a width of 19 ps and can be explained by the shortening of the pump pulse due to the exponential regime of SRS amplification, namely, $\exp(P_p gL/S_{\text{eff}})$. (We note that, when $P_p \simeq P_{cr}$, the SRS pulse is much more unstable than the pump pulse; the autocorrelation function of the first Stokes component of SRS in Fig. 3 was therefore recorded with a long averaging time constant.) Further increase in P_p was found to be accompanied by the cascade generation of the second Stokes component of SRS, which "burnt" a hole in the pulse of the first Stokes component.

Figures 4a and b show the autocorrelation functions for frequency shifts $\Delta\nu = \nu_a - \nu_p$ into the anti-Stokes region amounting to 36 and 104 cm^{-1}. The autocorrelation functions were recorded for $P_p > P_{cr}$. It is clear that the width of the central peak decreases with increasing $\Delta\nu$. At the same time, the intensity of the side peaks decreases and vanishes altogether for $\Delta\nu = 104\ \text{cm}^{-1}$, which is an indication of complete discrimination against one of the pulses. As a result, spectral filtering enables us to select a single-mode pulse of $\tau_p \simeq 2.5$ ps (3.6 ps wide autocorrelation function at half-height).

When the spectral interval was selected in the Stokes region relative to the pump region with $\Delta\nu = \nu_p - \nu_s < 300$ cm^{-1}, the autocorrelation function was practically the same as those shown in Fig. 2. We note that, in this region, the pedestal due to the pump overlaps the wing of the Stokes component of SRS.

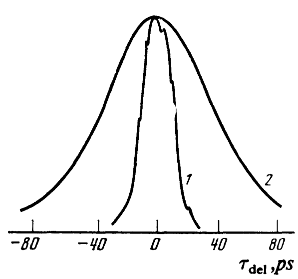

FIG. 3. Autocorrelation function at the frequency of the first Stokes component of SRS (curve 1) and of the pump at entry to the fiber (curve 2).

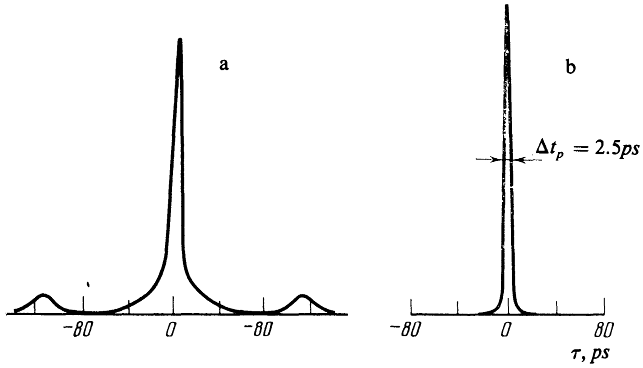

FIG. 4. Autocorrelation functions of radiation leaving an optical fiber in the anti-Stokes region relative to the pump: a—$\Delta\nu = \nu_a - \nu_p = 36\ \text{cm}^{-1}$, b—104 cm^{-1}.

2. DYNAMICS OF FREQUENCY SCANNING DURING SRS IN A SINGLE-MODE OPTICAL FIBER

Our results can be explained by the appearance of sharp changes in the phase structure of pump and SRS pulses, and the correspondingly large spectral broadening of the pump radiation when energy is transferred rapidly between waves at the fundamental and Raman generation frequencies (saturation of SRS gain).

The mathematical model of the SRS process in single-mode optical fibers, which takes into account the dependence of the refractive index on the radiation intensity and, in general, an arbitrary number of high-order components, is developed in Ref. 20 in first-order dispersion theory. In second order in this theory, when group velocity dispersion within the spectrum of each of the interacting waves is taken into account, the dynamics of the generation of ultrashort SRS pulses is described by a set of nonlinear Schrödinger-type equations which, in the simplest case of interaction between the pump and the first Stokes component, assumes the following form:

$$2i\frac{\partial E_p}{\partial z} = -\frac{\partial^2 k_p}{\partial\omega^2}\frac{\partial^2 E_p}{\partial\tau^2} + \frac{n_2(\omega_p)k_p\alpha_{pp}}{n_0}|E_p|^2E_p + \frac{n_2(\omega_s)k_p\alpha_{ps}}{n_0}|E_s|^2E_p - ig\alpha_{ps}\frac{\omega_p}{\omega_s}QE_s - i\gamma(\omega_p)E_p \qquad (1)$$

$$2i\left[\frac{\partial E_s}{\partial z} + \left(\frac{1}{v_s} - \frac{1}{v_p}\right)\frac{\partial E_s}{\partial\tau}\right] = -\frac{\partial^2 k_s}{\partial\omega^2}\frac{\partial^2 E_s}{\partial\tau^2} + \frac{n_2(\omega_s)k_s\alpha_{ss}|E_s|^2E_s}{n_0} + \frac{n_2(\omega_p)k_s\alpha_{sp}}{n_0}|E_p|^2E_s + ig\alpha_{sp}Q^*E_p - i\gamma(\omega_s)E_s + N(z,\tau). \qquad (2)$$

where $E_{p,s}(z,\tau)$ are the complex amplitudes of the pump and first Stokes components, Q is the amplitude of the wave of molecular vibrations of the medium, $\tau = t - z/v_p$ is the "running" time associated with the pump, $v_{p,s}$ are the group velocities of pulses at the fundamental and Raman frequencies, g is the stationary gain of the Stokes wave at the maximum of the amplification line, $\gamma(\omega_{p,s})$ are the linear optical loss coefficients, $\alpha_{mn} = \langle\psi_{mn}^{22}\rangle/\langle\psi_m^2\rangle$ is a geometric factor that appears when the nonlinear cubic susceptibility is averaged over the cross section of the fiber, $\psi_{m,n}(r)$ are spatial functions describing the radial field distribution in the optical fiber in the linear approximation, n_0 is the effective linear refractive index of the fiber material, determined, in particular, by the structure of the refractive-index profile, $d^2k/\partial\omega^2$ is a coefficient representing group velocity dispersion of the light pulse, which depends both on the dispersion of the fiber material and the waveguide dispersion, and $N(z,\tau)$ is the source of spontaneous field fluctuations at the Raman frequency.[25]

In the general case of inhomogeneous broadening of the SRS line with a Gaussian distribution function

$$G(\omega_0-\Omega)=g(0)\exp\left\{-\frac{(\omega_0-\Omega)^2}{\Delta_0^2}\right\}$$

the complex amplitude of the wave of molecular vibrations of the medium is related to the field by the following expressions:

$$Q=\int_{-\infty}^{\infty} q(\tau,z,\omega_0)G(\omega_0-\Omega)\,d\omega_0,$$

$$q(\tau,z)=\frac{1}{T_2}\exp\left\{-\left(\frac{1}{T_2}+i\Delta\right)\tau\right\}\int_{-\infty}^{\tau} E_p(z,t)E_c^*(z,t)\times\exp\left\{\left(\frac{1}{T_2}+i\Delta\right)t\right\}dt, \quad (3)$$

where $\Delta=\omega-\omega_0$, ω_0 is the frequency of a vibrational transition in an individual molecule, and T_2 is the transverse relaxation time defining the homogeneous width of the spontaneous scattering line.

The broad SRS lines in quartz optical fibers can be used, in principle, to generate subpicosecond SRS pulses. Below, we shall consider the generation of SRS pulses in an optical fiber in the "long" pump pulse limit, assuming that the condition $\tau_p \gg T_2$ is satisfied. The nonlinear part of the polarization due to SRS can then follow quasistatically the space-time variation of the pulse envelopes at the fundamental and Raman generation frequencies, and the original set of equations (1)–(3) assumes the following form in terms of dimensionless variables:

$$i\frac{\partial E_p}{\partial z}=-\frac{P_p}{2}\frac{\partial^2 E_p}{\partial\tau^2}+R_{pp}|E_p|^2E_p+R_{ps}|E_s|^2E_p-i\frac{\omega_p}{\omega_s}|E_s|^2E_p-i\gamma_pE_p, \quad (4)$$

$$i\left(\frac{\partial E_s}{\partial z}+\mu\frac{\partial E_s}{\partial\tau}\right)=-\frac{P_s}{2}\frac{\partial^2 E_s}{\partial\tau^2}+R_{ss}|E_s|^2E_s+R_{sp}|E_p|^2E_p+i|E_p|^2E_s-i\gamma_sE_s+N_s(z,\tau), \quad (5)$$

where the variables are normalized as follows: $E_{p,s}=E_{p,s}/|E_{p0}|$; $\tau=\tau/\tau_{p0}$, where E_{p0} and τ_{p0} are, respectively, the amplitude and length of the incoming pump pulse, $z=z/z_g$,

$$z_g=2/g\alpha_{ps}|E_{p0}|^2=cn_0/4\pi g\alpha_{ps}I_{p0}$$

is the nonlinear SRS amplification length, and

$$P_p=\operatorname{sign}\left(\frac{\partial^2k_p}{\partial\omega^2}\right)\frac{z_g}{z_g},$$

$$P_p=\left|\frac{\partial^2k_s}{\partial\omega^2}\right|\left|\frac{\partial^2k_p}{\partial\omega^2}\right|^{-1}P_p\operatorname{sign}\left(\frac{\partial^2k_s}{\partial\omega^2}\right),$$

$$R_{m,n}=\frac{z_g}{z_{nl}}\alpha_{m,n}\frac{k_m}{k_n},\quad \mu=\frac{z_g}{z_{\rm coh}},\quad z_d=\tau_{p0}^2\left(\frac{\partial^2k_p}{\partial\omega^2}\right)^{-1},$$

$$z_{nl}=z_g\frac{n_0g}{n_2k_p},\quad z_{\rm coh}=\tau_{p0}\left(\frac{1}{v_s}-\frac{1}{v_p}\right)^{-1},$$

where z_d, z_{n1}, and $z_{\rm coh}$ are, respectively, the dispersion, nonlinear, and group-delay lengths of the pulses.

When the difference between the group velocities of the pulses at the fundamental and Raman generation frequencies can be neglected ($v_p=v_s$), the phase structure of the interacting waves can be described to first order in dispersion theory by the following simple relationships for $n_2(\omega_p)\equiv n_2(\omega_s)$, $\alpha_{mn}=\alpha$, and $\gamma_{p,s}=0$:

$$\varphi_p(z,\tau)=\varphi_p(0,\tau)-2R_{pp}zF(z,\tau), \quad (6)$$

$$\varphi_s(z,\tau)=\varphi_s(0,\tau)-2R_{ss}zF(z,\tau), \quad (7)$$

where $R_{ss}=(k_s/k_p)R_{pp}$, and

$$F(z,\tau)=|E_p(0,\tau)|^2+\frac{\omega_p}{\omega_s}|E_s(0,\tau)|^2+\frac{\omega_s-\omega_p}{\omega_p}\ln\left\{\left[|E_p(0,\tau)|^2+\frac{\omega_p}{\omega_s}|E_s(0,\tau)|^2\exp\left[2z\left(|E_p(0,\tau)|^2+\frac{\omega_p}{\omega_s}|E_s(0,\tau)|^2\right)\right]\right]\left(|E_p(0,\tau)|^2+\frac{\omega_p}{\omega_s}|E_s(0,\tau)|^2\right)^{-1}\right\}.$$

It is clear from (6) and (7) that, in the absence of group velocity dispersion, the dynamics of the instantaneous frequency spectrum in the pulses

$$\omega(t)=\omega+\Delta\omega(t),\quad \Delta\omega(t)=-\partial\varphi/\partial\tau$$

is determined by the sum of intensities at the fundamental and Raman frequencies as a function of time. Since, to first order in dispersion theory, the sum of the instantaneous intensities

$$|E_p(z,\tau)|^2+\frac{\omega_p}{\omega_s}|E_s(z,\tau)|^2=|E_p(0,\tau)|^2+\frac{\omega_p}{\omega_s}|E_s(0,\tau)|^2$$

is a constant of motion for (4) and (5), it is readily seen that the transformation of the Stokes components into one another involves transferring the pump phase modulation to the Stokes frequency, in connection with which no discontinuities arise in the phase or frequency structure, and frequency discontinuities are associated only with the generation of the Stokes components of SRS. In contrast to this case, allowance for the difference between the group velocities v_p and v_s leads to qualitatively new features in the structure of phase and frequency modulation of pump and SRS pulses. If we look upon the term containing the parameter μ on the left-hand side of (5) as a small perturbation, we can readily show that self-phase modulation will be determined in this case not only by the profile, but also by the steepness of the leading edges of the interacting pulses. As the pump pulse energy is transferred to the Raman-frequency pulse in the region of strong energy transfer between neighboring components, an intensity burst is produced on the leading ($v_s > v_p$, positive group velocity dispersion) and trailing edges of the SRS pulse, whose height may exceed the pump-pulse amplitude. This type of self-steepening of the leading ($v_s > v_p$) and trailing ($v_s < v_p$) Raman pulse edges causes an additional jump to appear in the phase and, consequently, in the frequency within the structure of the pump pulse.

Figure 5 shows the nonlinear evolution of the narrow intensity bursts $\Delta\omega_p(z,\tau)$ in the frequency modulation of the pump pulse $\omega_p(z,\tau) = \omega_p(0) + \Delta\omega_p(z,\tau)$, calculated from (4), (5) on the assumption that the initial Stokes wave $N_s(z,\tau)$ was strictly monochromatic: $N_s(z,\tau) = \text{const}, E_{sp}$, where E_{sp} is the spontaneous scattering amplitude in the optical fiber. This idealization enables us to follow in detail the evolution of frequency modulation in pulses at the fundamental and Raman frequencies, which arises from the combined effect of self-phase modulation and dispersion of group velocity of pulses during SRS. In the weak scattering region, and so long as the SRS intensity is low enough: $I_s(z,\tau) \ll I_p(0,\tau)$, the exponential nature of the gain at the Raman frequency is responsible for the appearance of a pulse of length $\tau_s(z_{th}) = \tau_p(z=0)G_{th}^{-1/2}$, where G_{th} is the threshold increment on the SRS gain (Fig. 5b). In this SRS regime, we have a linear increase with increasing distance in the frequency scanning range within the pump pulse that overlaps the SRS pulse. At the end of the linear regime of SRS generation, $z = z_{th}$ (where $z_g z_{th} = 16/gI_{p0}$; Ref. 23), the spectral broadening of the pump pulse (for $\mu \ll 1$) is determined by the following relation among the parameters of the problem:

$$\Delta\omega\tau_{p0} = 2z_{th}\, z_g/z_{nl} = 2z_{th}\, n_2(\omega_p)\, k_p/n_0 g.$$

We note that, since in the region of linear scanning ($0 < z \leqslant z_{th}$) the frequency sweep is determined by the time structure of the pump pulse, the frequency within the entire SRS pulse varies linearly for $z = z_{th}$ (Fig. 5b). Hence, the SRS pulse can be used for the subsequent compression in a medium with negative group velocity dispersion.

As the pump pulse continues to propagate in the optical fiber, an energy transfer avalanche begins between the waves for sufficiently large scattering lengths $z > z_{th}$, and is accompanied by the burning of a hole in the central part of the pump pulse (Figs. 5c–f; see also Fig. 2). The frequency scanning law is then no longer linear (Figs. 5c–f). We emphasize that this effect is determined by the presence of a difference between the pulse group velocities ($v_s \neq v_p$). In the case of nonlinear scattering (Figs. 5c–f), the leading edge of the SRS pulse steepens (we are considering the spectral region of SRS excitation corresponding to ($\partial v/\partial\lambda > 0$) as it enters the undepleted portion of the pump pulse, the sharp spikes appear in the structure of the frequency-modulated scattered and transmitted pulses. As the pulse continues to propagate in the optical fiber, the width of the valley in the central part of the pump pulse increases, and the region of sharp changes in frequency modulation becomes concentrated near the leading and trailing edges of the pump and Raman

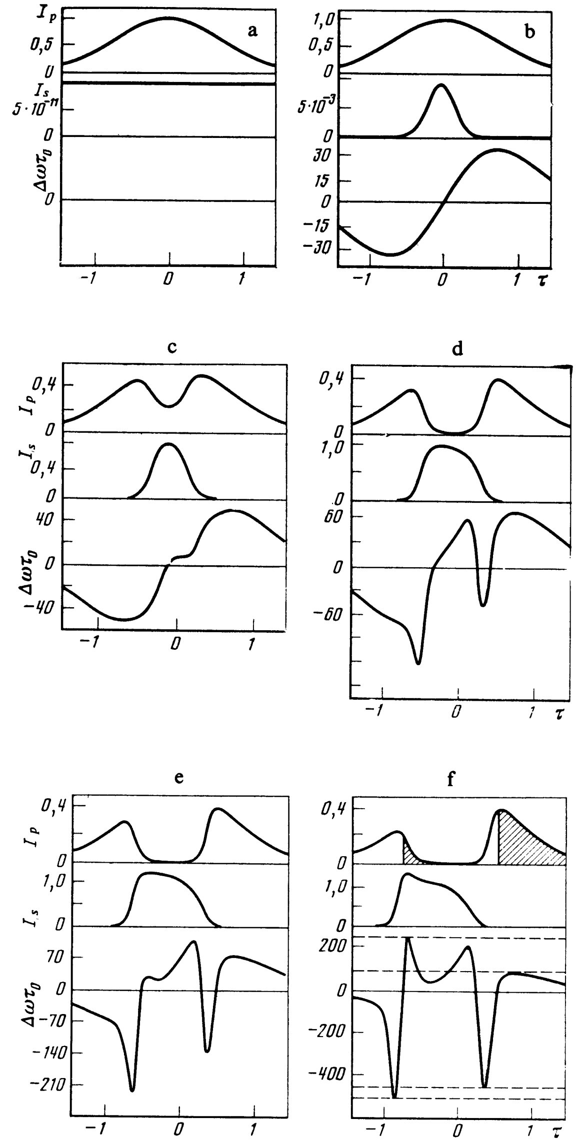

FIG. 5. Evolution of SRS pulses in a single-mode optical fiber pumped by a Gaussian pulse. Figures show the time envelopes of the pump pulse $I_p(\tau)$, the first Stokes component of SRS $I_s(\tau)$, and the variation in the instantaneous frequency of the pump field $\omega_p(z,\tau) - \omega_p(0) = \Delta\omega_p(z,\tau)$ for the following values of the dimensionless parameter $z = z/z_g$: 0 (a), 8 (b), 12 (c), 16 (d), 20 (e), and 32 (f). The main parameters of the problem are: $\mu = 2\times10^{-2}$; $\theta = z_{\rm coh}/z_d = [(\omega_p - \omega_s)\tau_0]^{-1}$; $(\omega_p - \omega_s)/2\pi s = 440\ \text{cm}^{-1}$; $\tau_0 = 60$ ps; $\partial^2 k/\partial\omega^2 = 2\times10^{-28}\ s^2/\text{cm}$, $n_2 = 3.2\times10^{-16}\ \text{cm}^2/\text{W}$; $g = 0.92\times10^{-11}$ cm/W. In an optical fiber in which the polarization is not constant, the parameters n_2 and g are replaced with averages over the polarization: $5n_2/6$ and $g/2$. The shaded areas in the last figure correspond to positive bursts in the frequency modulation structure of the pump pulse $\Delta\omega_p(z,\tau) \geqslant 0$. Horizontal lines in the lower part of this figure show the spectral regions of variation of the pump pulse frequency in which single pulses can be selected by spectral filtering.

pulses (see Figs. 5d–f). In Fig. 5f, the shaded areas on the pump pulse correspond to positive and negative bursts of frequency modulation. It follows from our calculations (cf. Fig. 5f) that spectral filtering with this kind of nonlinear frequency variation within the pump pulse in the anti-Stokes region relative to the pump ($\Delta\omega(z,\tau) > 0$) can be used to select a single pulse at a frequency shifted relative to ω_p

whereas, in the Stokes region ($\Delta\omega(z,\tau) < 0$), both negative bursts of frequency modulation of the pump pulse contain sufficient energy for two pulses to be selected at a frequency shifted toward the Stokes region.

In the general case, the SRS process in an extended channeling medium (optical fiber) must be described with allowance for noise sources that are distributed along the path and describe spontaneous Raman scattering.[25] Since spontaneous scattering is, by its very nature, a Gaussian fluctuation process, the SRS intensity will fluctuate and contain spikes with characteristic correlation time determined by the spontaneous scattering linewidth. The initial stage of SRS generation is dominated by regenerative narrowing of the spectrum and dispersive broadening of fluctuation bursts of SRS. As the pump intensity approaches the threshold value for the generation of SRS, the Raman-frequency radiation is found to contain individual fluctuation intensity spikes corresponding to the maximum of the pump pulse. The evolution of the Raman-frequency pulses during the linear stage of amplification of spontaneous RS noise is essentially analogous to that which we have obtained in numerical experiments on generation dynamics in the Raman fiber optic laser[26] where, however, we did not take into account effects connected with the dependence of the refractive index of the fiber material on radiation intensity which, as was shown in Ref. 20, can substantially modify the overall picture of SRS and of the optical fiber.

In the present research, we have carried out numerical experiments on the dynamics of SRS generation in an optical fiber, beginning with spontaneous RS. In contrast to Ref. 26, we have taken into account effects due to both SRS and the dependence of the fiber refractive index on radiation intensity. The calculations were performed in a wide range of parameter values, including the dispersion of group velocities of interacting pulses. The aim of the calculations was to provide a physical explanation of the effects described in the experimental part of this work.

Figures 6 and 7 show these numerical calculations and illustrate the development of SRS generation in a single-mode optical fiber. The numerical experiment enables us to show how individual physical processes described by (4) and (5) affect the evolution of the temporal and spectral structure of pulses during SRS in an optical fiber. These "numerical oscillograms" compare the temporal and frequency characteristics of pump pulses and the first two Stokes components of SRS, computed for different values of the parameter μ. Actually, this formulation of the problem corresponds to the experimental situation in which SRS generation in the optical fiber is due to a frequency-tunable pump pulse. Moreover, the dispersion in the optical fiber can be controlled by suitably choosing the refractive-index profile.[27] Comparison of Figs. 6 and 7 illustrates the dependence of SRS generation on group velocity dispersion and self-phase modulation.

In the region of low-intensity scattering we have $I_s(z,\tau) \ll I_{p0}$ and we can see individual groups of pulses at the Raman frequency with the envelope $\tau_{gr}(z) \sim \tau_{p0}/G_{th}^{1/2}$ (Fig. 6a; cf. Fig. 5). The nonlinear dynamics of pulses inside the envelope $\tau_{gr}(z)$ is determined by the exponential nature of Raman amplification: $I_s(z,\tau) \sim \exp[gzI_p(z,\tau)]$. As the length of the fiber increases, there is an increase in the Raman gain. For intensities $I_p \sim I_s$, narrow valleys appear in the central portion of the pump pulse and correspond to

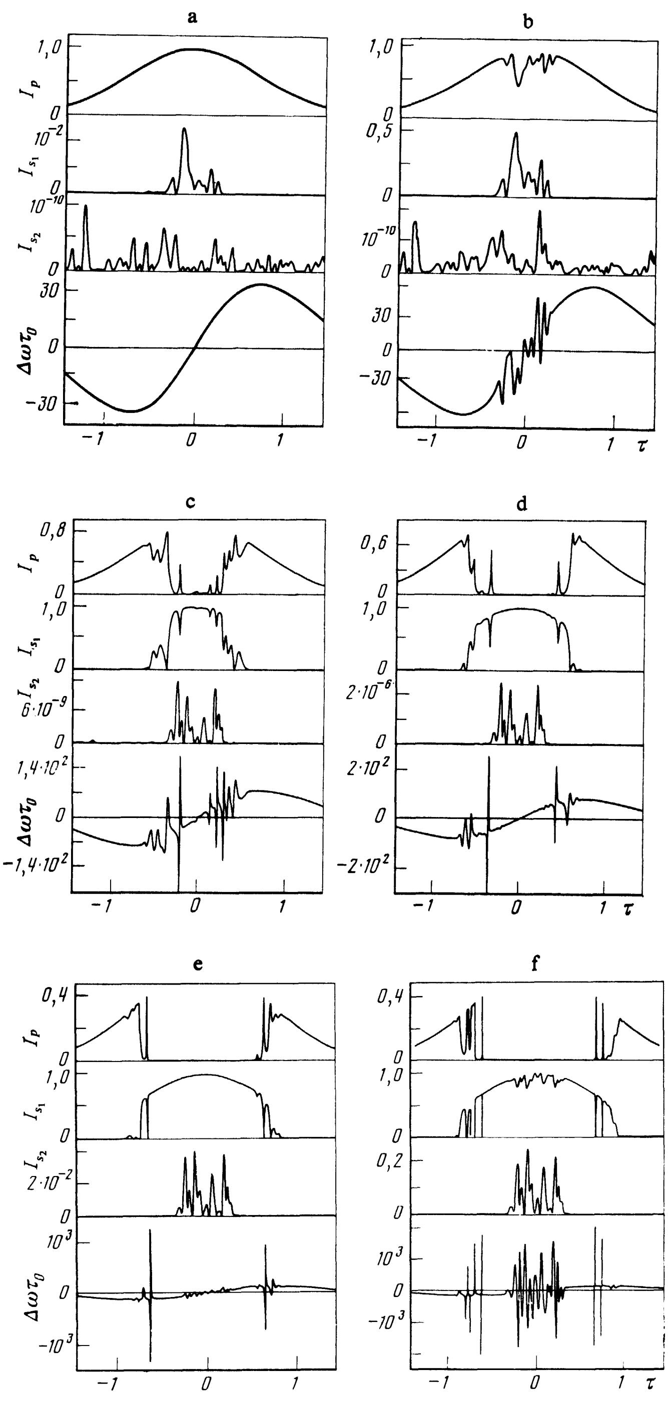

FIG. 6. "Numerical oscillograms" illustrating the nonlinear evolution of the first two Stokes components of SRS I_s, I_{ss} in an optical fiber and the distortion of the frequency sweep law within the pump pulse $\Delta\omega_p(z,\tau)$ for the following values of the dimensionless fiber length $z = z/z_g$: 8 (a), 10 (b), 12 (c), 16 (d), 18 (e), and 20 (f). The ratio of the lengths of coherent pulse interaction z_{coh}, SRS gain z_g, and dispersion z_d is as follows: $\mu = \mu_1 = z_g/z_{coh} = 10^{-4}$; $\theta = z_{coh}/z_d = 2\times10^{-4}$.

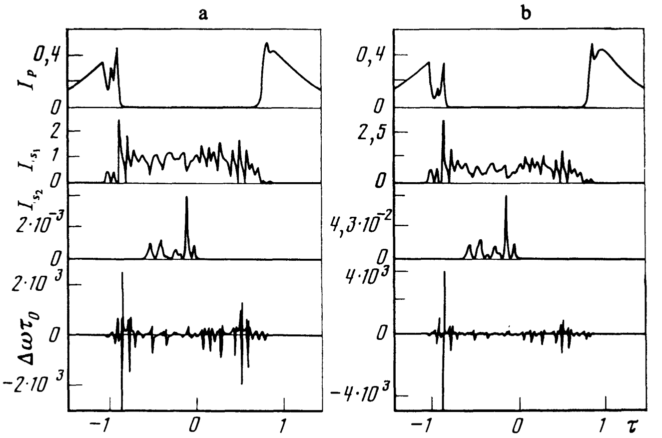

FIG. 7. "Numerical oscillograms" of SRS generation dynamics in an optical fiber, and of the nonlinear dynamics of frequency scanning in the pump pulse, $\Delta\omega_p(z,\tau)$, computed for $\mu = \mu_2 = 100\mu_1$; $\theta = 2\times10^{-4}$ km, and $z = 27$ (a) and 32 (b).

fluctuation spikes at the Stokes frequency (Fig. 6b). This instant of time is critical. Firstly, the "mapping" of the fluctuational time structure, which corresponds to SRS, on the pump pulses produces a sharp broadening of the pump pulse spectrum and leads to the appearance of a broad pedestal. Secondly, considerable distortions of the frequency modulation structure of the pump pulse appear at this time. This effect prevents further compression of the pump pulse in dispersive compression systems (Figs. 6d–f). As the intensity of the first Stokes component increases, pulses also appear at the frequency of the second component, as shown in Figs. 6c and d. The noise structure of the second component is then "reflected" in the first, as the central portion is "eaten up" (Figs. 6e–f). New noise spikes appear at this time in the frequency modulation of the pump, and the spectrum of the first Stokes component is additionally broadened.

Figure 7 shows the qualitatively new features of the time envelope and the frequency modulation structure of the pump, which arise as a result of the group velocity dispersion effect at the fundamental and Raman generation frequencies. Figures 7a–b show the main stages of evolution of SRS generation in this case. It is clear that an increase in the group velocity difference leads to fundamentally new features in the structure of the time envelopes of pulses and in frequency modulation. First, a narrow pulse with a sharp leading edge and intensity exceeding that of the pump appears at the Stokes frequency. Second, as the fiber length increases, a narrow spike corresponding to the leading "splinter" of the pump is formed on the phase modulation structure of the pump. Numerical experiments were performed for a large number of forms of the random seeding fields describing the spontaneous scattering at Raman frequencies and show that, in all cases, a similar frequency structure of the pump pulse is formed. This in turn means that it is possible to select experimentally a single ultrashort pulse, at a frequency shifted relative to the pump, by exploiting the spectral filtering of the pump radiation (see Fig. 4b) whereas for $\mu_1 = 10^{-4} = 10^{-2}\mu_2$ (cf. Figs. 6f and 7b) this selection is, in general, impossible.

We note that, in actual experiments, ultrashort pulse selection in the Stokes region relative to the pump pulse is prevented by the overlap of the wing of the Stokes component of SRS and the pump spectrum (Fig. 6).

The above results were obtained on the assumption that the total frequency scanning interval within the pump pulse is less than the width of the spontaneous Raman spectrum. If this were not so, self-phase modulation effects and group velocity dispersion of pulses channeled by the optical fiber may, in general, suppress the SRS process.[20,28]

Laser pulses are often not spectrally bounded in real physical experiments. We have examined the dynamics of SRS generation in an optical fiber pumped by partially coherent pulses of given shape and random time and phase substructure:

$$E_p(z=0,\ \tau)=E_{p0}(\tau)A(\tau)\ \exp\left[i\varphi_{p0}(\tau)\xi(\tau)\right],$$

where $E_{p0}(\tau)$ and $\varphi_{p0}(\tau)$ are the smooth pulse and phase envelopes, and $A(\tau)$ and $\xi(\tau)$ are random complex functions. We have investigated the basic dynamics of SRS generation in an optical fiber in the field of a Gaussian pulse $E_{p0}(\tau) = E_0\exp(-\tau^2/2)$ with a random time substructure and a random phase modulation corresponding to the pulse profile of the form $\varphi(\tau) = E_0^2\exp(-\tau^2)$ and $\xi(\tau) = |A(\tau)|^2$. The exponential dependence of SRS gain on pump intensity can lead to a considerable simplification of the time structure of radiation during the multistage generation of Stokes components.[26] An asymmetric Raman pulse can appear on the time structure at the Stokes frequency during the development of the generation process, and can burn up the pump during propagation in the optical fiber. This is accompanied by dramatic simplification of the time structure of the frequency modulation of the pump pulse. A strong, narrow SPM spike appears in the anti-Stokes region and corresponds to an ultrashort SRS pulse whereas, in the Stokes region, there are two frequency modulation bursts corresponding to steep leading edges of SRS pulses and the burnt-through portion of the pump pulse. This result of the numerical experiment leads to the conclusion that the spectral filtering of radiation during stimulated Raman scattering in an optical fiber can be used to select single ultrashort pulses even when the pump is an ordinary multimode laser with a fluctuating time structure at the cavity frequency.

We now summarize our main results.

1. We have identified the physical mechanism responsible for the generation of the intense continuum produced when optical fibers are pumped by lasers. The treatment that we propose is based on the combined operation of stimulated Raman scattering, self-phase modulation, and group velocity dispersion velocity when SRS is seeded by spontaneous noise.

2. Spectral filtration of the continuum can be used to generate ultrashort pulses at a frequency shifted relative to the pump.

We note that, if simultaneously with the pump, we introduce radiation of arbitrary frequency that does not participate in the SRS process, the phase cross modulation effect will ensure that our method of spectral filtration can be used to select ultrashort pulses of any wavelength convenient for subsequent amplification, which provides us with frequency-tunable, powerful ultrashort pulses.

3. The codes we have developed enable us not only to calculate the nonlinear evolution of SRS pulses, but also to use experimental data on pulse lengths obtained by the method of spectral filtering to solve the inverse problem and determine the mean length of fluctuations at Raman frequencies. The latter is an independent problem in the case of disordered media with considerable inhomogeneous broadening.

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