Optical pulse propagation and holographic storage in a coupled-resonator optical waveguide

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We propose a method of storage and reconstruction of a classical light pulse based on photorefractive holography in a coupled-resonator optical waveguide (CROW). Pulse propagation in a CROW is described in the context of the tight-binding approximation; the use of a CROW results in a large reduction of the group velocity, which is important for spatial compression of the optical pulses. Further, the efficiency of the photorefractive effect is enhanced in a CROW, enabling quasistatic holographic grating formation using much lower intensity optical pulses. We describe in detail the formation of a photorefractive index grating in a CROW via interference with a reference pulse and the subsequent holographic reconstruction of the signal pulse.

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Single pulse holographic recording requires strong nonlinear properties which are not available with ordinary materials, since the hologram has only a short time and a limited amount of energy in which to form. The highly nonlinear properties of ultracold atomic vapors form the basis of one approach which was demonstrated recently [1]. In this paper, we investigate the possibility of recording and reconstructing holograms of optical pulses using the recently proposed coupled-resonator optical waveguide (CROW) [2]. A CROW consists of an array of weakly coupled high-Q resonators, leading to very high optical intensities even at moderate (propagating) power levels—exactly what is required for holographic recording. In the paper, we will find that in spite of the discrete localization of an optical field at the individual resonators, it is still possible to reconstruct faithfully the signal pulse which is recorded in the hologram.

The particular structural realization of the CROW that we consider is shown in Fig. 1, in which the individual resonators consist of defect cavities embedded in a two-dimensional (2D) periodic structure (a 2D photonic crystal [3]). In our later discussion, the material of index n_2 will be assumed to be photorefractive (e.g., GaAs) and the material of index n_1 will be assumed to be air, for simplicity. Note that the defect cavity is then composed of photorefractive material, and the simultaneous presence of two optical fields in this region will induce a photorefractive index grating which can be used for holography. In order to describe the holographic storage and reconstruction of an optical pulse, we first need to understand how an optical pulse with a known free-space description propagates in a CROW, i.e., the form of the fields that write the index grating.

I. PULSE PROPAGATION IN A CROW

We briefly review the theory of optical pulse propagation in the simplest of waveguides—one described satisfactorily by a linear dispersion relationship [4]. The results and notation established here will be of importance in the subsequent discussion of pulse propagation in a CROW.

Consider an input pulse $\mathcal{E}(z=0,t)$ described by

$$\mathcal{E}(z=0,t)=e^{i\omega_0 t}E(z=0,t) \tag{1}$$

$$=e^{i\omega_0 t}\frac{1}{2\pi}\int d\Omega\,\tilde{E}(z=0,\Omega)e^{i\Omega t}, \tag{2}$$

where $\tilde{E}(z=0,\Omega)$ is the Fourier transform of the envelope $E(z=0,t)$:

$$\tilde{E}(z=0,\Omega)=\int dt\,E(z=0,t)e^{-i\Omega t}. \tag{3}$$

The field at a distance z, expressed as $\mathcal{E}(z,t)$, is obtained by multiplying each frequency component $(\omega_0+\Omega)$ by $\exp[-ik(\omega_0+\Omega)z]$,

$$\mathcal{E}(z,t)=e^{i\omega_0 t}\frac{1}{2\pi}\int d\Omega\,\tilde{E}(z=0,\Omega)e^{i\Omega t}e^{-ik(\omega_0+\Omega)z}, \tag{4}$$

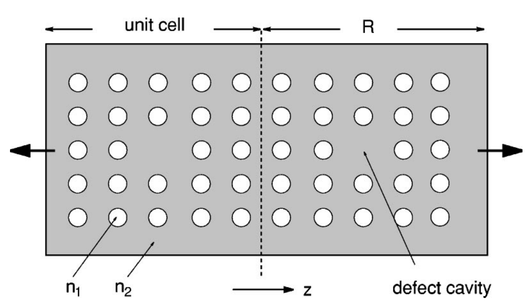

FIG. 1. Schematic of a photorefractive CROW realized by a coupling of the individual defect cavities in a 2D photonic crystal, comprised of a photorefractive dielectric medium with high refractive index n_2 and a nonphotorefractive dielectric medium of low refractive index n_1. The structure is periodic in the $\hat{z}$ direction, with a spatial period R.

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where $k(\omega)$ is the wave vector at the optical frequency ω. (We use the sign convention of Ref. [4].)

If we expand $k(\omega_0+\Omega)$ near ω_0 as

$$k(\omega_0+\Omega)=k(\omega_0)+\left.\frac{dk}{d\omega}\right|_{\omega_0}\Omega+\cdots\equiv k_0+\frac{1}{v_g}\Omega+\cdots, \tag{5}$$

where v_g is the group velocity of the pulse, and substitute this relationship into Eq. (4), retaining terms up to the linear terms in Ω, we obtain

$$\begin{aligned}\mathcal{E}(z,t)&=e^{i(\omega_0 t-k_0 z)}\frac{1}{2\pi}\int d\Omega\,\tilde{E}(z=0,\Omega)e^{i\Omega(t-z/v_g)}\\&=e^{i(\omega_0 t-k_0 z)}E(z=0,t-z/v_g).\end{aligned} \tag{6}$$

This is the well-known result (Ref. [5], pp. 322–326) that a pulse propagates unchanged in shape in a weakly dispersive medium, apart from an overall phase factor, and that the velocity of propagation is given by the group velocity of the pulse v_g defined from the dispersion relationship as in Eq. (5).

As discussed by Yariv *et al.* [2], we can describe a CROW comprised of weakly coupled resonators with the tight-binding approximation commonly used in solid-state physics to describe electronic states in semiconductors with impurity doping (Ref. [6], Chap. 10). There are slight differences in the treatment presented here as compared to Ref. [2], in order to make the correspondence between the usual description of the tight-binding method in solid-state physics and the above description of pulse propagation more direct. The time-independent waveguide mode (eigenmode) of an infinitely long CROW $\phi_k(z)$ with wave vector k is a linear combination of the (normalized) high-Q modes $\psi(z)$ of a large number of identical resonators located along the z axis with inter-resonator spacing R:

$$\phi_k(z)=\sum_n e^{-iknR}\psi(z-nR). \tag{7}$$

We write k_0 as the wave vector corresponding to the center optical frequency ω_0 as defined by Eq. (5), and normalize the eigenvectors in a CROW comprised of M identical resonators:

$$M\int dz\,|\psi(z)|^2=1. \tag{8}$$

Equation (7) may be written in terms of Ω rather than k by expanding $k(\omega_0+\Omega)$ in a Taylor series near ω_0 as in Eq. (5):

$$\phi_\Omega(z)=\sum_n e^{-ik_0 nR}\psi(z-nR)e^{-i\Omega nR/v_g}. \tag{9}$$

For an input pulse of the form of Eq. (1), the field at a distance z, analogous to Eq. (4), is given by

$$\begin{aligned}\mathcal{E}(z,t)&=e^{i\omega_0 t}\frac{1}{2\pi}\int d\Omega\,\tilde{E}(z=0,\Omega)\phi_\Omega(z)e^{i\Omega t}\\&=e^{i\omega_0 t}\sum_n e^{-ik_0 nR}\psi(z-nR)\\&\quad\times\frac{1}{2\pi}\int d\Omega\tilde{E}(z=0,\Omega)e^{i\Omega[t-(nR)/v_g]}.\end{aligned} \tag{10}$$

The term on the last line of the above expression is merely the n-dependent shifted replica of the original input envelope; it follows from Eq. (2) that

$$\mathcal{E}(z,t)=e^{i\omega_0 t}\sum_n e^{-ik_0 nR}\psi(z-nR)E\left(z=0,t-\frac{nR}{v_g}\right). \tag{11}$$

In our description of the holographic process, we will need an expression for the spatial Fourier transform (in K space) of $\mathcal{E}(z,t)$,

$$\tilde{\mathcal{E}}(K,t)\equiv\int\mathcal{E}(z,t)e^{-iKz}dz=e^{i\omega_0 t}\sum_n I_n(K,t), \tag{12}$$

where

$$\begin{aligned}I_n(K,t)&\equiv e^{-ik_0 nR}\left[\int dz e^{-iKz}\psi(z-nR)\right]E\left(z=0,t-\frac{nR}{v_g}\right)\\&=e^{-i(K+k_0)nR}\tilde{\psi}(K)E\left(z=0,t-\frac{nR}{v_g}\right).\end{aligned} \tag{13}$$

Therefore,

$$\tilde{\mathcal{E}}(K,t)=e^{i\omega_0 t}\sum_n e^{-i(k_0+K)nR}\tilde{\psi}(K)E\left(z=0,t-\frac{nR}{v_g}\right). \tag{14}$$

Equations (11) and (14) form the main conclusions of this section, and describe pulse propagation in any medium (particularly a CROW) for which the eigenmodes are described by the tight-binding assumption [Eq. (7)], and the dispersion relationship is approximately linear as in Eq. (5). It is known that away from the band edges, such a dispersion relationship fits a CROW quite well [2]. The formulation of Eqs. (11) and (14) in a structure of finite length is discussed in Ref. [7].

II. PHOTOREFRACTIVE HOLOGRAPHY IN A CROW

If we design the CROW in a photorefractive medium, we can form a dynamic hologram of the signal pulse via interference with a reference pulse as shown in Fig. 2; the induced index grating persists in the photorefractive medium after the pulses have propagated away and contains all the necessary amplitude and phase information to (classically) reconstruct the signal pulse [8]. In this paper we will carry out an analysis for the scenario depicted in Fig. 2(a).

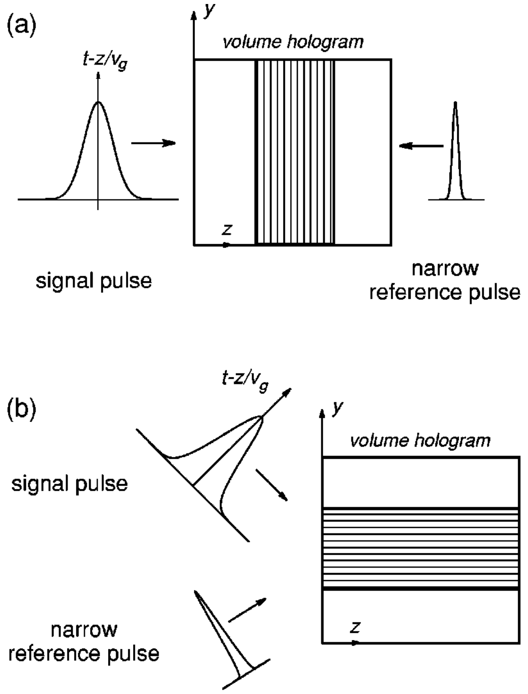

FIG. 2. Schematic diagram of the volume holography using pulse collision recording in (a) a counterpropagating geometry (reflection-type holography) and (b) a copropagating geometry (transmission-type holography). After Fig. 15.3 in Fainman *et al.* [10].

We will use the symbols t' and t to denote the temporal coordinates at the time of writing the hologram and at the time of reconstruction, respectively. In general, $t=t'-T$ for some time interval T. Based on the geometry shown, we assume that two pulses $\mathcal{E}_1(z=0,t')$ and $\mathcal{E}_2(z=L,t')$ are input at the two opposite ends of a CROW. The pulses propagate in opposite directions with wave vectors k_1 and $-k_2$ and group velocities v_1 and v_2 respectively. The total field $\mathcal{E}(z,t')$ is given by the sum of the fields due to these two pulses, and, in K space, can be written as

$$\tilde{\mathcal{E}}(K,t')=e^{i\omega_1 t'}\sum_n e^{-i(k_1+K)nR}\tilde{\psi}(K)E_1\left(z=0,t'-\frac{nR}{v_1}\right)+e^{i\omega_2 t'}\sum_m e^{-i(-k_2+K)mR}\tilde{\psi}(K)E_2\left(z=L,t'-\frac{mR}{v_2}\right), \tag{15}$$

analogous to the K-space representation of a single pulse, [Eq. (14)].

The holographic grating $\delta n(z,t')$ is produced by the interference pattern of the spectral components of this field (for example, at K_1 and K_2) weighted by a complex proportionality coefficient $\delta n(K_1,K_2)$ which represents the photorefractive coupling coefficient between two plane waves defined by K_1 and K_2. This coefficient depends on the material properties, the orientation of the medium, and the polarization of the waves [8,9]. Using the inverse Fourier transform to write the K space field [Eq. (15)] in terms of z, the grating is

$$\begin{aligned}\delta n(z,t')=&\frac{1}{F_0}e^{i(\omega_1-\omega_2)t'}\sum_{n,m}\int\frac{dK_1}{2\pi}\frac{dK_2}{2\pi}e^{-i(K_1-K_2)z}\\&\times\delta\hat{n}(K_1,K_2)\,e^{-i(k_1+K_1)nR}e^{-i(-k_2+K_2)mR}\\&\times E_1\left(z=0,t'-\frac{nR}{v_1}\right)E_2^*\left(z=L,t'-\frac{mR}{v_1}\right)\\&\times\tilde{\psi}(K_1)\tilde{\psi}^*(K_2)+\text{c.c.},\end{aligned} \tag{16}$$

where F_0 is the total optical power. We assume that this grating persists temporally, so that at a later time t, $\delta n(z,t)=\delta n(z,t')$ and we can relabel the temporal coordinate t' to t in Eq. (16) to describe the reconstruction process. We will only focus on the term in Eq. (16) written out in full, with the remark that analogous results hold for the complex conjugate term (written in Eq. (16) as c.c.)—this term will ultimately give rise to a field propagating in the direction opposite to that of $\mathcal{E}_1$, and is not of interest in this discussion.

We use a backward-propagating reference pulse $\mathcal{E}_r(z,t)$ to illuminate the grating, and preserve its distinction from the reference pulse at the time of grating formation $\mathcal{E}_2(z,t)$ to maintain the generality of this discussion. At a later stage, we will assume that these two pulses are in fact identical, and simplify the expressions appropriately. We can write $\mathcal{E}_r(z,t)$ in the Fourier domain using Eq. (14):

$$\tilde{\mathcal{E}}_r(K_r,t)=\int\frac{dK_r}{2\pi}e^{i\omega_r t}\sum_p e^{-i(-k_r+K_r)pR}\times\tilde{\psi}(K_r)E_r\left(z=L,t-\frac{pR}{v_r}\right). \tag{17}$$

Upon illumination by $\mathcal{E}_r(z,t)$, the polarization driving the propagation equation [8] for the reconstructed field $\mathcal{E}_c(z,t)$ is given by

$$\mathcal{P}_c(z,t)=\delta n(z,t)\mathcal{E}_r(z,t). \tag{18}$$

The evolution of the reconstructed field $\mathcal{E}_c(z,t)$ will follow that of the original signal field $\mathcal{E}_1(z,t)$ if this polarization [Eq. (18)] can be shown to be proportional to $\mathcal{E}_1(z,t)$; the multiplicative constant in this relationship includes the third-order susceptibility $\chi^{(3)}$ rather than the linear susceptibility $\chi^{(1)}$, since the grating $\delta n(z,t)$ given by Eq. (16) is proportional to the product of two optical fields [8].

We can multiply both sides of Eq. (18) by $\exp(-iK_c z)$, and integrate over z to write Eq. (18) in Fourier-transformed K space. In doing so, we use Eqs. (16) and (17), which describe the grating and the reference pulse, to obtain

$$\tilde{\mathcal{P}}_c(K_c,t)=\frac{1}{F_0}e^{i(\omega_1-\omega_2+\omega_r)t}\sum_{n,m,p}\int\frac{dK_1}{2\pi}\frac{dK_2}{2\pi}\frac{dK_r}{2\pi}\delta\hat{n}(K_1,K_2)\left\{\int dz e^{i(K_1-K_2+K_r-K_c)z}\right\}$$
$$\times\; e^{-i(k_1+K_1)nR}e^{-i(-k_r+K_r)pR}e^{i(-k_2+K_2)mR}\tilde{\psi}(K_1)\tilde{\psi}(K_r)\tilde{\psi}^*(K_2)E_1\left(z=0,t-\frac{nR}{v_1}\right)$$
$$\times E_r\left(z=L,t-\frac{pR}{v_r}\right)E_2^*\left(z=L,t-\frac{mR}{v_2}\right). \tag{19}$$

The phase-matching integral over z can be approximated by $2\pi\delta(K_1-K_c)\delta(K_2-K_r)$, and we can carry out the integrals over K_2 and K_1. We write $K_r\equiv K'$ and $K_c\equiv K$ to generalize the notation. In order to focus the discussion on a holographic reconstruction of the signal pulse, we now assume that $E_r(z=L,t)=E_2(z=L,t')|_{t'=t}$, $k_r=k_2$, $\omega_r=\omega_2$, and $v_r=v_2$, i.e., we use a replica of the reference write-in pulse $\mathcal{E}_2$ (in the original temporal coordinate t') as the reference reconstruction pulse. Then, Eq. (19) becomes

$$\tilde{\mathcal{P}}_c(K_c,t)=e^{i\omega_1 t}\sum_n e^{-i(k_1+K_1)nR}\tilde{\psi}(K_1)E_1\left(z=0,t-\frac{nR}{v_1}\right)\left[\frac{1}{2\pi F_0}\int\frac{dK'}{2\pi}\delta\hat{n}(K,K')\right.$$
$$\left.\times\sum_{m,p}e^{-i(-k_2+K')pR}e^{i(-k_2+K')mR}\tilde{\psi}(K')\tilde{\psi}^*(K')E_2\left(z=L,t-\frac{pR}{v_2}\right)E_2^*\left(z=L,t-\frac{mR}{v_2}\right)\right]. \tag{20}$$

The term in square brackets in Eq. (20) can be written as

$$[\cdots]=\frac{1}{2\pi F_0}\int\frac{dK'}{2\pi}\delta\hat{n}(K,K')G(K',t)G^*(K',t), \tag{21}$$

where

$$G(K',t)=\sum_p e^{-i(-k_2+K')pR}\tilde{\psi}(K')E_2\left(z=L,t-\frac{pR}{v_2}\right). \tag{22}$$

We can multiply $G(K',t)$ by $\exp(i\omega_2 t)$ without changing Eq. (21).

Then, using the relationship established in Eq. (14),

$$\tilde{\mathcal{P}}_c(K,t)=\tilde{\mathcal{E}}_1(K,t)\times h(K,t), \tag{23}$$

where

$$h(K,t)=\frac{1}{2\pi F_0}\int\frac{dK'}{2\pi}\delta\hat{n}(K,K')|\tilde{\mathcal{E}}_2(K',t)|^2. \tag{24}$$

We assume that the photorefractive properties of the CROW characterized by $\delta\hat{n}$ are spectrally nonselective:

$$\delta\hat{n}(K,K')\equiv\delta\hat{n}\,\delta(K-K')\qquad\text{for all }\;K\text{ and }K'. \tag{25}$$

Further, we assume that the (backward-propagating) reference pulses are intense and narrow, i.e., the input free-space reference pulse $\mathcal{E}_2(z=L,t)$ is given by

$$\mathcal{E}_2(z=L,t)=e^{i\omega_2 t}E_0\delta(t-t_0). \tag{26}$$

The corresponding field in the CROW is given by Eq. (14):

$$\tilde{\mathcal{E}}_2(K',t)=E_0e^{i\omega_2 t}\sum_m e^{i(-k_2+K')mR}$$
$$\times\tilde{\psi}(K')\,\delta\left(t-t_0-\frac{mR}{v_2}\right). \tag{27}$$

Next, if the individual resonator modes are highly localized, $\psi(z)=\hat{\psi}\delta(z-z_0)$, we can write

$$\tilde{\mathcal{E}}_2(K',t)\tilde{\mathcal{E}}_2^*(K',t)=|E_0|^2|\hat{\psi}|^2\sum_{m,m'}e^{-i(-k_2+K')(m-m')R}$$
$$\times\left[\delta\left(t-t_0-\frac{mR}{v_2}\right)\delta\left(t-t_0-\frac{m'R}{v_2}\right)\right]. \tag{28}$$

For cumulative power transfer over a finite time interval, as given by integrating the above expression over a region of t comparable to or greater than R/v_2, the term in square brackets in Eq. (28) can be replaced by the Kronecker delta $\delta_{mm'}$, and the result is

$$\tilde{\mathcal{E}}_2(K',t)\tilde{\mathcal{E}}_2^*(K',t)=|E_0|^2|\hat{\psi}|^2M, \tag{29}$$

where M is the number of resonators. Using the normalization relationship Eq. (8), we can simplify Eq. (24),

$$h(K,t)=\frac{1}{(2\pi)^2F_0}\delta\hat{n}|E_0|^2, \tag{30}$$

which is a constant $\equiv\hat{h}$, and therefore,

$$\tilde{\mathcal{P}}_c(K,t)=\hat{h}\tilde{\mathcal{E}}_1(K,t). \quad (31)$$

This shows that the polarization term driving the evolution of the reconstructed pulse is indeed proportional to the input signal pulse, as it would be for the input signal pulse itself. Note that the scaling constant is dependent on the intensity of the reference pulse, but for weak signal pulses, the total field intensity is dominated by the reference and $F_0 \approx |E_0|^2$ so that the two factors cancel each other.

III. DISCUSSION

The earlier sections have described pulse propagation in a CROW, and have shown that photorefractive holography via short and intense reference pulses in a weakly coupled CROW with spectrally nonselective photorefractive properties exactly reconstructs the signal pulse. In this section we point out two features of the CROW that make it particularly suitable for such holographic pulse storage and reconstruction processes.

As discussed by Yariv *et al.* [2], it has been shown that, in a weakly coupled CROW, the dispersion relation for a waveguide mode is approximately

$$\omega_k=\omega_\Delta\left[1+\frac{\Delta\alpha}{2}+\kappa\cos(kR)\right], \quad (32)$$

in terms of the single-resonator mode frequency ω_Δ, a coupling factor κ, and an overlap integral $\Delta\alpha$. (Although the single defect cavity modes are actually doubly degenerate, the two resultant CROW bands have opposite polarities and cannot couple to each other; therefore, the dispersion relation of each band has the same form as the above expression [11].) In an earlier analysis, we have assumed that the central wave vector k_0 corresponds to a linear section of this curve; an analysis using the form of Eq. (32) is presented elsewhere. The k-dependent group velocity (Ref. [8], p. 37) is given by

$$v_g(k)=\frac{d\omega_k}{dk}=-\omega_\Delta R\kappa\sin(kR), \quad (33)$$

which can be made quite small (e.g., $v_g \approx 10^{-3}$ times the group velocity in a medium with spatially uniform dielectric properties and refractive index n_2) for a weakly coupled CROW [11], and for sufficiently narrow band pulses. Because of this reduction in group velocity, an optical pulse propagating in the CROW is compressed by a factor $1/v_g$ relative to its spatial extent in free space. Spatial compression of pulses is also observed in electromagnetically induced transparency [12], where the slow group velocity is the result of the steep slope of the refractive index. In the case of the CROW, the slow group velocity is a consequence of the weak evanescent coupling between the individual high-Q modes that comprise the propagating eigenmode of the waveguide. A CROW designed for a factor of 10^{-3} reduction in group velocity will permit a Gaussian pulse of temporal duration 150 ps to be completely contained in a waveguide of length 100 μm.

Photorefractive holography of single pulses has been difficult because of the low efficiency of the process, and usually, multiple write-in procedures of thousands of repeated pulses are necessary to obtain a sufficiently strong quasi-steady state grating. In a CROW, the highly concentrated optical field can also enhance this aspect of the photorefractive effect. The propagating power flux P in a CROW is proportional to the group velocity of the CROW band [2],

$$P=\frac{1}{8\pi R}v_g E_0^2, \quad (34)$$

and, therefore, we can obtain a higher optical field for a given power flux because of the reduction in group velocity. Consequently, the time constant which determines the photorefractive response time (and which varies linearly with the intensity [8]) is reduced by a factor $v_g \approx 10^{-3}$ compared to the group velocity in a medium with spatially uniform dielectric properties and refractive index n_2. The quasi-steady-state equilibrium is reached with orders of magnitude lower intensities in a photorefractive CROW as compared to a photorefractive bulk medium. As pointed out by Yeh [13], the fundamental limit on the speed of the photorefractive effect depends on the intensity rather than the phenomenological parameters induced by doping or heat treatment.

In summary, we have analyzed optical pulse propagation in a coupled-resonator optical waveguide (CROW), and propose a method for the storage and reconstruction of optical pulses using photorefractive holography in a CROW. The advantages of this method include a substantial reduction in the group velocity, leading to a spatial compression of the signal pulse so that it may be contained in a relatively short waveguide compared to the spatial extent of the pulse in free space. The highly localized field distribution enhances the photorefractive effect, and we have examined in detail the process of the formation of the grating and the reconstruction of the signal pulse by holography. There are many possible applications of such room-temperature, compact, nondestructive, and low-intensity pulse storage mechanisms; two important ones are buffers for optical switches and correlators for optical measurement devices.

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