Direct measurement of the nonlinear phase shift between the orthogonally polarized states of a single-mode fiber

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We measure the phase shift induced by the optical Kerr effect between the two orthogonally polarized states of a birefringent single-mode fiber. The associated noise, which can arise whenever amplitude fluctuations of the source are present, is discussed.

The optical Kerr effect, which is associated with a nonlinear contribution to the refractive index proportional to the intensity of the propagating field,[1] is responsible for a number of nonlinear processes whose practical relevance has become more apparent in the frame of optical-fiber propagation. Basically, the nonlinear refractive index induces an intensity-dependent self-phase modulation of the propagating field, which in turn can give rise to a number of associated phenomena, for example, frequency broadening[2] and pulse reshaping.[3,4] In particular, the Kerr effect in fibers has led to the development of devices such as ultrafast optical shutters,[5] intensity discriminators,[6,7] and soliton lasers.[8] More precisely, the shutter and the discriminator are based on the nonlinear phase shift suffered by the two orthogonally polarized states of a single-mode birefringent fiber and are due either to the presence of a strong pump field directed along one of the fiber principal axes (shutter) or to the different powers carried by the two modes in the absence of pump (discriminator).

We have directly measured the nonlinear phase shift, which is responsible in particular for the mechanism of intensity discrimination named above, at the relatively low powers associated with a cw laser source. This power-dependent birefringence can actually act as a source of noise in an intensity-modulated communication system.[9]

The experimental setup is shown in Fig. 1. The light from a cw linearly polarized argon laser emitting a power of about 1.3 W at λ = 514.5 nm is coupled to a polarization-maintaining birefringent fiber[10] about 60 m long, with a 2.8-μm-diameter d, an index difference Δn = 0.0089, a nominal loss at 514.5 nm of 40 dB/km, and a beat length at the same wavelength of 1.57 cm.[5] Although the normalized frequency assumes a value V = 2.76, the excitation conditions are such as to provide single-mode operation. The nonlinear self-induced phase difference between the two polarization states increases with the difference between the corresponding powers (in fact, for an ideal circular fiber, equal powers would imply no effect for obvious symmetry reasons, so that, for slightly elliptical fibers, the same condition would give rise to a small effect). In view of this fact, we excite the two states with a power ratio 25:1. After the light source, there is a chopper whose slits are filled with a variety of neutral-density filters to modulate the power. This arrangement allows us to excite the fiber with variable powers, and, because of the low duty cycle employed, prevents undesired thermal effects (the fiber is also enclosed in a thermally controlled cavity). Part of the light output is sent to a calibrated silicon cell in order to determine the exit power. Since the nonlinear effect is observed through the beat signal between the two orthogonal fields, their amplitudes after the fiber exit are equalized by means of a linear polarizer in order to improve the sensitivity of the measurement (this is accomplished by rotating the analyzer axis by an angle $\pi/2$ with respect to the polarization direction of the input field). The far-field intensity after the polarizer is detected at a fixed position **R** by an avalanche photodiode (APD) whose output is displayed on an oscilloscope. The field at **R** is the sum of two contributions:

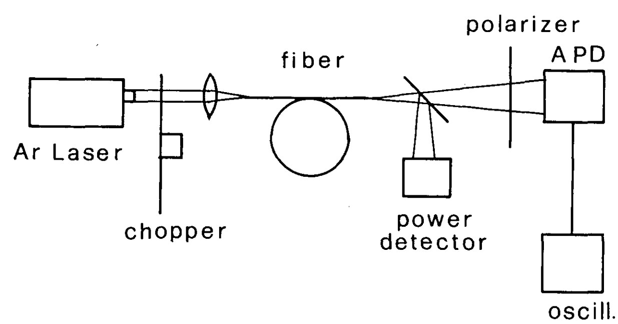

Fig. 1. Experimental setup.

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$$\mathbf{E}_1(\mathbf{R}, t) = \mathbf{E}_0(\mathbf{R}) \exp[i(\omega_0 t + \theta_1{}^L + \theta_1{}^{NL})],$$
$$\mathbf{E}_2(\mathbf{R}, t) = \mathbf{E}_0(\mathbf{R}) \exp[i(\omega_0 t + \theta_2{}^L + \theta_2{}^{\mathrm{NL}})], \quad (1)$$

where $\theta_i{}^L$ and $\theta_i{}^{\mathrm{NL}}$ represent, respectively, the linear and the nonlinear contributions to the phase of the ith field. The associated intensity is of the form

$$I(\mathbf{R}, T; P) = I_0\{1 + \cos[\Delta\theta_T{}^L + \Delta\theta^{\mathrm{NL}}(\mathrm{P})]\}, \quad (2)$$

where T indicates the parametric dependence on temperature associated with slow environmental changes, $\Delta\theta^L = \theta_1{}^L - \theta_2{}^L$, and $\Delta\theta^{\mathrm{NL}}(P) = \theta_1{}^{\mathrm{NL}} - \theta_2{}^{\mathrm{NL}}$. Both I_0 and $\Delta\theta^{\mathrm{NL}}$ depend on the injected power P, that is, on the particular filter that the beam crosses when measuring $I(\mathbf{R}, T; P)$, while $\Delta\theta_T{}^L$ is obviously independent of P. Equation (2) furnishes $\Delta\theta_T{}^L + \Delta\theta^{\mathrm{NL}}(P)$ as a function of $I(\mathbf{R}, T; P)/I_0(P)$. In order to get rid of the undesired quantity $\Delta\theta_T{}^L$, we proceed as follows. When we let the cavity temperature change slowly with time, the linear phase $\Delta\theta_T{}^L$ undergoes variations slow with respect to the chopper rotation period. Consequently, the values of $I(\mathbf{R}, T; P)$ displayed on the oscilloscope that correspond to various powers of P fed into the fiber during a chopper turn refer to a common value of $\Delta\theta_T{}^L$. On the other hand, $I(\mathbf{R}, T; P)$ ranges, for any fixed value of P, from zero to $2I_0(P)$, so that the variations described above allow us to obtain the values of $I_0(P)$ for any P. We now choose the set of measured values $I(\mathbf{R}, \Psi; P)$ corresponding to the chopper turn in which $I(\mathbf{R}, \Psi; P_m) = 0$, Ψ being the temperature for which this equation is satisfied and P_m the power launched through the most-attenuating filter. For this set of measurements, Eq. (2) implies that

$$\Delta\theta_\Psi{}^L + \Delta\theta^{\mathrm{NL}}(P_m) = \pi \quad (3)$$

and

$$I(\mathbf{R}, \Psi; P) = I_0(P)[1 - \cos \Delta\phi(P)], \quad (4)$$

where $\Delta\phi(P) = \Delta\theta^{\mathrm{NL}}(P) - \Delta\theta^{\mathrm{NL}}(P_m)$. Inverting Eq. (4) allows us to determine $\Delta\phi$ as a function of P.

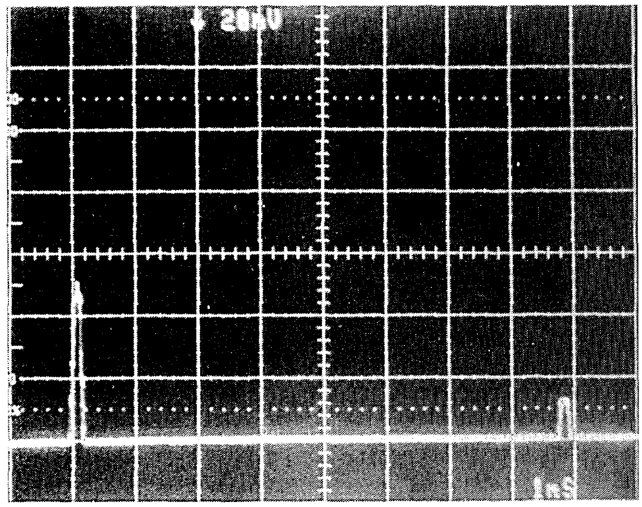
In Figs. 2a and 2b two particular scope displays are shown, illustrating $I(\mathbf{R}, \Psi; P_m)$ and $I(\mathbf{R}, \Psi; P_M)$ (P_M indicates the power through the least-attenuating filter). Figure 2a refers to a generic temperature, while Fig. 2b refers to $T = \Psi$; the presence of the nonlinear effect is shown by the nonvanishing of $I(\mathbf{R}, \Psi; P_M)$. The experimental values of $\Delta\phi$ follow a linear behavior as a function of $\Delta P = P - P_m$, as is clearly seen from Fig. 3, where they are shown as a function of the exit power $\Delta P_{\mathrm{out}} = P_{\mathrm{out}} - P_{m,\mathrm{out}}$, together with the best-fitting straight line.

The above results can be interpreted by investigating the evolution of the two copropagating polarization states in the presence of Kerr effect, according to which the refractive index n consists of a linear part n_1 and a nonlinear one:

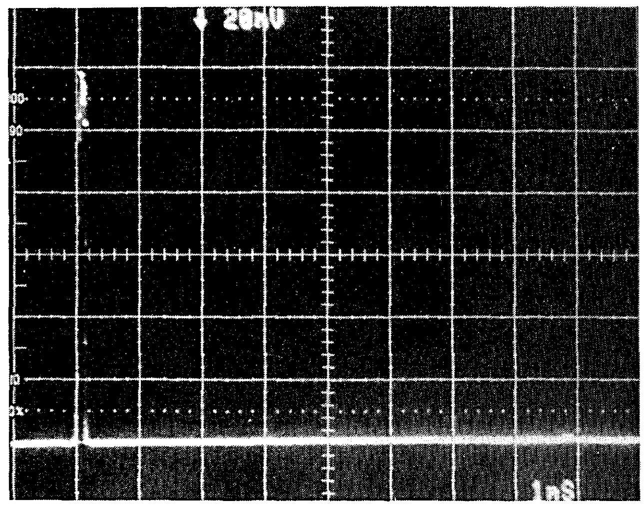
$$n = n_1 + n_2|E|^2, \quad (5)$$

n_2 being the nonlinear refractive-index coefficient. To this end, let us write the electric field inside the fiber as

$$\mathbf{E}(\mathbf{r}, z, t) = E_1(\mathbf{r}) \exp(i\omega_0 t - i\beta_1\omega_0 z)\Phi_1(z, t)\hat{x} + E_2(\mathbf{r}) \exp(i\omega_0 t - i\beta_2\omega_0 z)\Phi_2(z, t)\hat{y}, \quad (6)$$



a



b

Fig. 2. Scope display of (left) the maximum and (right) the minimum output intensities: a, at a generic temperature T; b, at T = T.

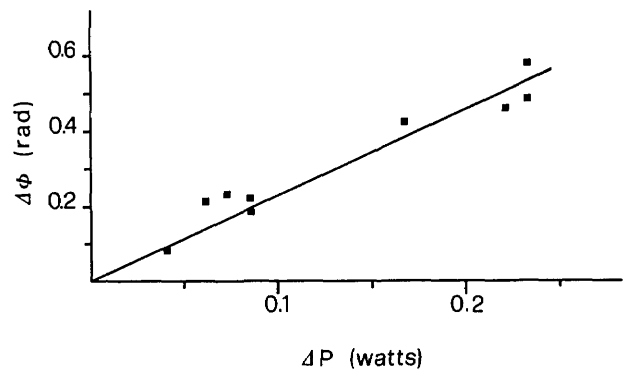


Fig. 3. Nonlinear phase shift $\Delta\theta$ as a function of $\Delta P = P - P_m$.

where $E_i(\mathbf{r})$ is the transverse configuration of the ith state and β_i is its propagation constant. By generalizing a previous result worked out for modes possessing the same linear polarization,[11] it is possible to obtain a set of coupled equations describing the evolution of the amplitudes $\Phi_1(z, t)$ and $\Phi_2(z, t)$, which reads, neglecting chromatic dispersion, as

$$\left\{\frac{\partial}{\partial z} + \frac{1}{V_1}\frac{\partial}{\partial t}\right\} \Phi_1 = -i\{R_{11}|\Phi_1|^2 + \tfrac{2}{3}R_{12}|\Phi_2|^2\}\Phi_1,$$
$$\left\{\frac{\partial}{\partial z} + \frac{1}{V_2}\frac{\partial}{\partial t}\right\} \Phi_2 = -i\{R_{22}|\Phi_2|^2 + \tfrac{2}{3}R_{12}|\Phi_1|^2\}\Phi_2, \quad (7)$$

where $V_i^{-1} = \mathrm{d}\beta_i/\mathrm{d}\omega$ evaluated at $\omega = \omega_0$ and

$$R_{ij} = \frac{\omega_0 n_2}{c} \iint E_i^2(\mathbf{r}) E_j^2(\mathbf{r}) \mathrm{d}x\,\mathrm{d}y. \quad (8)$$

The set of Eqs. (7) can be easily solved,[11] and the nonlinear phase difference $\Delta\theta^{\mathrm{NL}}(P)$ between Φ_1 and Φ_2, which is found by evaluating $\Phi_1\Phi_2^*$, turns out to be, in MKS units,

$$\Delta\theta^{\mathrm{NL}}(P) = \tfrac{2}{3}k(n_2/n_1)Z_0PL/A, \quad (9)$$

where $k = \omega_0/c$, Z_0 is the vacuum impedance, A is (approximately) the common effective area of the two orthogonal configurations, and L is the fiber length.

We can now compare the theoretical values $\Delta\phi(P) = \Delta\theta^{\mathrm{NL}}(P) - \Delta\theta^{\mathrm{NL}}(P_m)$ obtained from Eq. (9) with the experimental ones reported in Fig. 3, provided that the fiber length L is replaced by the effective length $L_{\mathrm{eff}} = [1 - \exp(-\alpha L)]/\alpha$ (where α is the attenuation coefficient) and the input power P is obtained from the measured exit power P_{out} by correcting for losses and Fresnel reflection according to the relation $P = P_{\mathrm{out}} \exp(\alpha L)/0.96$. If we write $n_1 = 1.46$ and $A = \pi(d/2)^2$ (which is a good approximation for the assumed value of the normalized frequency V), we find satisfactory agreement between the theoretical expression and the experimental data for a value of the nonlinear refractive-index coefficient $n_2 \cong 0.9 \times 10^{-22}$ (m/V)^2. This value (multiplied by a factor of 2 in order to compare it with those available in the literature, where n_2 is usually defined through the relation $n = n_1 + n_2|\mathbf{E}|^2/2$) is 1.6 $\times 10^{-13}$ esu (shifting to CGS units) and appears to be in reasonable agreement with those obtained by means of other methods.[2]

We wish finally to mention the role that the nonlinear phase shift can play as a source of noise, whenever the exciting field undergoes amplitude fluctuations (or variations) and a polarization-dependent loss mechanism is present in the fiber. By identifying, for simplicity, this mechanism with that associated with the presence of a linear-polarization analyzer at the fiber output, the power after the analyzer can be written as

$$P = P_1 \cos^2 \alpha + hP_1 \sin^2 \alpha + h^{1/2}P_1 \sin(2\alpha) \cos \Gamma, \quad (10)$$

where h is the small ratio between the power unavoidably present in polarization state 2 and that in state 1 (P_1), α the angle between $\hat{x}$ and the polarizer axis, and Γ is the phase delay between the two states. According to Eq. (9), Γ contains a nonlinear term following the time fluctuations (or modulations) of $P_1(\mathrm{t})$, which, for a long fiber link, can become relevant even at low power (for a telecommunication link 10 km long, $\Delta\theta^{\mathrm{NL}} \cong 1$ rad for P of the order of few milliwatts) and thus give rise to nonnegligible amplitude noise. The noise described is analogous to the phase noise generated in angle-modulated fiber communication systems by power fluctuations of the source.[12]

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