Bistability and symmetry breaking in distributed coupling of counterpropagating beams into nonlinear waveguides

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By using a modal approach, the distributed coupling of two inhomogeneous, transient waves with opposite angles of incidence into nonlinear waveguides is studied. We show that cross modulation and the coherent interaction via a dynamic grating, nonlinearly induced by the excited counterpropagating leaky waves, provide a longitudinal feedback. This feedback is the necessary prerequisite for bistability to occur. We show that bistability may even appear when one incident beam is weak. Consequently, modulation of the strong signal beam by a weak control pulse can be achieved. Furthermore, we demonstrate that the symmetry of the input signal can be broken beyond a certain intensity and that set-reset flip-flop operation with finite beams is feasible. We demonstrate that the mechanism responsible for this flip-flop operation differs considerably from that acting in the plane-wave case. The optical nonlinearities may be caused by either virtual or real carrier excitation, which models the behavior of semiconductors as well as nonlinear polymers or organic dyes well below or close to a resonance, respectively. For real carrier excitation the interplay of carrier relaxation, diffusion, and the propagation of the excited fields is taken into account.

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I. INTRODUCTION

The theoretical modeling of the optical response of nonlinear planar resonators excited by a pulsed beam has attracted considerable interest [1–10]. The modal theory, formulated by Ulrich [11] for the linear situation and extended the nonlinear case by other authors [5–10] later on proves fairly versatile and provides a clear physical picture of the nature of the optical response [10].

Since the geometry of the resonator enters the modal theory merely by the location and the linewidth of the resonance [6–10], the optical response of Fabry-Pérot resonators with metallic or dielectric (Bragg) mirrors as well as prism and grating couplers may be described by that approach. Adopting this point of view, Fabry-Pérot resonators with dielectric mirrors, operated under oblique incidence as well as prism couplers, represent attenuated total internal reflection (ATR) configurations. The former variant is nothing other than a Bragg waveguide where the guided leaky wave may be excited directly from the air [12].

In contrast to plane-wave models, which predict optical bistability under appropriate conditions and for any angle of incidence, mere switching occurs for oblique incidence provided that the beam width is finite [6–8,10] and the nonlinearity is local. The physical explanation is rather straightforward if one studies the behavior of the leaky waves that are excited near any resonance of the configuration. In the plane-wave case these leaky waves are excited everywhere along the structure and their amplitudes are correlated by definition in both directions, providing an inherent longitudinal feedback, which is necessary for optical bistability to occur.

The situation is different for an excitation with a finite beam where the information about the amplitude of the excited leaky wave flows exclusively in the direction of propagation. Thus the longitudinal feedback disappears provided that the response of the nonlinear material is local.

Keeping in mind the analogy between the different resonant devices, it is not surprising that the physics of switching and bistability identified for Fabry-Pérot resonators under oblique incidence coincides with that of a nonlinear prism coupler extensively discussed by Stegeman *et al.* [13] and Vitrant *et al.* [14]. There, it has been shown that bistability may appear only if either the material response is nonlocal (due to diffusion or thermal effects) or an additional longitudinal feedback is introduced by a weak counterpropagating wave. Particularly, the counterpropagating wave was assumed to be excited by reflection of the excited guided wave at the prism edge [13]. Due to the small reflectivity, this feedback is rather weak. Only when the detuning of the input field from the ATR resonance is extremely large [13] is this feedback sufficient to yield bistability. Since the detuning has to be compensated for by the nonlinear susceptibility changes in the waveguide a high switching power, as well as a sufficiently large saturation level of the respective nonlinear process are required.

The other feedback mechanism discussed was nonlocality of the nonlinear response of the material caused, e.g., by diffusion [13,14]. It may produce optical bistability when the diffusion length is comparable to the decay length of the excited leaky wave, which corresponds to the coupling length of the ATR configuration. This condition is always met for thermal nonlinearities, but thermal effects are considered to be too slow for optical information processing. For carrier-induced nonlinearities, as, e.g., in semiconductors, there is a similar tradeoff because the diffusion length scales with the square root of the relaxation time of the carriers. Hence large diffusion lengths are always accompanied by slow relaxation and consequently a small repetition rate of the device envisaged. However, by increasing the detuning [14] one can compensate for a small diffusion length, provided that the saturation value of the nonlinearity is large enough.

The present paper takes advantage of the idea to introduce longitudinal feedback by a counterpropagating wave. The operation speed of this mechanism is limited only by the group velocity of the excited leaky waves and has practically no inherent delay. Hence it seems to be preferable for fast optical information processing. Nevertheless, the influence of undesired diffusion effects needs to be estimated.

In contrast to [13], we assume that the backward propagating wave is excited by an additional input beam, launched into the guide with the opposite angle of incidence with respect to the primary one. This scheme provides an additional degree of freedom to control the optical response of the system. Similar investigations were carried out first by Haelterman [15,16], but restricted to the plane-wave approach and instantaneous Kerr nonlinearities. As an interesting new effect symmetry breaking was predicted for input waves with equal intensities.

The primary aim of our paper is twofold, namely, to study the various effects for finite beam width and to cover both resonant and nonresonant nonlinearities. First, we investigate whether bistability can be observed for finite beam widths and transient excitation, taking exactly into account the interaction of the two waves via cross-phase modulation and a transient nonlinear grating. The prerequisites for optical bistability to occur, as, e.g., the critical ratio between the amplitudes of both input beams, are identified. Additionally, the influence of carrier diffusion is discussed. Since we are studying the temporal dynamics of the response, the stability of the field structures can be directly identified. Moreover, we show that effects such as, e.g., filamentation, leading to the formation of high transmission domains, play an important role for finite beam widths.

Furthermore, we study how the output characterisitic of one beam can be controlled by the counterpropagating one. This interplay may offer opportunities for a definite response control. Eventually, we investigate symmetry-breaking effects for finite input beams with equal intensities.

It is commonly agreed that direct semiconductor materials are promising candidates for nonlinear optical devices such as nonlinear waveguides [17] as well as planar resonators (see, e.g., [18]). Two advantageous frequency domains have been identified. One can use frequencies of the optical field below half the fundamental band gap where the nonlinearity is caused by virtual carrier excitation. In this frequency domain the nonlinearity is extremely fast (quasi-instantaneous) and one- as well as two-photon absorption processes can be avoided [17,19]. By contrast, for frequencies near the fundamental gap, the nonlinearity is due to real carrier excitation [18,20]. Hence the nonlinearity is some orders of magnitude larger, but also some orders of magnitude slower. Additionally, considerable absorption losses occur. Consequently, one prefers frequencies below half the gap if the guided field is used as the signal and very short response times are required. On the contrary, frequencies near the gap can be employed if the reflected-transmitted signal is exploited and low switching powers are desirable. In order to cover these different situations our investigations are carried out simultaneously for both types of nonlinearities.

Similar considerations hold also for other nonlinear materials in which the nonlinearity is due to electronic excitations, such as, e.g., nonlinear polymers or organic dye molecules [21]. The main difference to the semiconductor case is the negligible role of carrier diffusion. All these materials exhibit saturable nonlinearities with a finite response time when excited at resonance and instantaneous nonlinearities in the nonresonant case [21].

The paper is organized as follows. In Sec. II we present the basic equations of the modal theory by taking into account cross modulation and the interaction via the nonlinearity induced grating. In Sec. III the optical response of our configuration is studied numerically for various excitation conditions.

II. MODAL THEORY FOR DISTRIBUTED COUPLING OF COUNTERPROPAGATING WAVES INTO NONLINEAR WAVEGUIDES

A. Field equations

We consider an ATR configuration consisting of either a prism-loaded waveguide or a Fabry-Pérot cavity with a multilayer cladding under oblique incidence of the exciting beam (Bragg waveguide) (see Fig. 1). In the latter case guiding is provided by reflection at Bragg mirrors rather than by total internal reflection. Since details of the underlying modal approach are discussed elsewhere [5–10], it shall suffice to give a concise derivation of the resulting basic equations which describe the evolution of the counterpropagating leaky waves as well as the nonlinear polarization. For the sake of convenience we restrict ourselves to one-dimensional beams and TE-polarized fields.

As fundamental dynamic variables, describing the evolution of the fields in the system, we use the amplitudes of the leaky waves at the film-substrate interface, referred to as a transmitted field [9]. The reflected field can be easily calculated by a linear relationship from the transmitted field [6,9]. By taking the Fourier transform of the transmitted (E_t) and the incident (E_{in}) field as well as of the nonlinear polarization (P_{NL}) with respect to the coordinate parallel to the interfaces (z) and the time and approximating the linear transmission function near the resonance by a Lorentzian, we may write for one Fourier component of the transmitted field

$$\hat{E}_t(\beta,\omega)=\frac{\tilde{\tau}}{\beta_r^2(\omega)-\beta^2}\left[\hat{E}_{\rm in}(\beta,\omega)-\mu_0\omega^2\int_{\rm NL} g(x)\hat{P}_{\rm NL}(x,\beta,\omega)dx\right], \quad (1)$$

where $\beta_r(\omega)$ is the complex resonance, $g(x)$ the nonresonant part of the Green's function, and $\tilde{\tau}$ the transfer coefficient of the configuration (for details see [9,10]). Note that the field at any point in the configuration is given by

$$\hat{E}(x;\beta,\omega)=f(x)\hat{E}_t(\beta,\omega)\ , \quad (2)$$

with $f(x)$ being the mode profile. Provided that the modulus of the resonance angle is large compared to the angular width of the resonances, we may separate the field, applying to $\beta=\pm\beta_r(\omega)$ resonances, and designate them as the forward and backward propagating fields, respectively. Hence we may restrict ourselves to two distinct domains in Fourier space and split (1) into two equations for the respective directions

$$\hat{E}_t^{f/b}(\beta,\omega)=\frac{\hat{\tau}}{\beta_r(\omega)\mp\beta}\left[\hat{E}_{\rm in}^{f/b}(\beta,\omega)-\mu_0\omega^2\int_{\rm NL} g(x)\hat{P}_{\rm NL}^{f/b}(x,\beta,\omega)dx\right], \quad (3)$$

where $\hat{\tau}=\tilde{\tau}/2\beta_r$. In expanding β_r near the central frequency ω_0 into a Taylor series up to first order and applying the inverse Fourier transform, we end up with two differential equations for the amplitudes of the forward and the backward propagating wave, which are slowly varying in time

$$\left[i\frac{1}{v_g}\frac{\partial}{\partial t}\pm i\frac{\partial}{\partial z}+\beta_0\right]E_t^{f/b}(z,t)+\hat{\tau}\mu_0\omega_0^2\times\int_{\rm NL} g(x)P_{\rm NL}^{f/b}(x,z,t)dx=\hat{\tau}E_{\rm in}^{f/b}(z,t)\ . \quad (4)$$

Here v_g is the group velocity of the leaky wave, $\beta_0=\beta_r(\omega_0)$, and the + sign applies to the forward propagating wave.

For the nonlinear polarization we may write

$$P_{\rm NL}(x,z,t)=\varepsilon_0\chi_{\rm NL}(x,z,t)f(x)[E_t^f(z,t)+E_t^b(z,t)]\ . \quad (5)$$

We separate rapidly oscillating terms from the fields and the polarization as, e.g.,

$$E_t^f(z,t)=\tilde{E}_t^f(z,t)e^{i\bar{\beta}z}\ ,\quad E_t^b(z,t)=\tilde{E}_t^b(z,t)e^{-i\bar{\beta}z}\ , \quad (6)$$

where $\bar{\beta}=k_0n_c\sin|\varphi|$ $(k_0=\omega_0/c)$ denotes the modulus of the mean propagation constant of the incident beams (determined by their angles of incidence φ). Now $\tilde{E}_t^f(z,t)$ and $\tilde{E}_t^b(z,t)$ vary slowly in space (z) and time (t).

For any nonlinear susceptibility that depends on the field intensity we may identify two major contributions.

(i) One is connected with a slowly varying term induced by the total guided intensity (self- and cross modulation). It represents the incoherent interaction of both fields and leads to a nonlinear shift of the resonances.

(ii) A second, coherent, term arises from the interference of forward and backward propagating fields (nonlinearly induced grating) leading to energy exchange between both fields.

In the framework of the modal theory we may separate the nonlinear susceptibility as

$$\chi_{\rm NL}(x,z,t)=\hat{\chi}(x)[\chi_0(z,t)+\chi_1(z,t)e^{2i\bar{\beta}z}+\chi_{-1}(z,t)e^{-2i\bar{\beta}z}]\ , \quad (7)$$

where χ_0 and $\chi_{\pm1}$ represent slowly varying envelopes with respect to z and t. $\chi_{\pm1}$ are the amplitudes of a nonlinearly induced grating that couples both waves. Since we consider intensity-driven nonlinearities, $\hat{\chi}(x)$ is mainly determined by the transverse intensity distribution $|f(x)|^2$ of the leaky waves and, provided that carriers are excited, by diffusion. Obviously, it depends on the kind of the nonlinearity and has to be specified later on. By

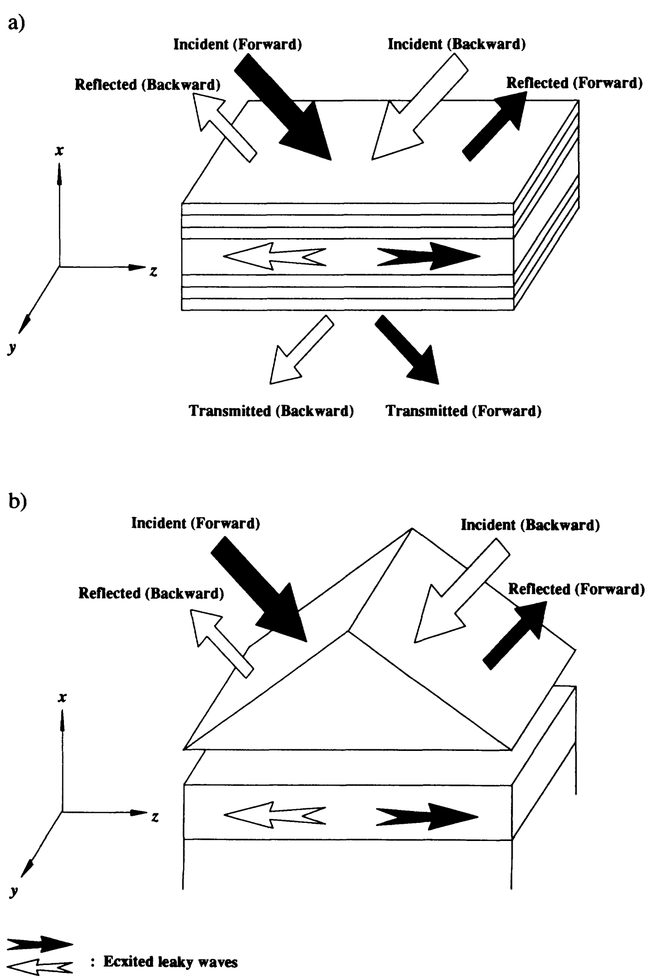

FIG. 1. Geometry of the ATR-like configurations under study: (a) the leaky Bragg waveguide and (b) the prism-loaded waveguide.

using (4)–(7) we are now in the position to introduce the slowly varying nonlinear polarizations in the forward and the backward direction as

$$\tilde{P}^f_{\rm NL}(x,z,t)=\varepsilon_0 f(x)\hat{\chi}(x)[\chi_0(z,t)\tilde{E}^f_t(z,t)+\chi_1(z,t)\tilde{E}^b_t(z,t)]\ , \tag{8a}$$

$$\tilde{P}^b_{\rm NL}(x,z,t)=\varepsilon_0 f(x)\hat{\chi}(x)[\chi_0(x,t)\tilde{E}^b_t(z,t)+\chi_{-1}(z,t)\tilde{E}^f_t(z,t)]\ , \tag{8b}$$

which give, with (4), the coupled field equations

$$\left[\frac{i}{v_g}\frac{\partial}{\partial t}+i\frac{\partial}{\partial z}+(\mathrm{Re}\beta_0-\bar{\beta})+i\,\mathrm{Im}\beta_0\right]\tilde{E}^f_t(z,t)+k_0\tilde{\chi}[\chi_0(z,t)\tilde{E}^f_t(z,t)+\chi_1(z,t)\tilde{E}^b_t(z,t)]=\hat{r}\tilde{E}^f_{\rm in}(z,t)\ , \tag{9a}$$

$$\left[\frac{i}{v_g}\frac{\partial}{\partial t}-i\frac{\partial}{\partial z}+(\mathrm{Re}\beta_0-\bar{\beta})+i\,\mathrm{Im}\beta_0\right]\tilde{E}^b_t(z,t)+k_0\tilde{\chi}[\chi_0(z,t)\tilde{E}^b_t(z,t)+\chi_{-1}(z,t)\tilde{E}^f_t(z,t)]=\hat{r}\tilde{E}^b_{\rm in}(z,t)\ , \tag{9b}$$

where the effective susceptibility is defined as $\tilde{\chi}=\hat{r}k_0\int_{\rm NL}g(x)f(x)\hat{\chi}(x)dx$.

B. Nonlinear susceptibility

1. Resonant excitation of carriers

The simplest model that contains the most prominent effects, such as saturation of the nonlinearity, a finite carrier relaxation time, and diffusion of carriers, is a two-level system with diffusion added [20]:

$$\left[T_1\frac{\partial}{\partial t}-L_d^2\frac{\partial^2}{\partial x^2}-L_d^2\frac{\partial^2}{\partial z^2}+1\right]\chi_{\rm NL}(x,z,t)+a[\chi_{\rm NL}(x,z,t)-\bar{\chi}]|E(x,z,t)|^2=0\ . \tag{10}$$

Here T_1 is the relaxation time of the carriers and L_d their diffusion length. $\bar{\chi}$ stands for the saturation value of the nonlinear susceptibility and a reflects the strength of carrier excitation by an optical field. This model is exact for the case of excitations in dye molecules. But it can also be regarded as a good approximation for semiconductors [20] or polymers excited slightly below their gap energy, if one assumes a linear relationship between the number of excited carriers and the induced susceptibility changes.

The actual diffusion length of the material defines two different regimes. If the diffusion length is much less than the thickness of the guide, diffusion is negligible and $\hat{\chi}(x)$ is determined by the intensity of the guided wave as

$$\hat{\chi}(x)\approx\bar{\chi}|f(x)|^2\ , \tag{11a}$$

where we have neglected the minor effect of saturation. For strong diffusion, occurring, e.g., in semiconductors ($L_d\approx 10$ μm), $\hat{\chi}(x)$ is smeared out by diffusion, which gives

$$\hat{\chi}(x)\approx\bar{\chi}\frac{1}{l}\int_{\rm NL}|f(x')|^2dx'=\mathrm{const}\ , \tag{11b}$$

where l is the thickness of the nonlinear material.

Note that (10) implies a scaling of the susceptibility with its saturation value $\bar{\chi}$. We included $\bar{\chi}$ into $\hat{\chi}(x)$ to fix the upper bound of the normalized slowly varying envelopes to unity.

Inserting (6), (7), and (11) into the equation of motion of the susceptibility (10), we obtain, for the incoherent contribution,

$$T_1\frac{\partial\chi_0(z,t)}{\partial t}-L_d^2\frac{\partial^2\chi_0(x,t)}{\partial z^2}+\chi_0(z,t)+a[\chi_0(z,t)-1)[|\tilde{E}^f_t(z,t)|^2+|\tilde{E}^b_t(z,t)|^2]+a\chi_1(z,t)\tilde{E}^{f*}_t(z,t)\tilde{E}^b_t(z,t)+a\chi_{-1}(z,t)\tilde{E}^f_t(z,t)\tilde{E}^{b*}_t(z,t)=0 \tag{12a}$$

and for coherent contribution (nonlinear grating)

$$T_1\frac{\partial\chi_1(z,t)}{\partial t}-L_d^2\frac{\partial^2\chi_1(z,t)}{\partial z^2}-4i\bar{\beta}L_d^2\frac{\partial\chi_1(z,t)}{\partial z}+[1+(2\bar{\beta}L_d)^2]\chi_1(z,t)+a[\chi_0(z,t)-1]\tilde{E}^f_t(z,t)\tilde{E}^{b*}_t(z,t)+a\chi_1(z,t)[|\tilde{E}^f_t(z,t)|^2+|\tilde{E}^b_t(z,t)|^2]=0\ . \tag{12b}$$

Since the absorptive contributions to nonlinearity are already included in $\bar{\chi}$, χ_0 represents a real quantity and $\chi_{-1}=\chi_1^*$ holds.

The major effect caused by diffusion consists in bleaching the induced grating $\chi_{\pm1}$. Phenomenologically it can be described by a reduced effective recombination time $T_{\rm eff}=T_1/[1+(2\bar{\beta}L_d)^2]$ of the carriers that induce the grating. It is determined by the ratio of the diffusion length and the period of the induced grating. In what follows we assume that χ_0,χ_1 vary slowly and neglect all derivatives with respect to z (slowly varying envelope approximation). This assumption holds if the diffusion length of the carriers is small compared to the decay length of the excited leaky waves and is well justified for semiconductor materials such as $\mathrm{Al}_x\mathrm{Ga}_{1-x}\mathrm{As}$ [22].

2. Nonresonant excitation

Provided that the mean frequency of the incident pulse is far from any material resonance, the susceptibility may be described by a simple Kerr model

$$\chi_{\rm NL}(x,z,t)=\alpha|E(x,z,t)|^2 \tag{13}$$

where α is the Kerr constant. Here we have to assume that $\hat{\chi}(x)\approx\alpha|f(x)|^2/|\alpha|$ and obtain

$$\chi_0(z,t)=|\alpha|[|\tilde{E}^f_t(z,t)|^2+|\tilde{E}^b_t(z,t)|^2]\ , \tag{14a}$$

$$\chi_1(z,t)=|\alpha|\tilde{E}^f_t(z,t)\tilde{E}^{b*}_t(z,t)\ , \tag{14b}$$

where χ_0 is real and $\chi_{-1}=\chi_1^*$. In conclusion, we have derived the complete set of equations governing the evolution of the optical fields (9) and the excitation in the nonlinear material (12) or (14).

C. Normalized equations

It is convenient to employ an appropriate normalization in order to minimize the number of free parameters and to cover different physical situations by using the decay length of the leaky wave

$$L=\frac{1}{\mathrm{Im}\beta_0}\ , \tag{15a}$$

the lifetime of the leaky wave

$$T_c=\frac{1}{\mathrm{Im}\beta_0}\left.\frac{\partial\beta_0}{\partial\omega}\right|_{\omega_0}=\frac{L}{v_g}\ , \tag{15b}$$

and the effective susceptibility

$$\chi_{\mathrm{eff}}=\frac{\hat{r}k_0^2}{\mathrm{Im}\beta_0}\int_{\mathrm{NL}}g(x)f(x)\hat{\chi}(x)dx\ . \tag{15c}$$

Now the dimensionless quantities $Z=z/L$, $T=t/T_c$, $u_{f/b}=\kappa\tilde{E}_t^{f/b}$, and $u_{f/b}^{\mathrm{in}}=(\hat{r}\kappa/\mathrm{Im}\beta_0)\tilde{E}_{\mathrm{in}}^{f/b}$ are introduced [5–10] where $\kappa=\sqrt{a}$ for resonant and $\kappa=\sqrt{|\alpha|}$ for nonresonant nonlinearities hold, respectively. We obtain, for the fields,

$$\left[i\frac{\partial}{\partial T}+i\frac{\partial}{\partial Z}+\Delta+i+\chi_{\mathrm{eff}}\chi_0(Z,T)\right]u_f(Z,T)+\chi_{\mathrm{eff}}\chi_1(Z,T)u_b(Z,T)=u_f^{\mathrm{in}}(Z,T)\ , \tag{16a}$$

$$\left[i\frac{\partial}{\partial T}-i\frac{\partial}{\partial Z}+\Delta+i+\chi_{\mathrm{eff}}\chi_0(Z,T)\right]u_b(Z,T)+\chi_{\mathrm{eff}}\chi_1^*(Z,T)u_f(Z,T)=u_b^{\mathrm{in}}(Z,T)\ , \tag{16b}$$

where $\Delta=(\mathrm{Re}\beta_0-\bar{\beta})/\mathrm{Im}\beta_0$ is the detuning of the incident field from the ATR resonance characterized by the complex propagation constant β_0 of the leaky wave.

For resonant nonlinear susceptibilities [see (12)], normalization results in

$$T_R\frac{\partial\chi_0(Z,T)}{\partial T}+\chi_0(Z,T)+[\chi_0(Z,T)-1]\times[|u_f(Z,T)|^2+|u_b(Z,T)|^2]+\chi_1(Z,T)u_f^*(Z,T)u_b(Z,T)+\chi_1^*(Z,T)u_f(Z,T)u_b^*(Z,T)=0\ , \tag{17a}$$

$$T_R\frac{\partial\chi_1(Z,T)}{\partial T}+\frac{1}{\delta}\chi_1(Z,T)+[\chi_0(Z,T)-1]u_f(Z,T)u_b^*(Z,T)+\chi_1(Z,T)[|u_f(Z,T)|^2+|u_b(Z,T)|^2]=0\ , \tag{17b}$$

where $T_R=T_1/T_c$ is the normalized relaxation time. Note that the lifetime of the nonlinearly induced grating is diminished due to diffusion by a factor of $\delta=T_{\mathrm{eff}}/T_1<1$. For the Kerr nonlinearity [see (14)] we obtain

$$\chi_0(Z,T)=|u_f(Z,T)|^2+|u_b(Z,T)|^2\ , \tag{18a}$$

$$\chi_1(Z,T)=u_f(Z,T)u_b^*(Z,T)\ . \tag{18b}$$

III. NUMERICAL RESULTS

In what follows, we solve the basic system of equations numerically for different excitation conditions. We use the explicit scheme [23], which is based on an integration of the partial differential equations along their characteristic lines. Effects such as nonlinearly induced switching delay, optical bistability, symmetry breaking, modulation and flip-flop operation are discussed. Both input fields are assumed to be Gaussian in space and time as

$$u_{f/b}^{\mathrm{in}}(Z,T)=A_{f/b}\exp\left[-\frac{Z^2}{W^2}\right]\exp\left[-\frac{T^2}{T_0^2}\right]\ . \tag{19}$$

To compare the respective response characteristic we have performed the calculations for resonant (carrier-induced) as well as nonresonant (Kerr) nonlinearities. Since carrier-induced effects are discussed less in the literature, most figures are related to that case.

A. Nonlinearly induced switching delay for one input beam

First we investigate the optical response without a backward propagating input beam. As reported in [6–8,10,13,14] we found that switching rather than bistability occurs as far as no nonlocality (e.g., by diffusion) is introduced. Bistability sets in when the diffusion length of the carriers exceeds a certain threshold value. For our configuration that threshold was found to be approximately 40% of the decay length L of the leaky wave. Even in semiconductors it is difficult to meet this requirement in an ATR-like configuration, where L is about 100 μm and the ambipolar diffusion length for carriers, e.g., in GaAs, is below 10 μm [22]. However, an interesting retardation effect arises, provided that carriers are excited and a critical bias amplitude is applied. The dynamics of the switching process is depicted in Fig. 2. The system is biased with an extremely long pulse and two control pulses are launched into the waveguides. The bias amplitude equals the amplitude where the system switches under stationary conditions. The most striking feature is a strong retardation effect in the excitation of the nonlinear leaky wave when the input amplitude approaches its critical value. It is evident from Fig. 2 that the relaxation time of the nonlinearity ($T_R=2$) as well as the decay time of the leaky wave ($T_c=1$) are approximately two orders of magnitude less than the buildup or decay times of the nonlinear leaky wave. Hence this delay may be regarded as an inherently nonlinear effect. It introduces a considerable amount of hysteresis into the switching characteristics for short pulse excitation which may even be confused with optical bistability [9]. On the other hand, this nonlinear delay may have a stabilizing effect on optical bistability if a small amount of additional longitudinal feedback is provided.

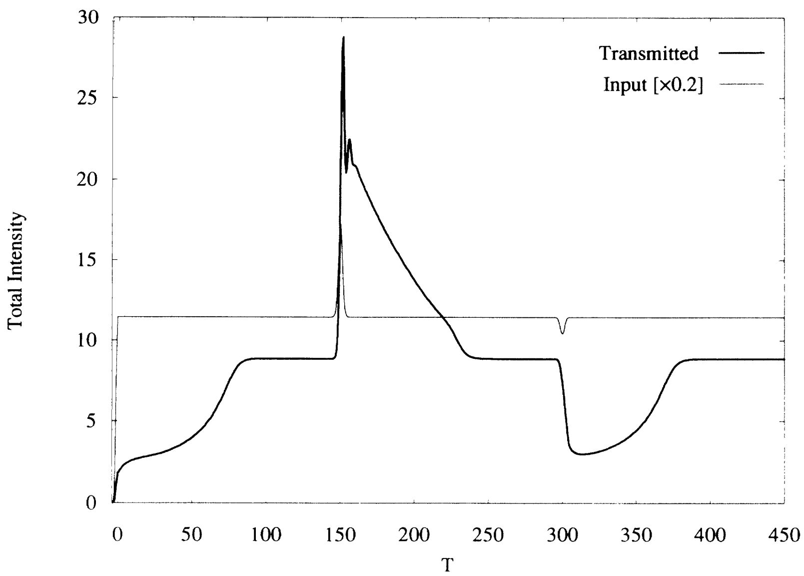

FIG. 2. Temporal evolution of the intensity of the excited leaky wave when a pulse sequence is applied as the input and a resonant nonlinearity is assumed. Diffusion is neglected. The bias amplitude is $A=2.1373$, which corresponds to the respective critical input intensity. The parameters are $\Delta=6$, $\chi_{\rm eff}=-10$, $T=2$, $W=10$, and $T_0=2$.

B. Optical bistability induced by a counterpropagating input beam

In order to introduce a longitudinal feedback into the system, which is necessary to obtain optical bistability, we have launched an additional weak stationary input beam with an opposite angle of incidence (backward beam). The pulse sequence is applied only to the forward input. The results for a Kerr nonlinearity are depicted in Fig. 3. The bistable character of the response for both counterpropagating leaky waves can be clearly recognized. Note that the intensity of the backward input wave may be as small as about 10% of that of the forward input (compare Fig. 5). We found numerically that this represents some lower threshold for the detuning considered. For resonant nonlinearities this threshold is slightly higher (approximately 15%).

By increasing the detuning, the threshold for the backward beam can be reduced further. For instance, we found a threshold as low as 1% for a detuning of $\Delta=15$ for a resonant nonlinearity with $\chi_{\rm eff}=-20$. Note, however, that the maximum detuning achievable in a given configuration is always limited by the saturation value of the nonlinearity.

The major difference between the two types of nonlinearity consists in the appearance of fluctuations in the high transmission state for the Kerr nonlinearity. The physical explanation of these oscillations can be found by looking to the field structures. In the Kerr case the guided field tends to form domains of different mean direction of propagation which propagate across the beam in the direction of the stronger wave (see Fig. 4). The formation of such structures is suppressed if the nonlinearity contains a memory effect which smears out the domain boundaries.

Diffusion plays only a secondary role for the situation

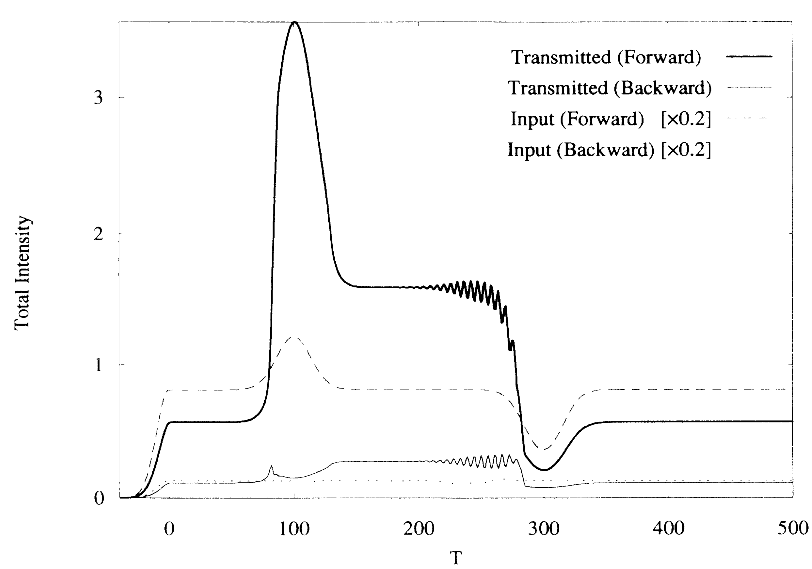

FIG. 3. Temporal evolution of the intensity of the excited leaky waves when a pulse sequence is used as the input in the forward direction and additional feedback is provided by a weak stationary beam in the backward direction for a Kerr nonlinearity. The parameters are $\Delta=-3$, $\chi_{\rm eff}=10$, $W=10$, $T_0=20$, $A_+=0.57$, and $A_-=0.23$.

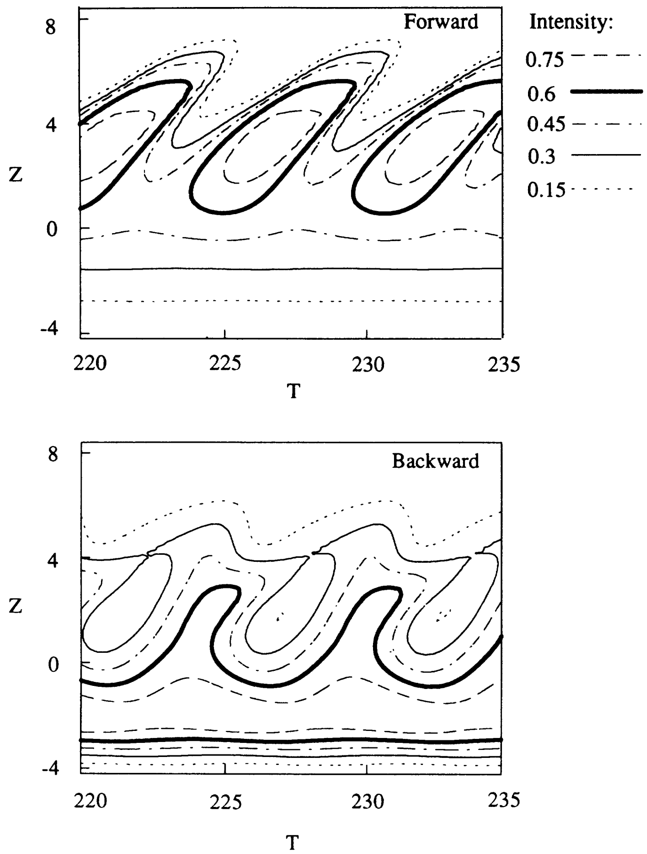

FIG. 4. Temporal evolution of the spatial intensity distributions of the forward and backward waves, respectively, in the case of Kerr nonlinearity (see Fig. 3). The time domain was selected to demonstrate the intensity fluctuations of the fields.

studied here. We found numerically that strong diffusion, i.e., a complete lack of any induced grating, increases the relative amount of power needed in the backward beam from approximately 10% to 20% of the forward beam at a detuning of $\Delta=6$. Note that the induced grating likewise disappears if the two input beams are incoherent. This is of particular interest for optical signal processing since independent sources may be used for signal and pump beams. The results obtained for strong diffusion can therefore be straightforwardly generalized to this case. Hence optical bistability, evoked by incoherent counterpropagating beams, may be anticipated.

The reverse situation, namely, that a strong constant field in the forward direction is modulated by a weak signal sequence in the backward direction, is of particular interest for signal processing because it might offer the opportunity to trigger a strong signal beam by a weak pulsed control beam. This effect can be clearly identified from Fig. 5 for a resonant nonlinearity. Note that the evolution of the excited guided waves is very similar to that obtained in the case of a constant input in the backward direction (compare with Fig. 3).

C. Symmetry breaking

Now we investigate the distributed coupling of two pulsed beams with opposite angles of incidence, but equal amplitudes, into the waveguide. Recently, it has been shown for the plane-wave case that the amplitudes of the counterpropagating leaky waves can become different beyond a certain input intensity (symmetry breaking) [15,16]. Furthermore, it turned out that beyond this bifurcation point, only the asymmetric solutions are stable. The aim of our investigation is to find out whether similar effects can be identified if the beam width is finite.

In the stationary, plane-wave case the system of equations (15)–(17) reduces to algebraic relations for the respective leaky wave intensities $I_{f/b}=|u_{f/b}|^2$ as functions of the input intensities $I^{\rm in}_{f/b}=|u^{\rm in}_{f/b}|^2$. In the Kerr case [see (15) and (16)] we obtain

$$I^{\rm in}_{f/b}=I_{f/b}|\Delta+i+\chi_{\rm eff}(I_{f/b}+2I_{b/f})|^2 , \tag{20}$$

whereas a nonlinearity caused by real carrier excitation

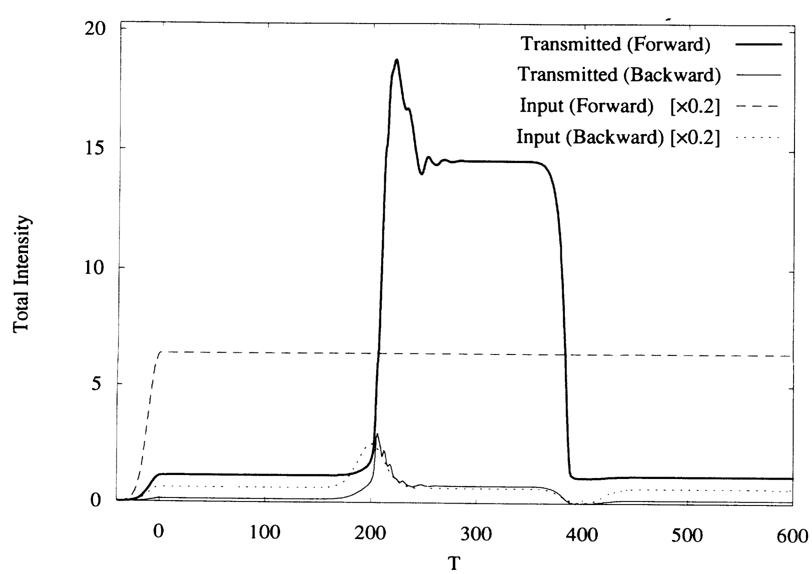

FIG. 5. Modulation of a strong stationary beam in the forward direction by a weak pulse sequence applied in the backward direction for a carrier-induced nonlinearity. The parameters are $\Delta=6$, $\chi_{\rm eff}=-10$, $T=2$, $W=10$, $T_0=20$, $A_+=1.6$, and $A_-=0.5$.

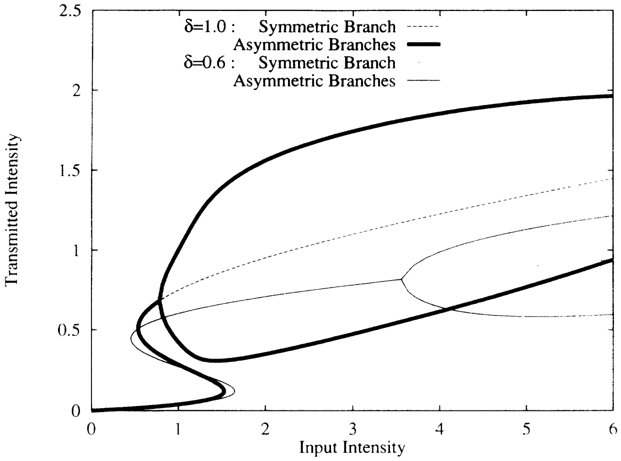

FIG. 6. Intensity of the guided wave versus the input intensity both in the forward direction for a stationary plane-wave excitation ($T_0 \rightarrow \infty$ and $W \rightarrow \infty$). The input intensity is the same in both directions. The parameters are $\Delta=6$ and $\chi_{\text{eff}}=-10$. The results are shown for negligible diffusion ($\delta=1$) and for the case that diffusion reduces the induced grating by 40% ($\delta=0.6$).

(15) and (16) leads to

$$I_{f/b}^{\text{in}}=I_{f/b}\left|\Delta+i+\chi_{\text{eff}}\frac{I_{f/b}+(1+\delta)I_{b/f}+\delta(I_f^2+I_b^2)}{1+(1+\delta)(I_f+I_b)+\delta(I_f^2+I_b^2)}\right|^2, \tag{21}$$

where $\delta<1$ denotes the reduction of the induced grating strength by diffusion. Note that for strong diffusion, i.e., $\delta=0$, the squared term on the right-hand side becomes completely equal for the forward and the backward direction, resulting in a complete symmetry between both fields. Hence we may conclude that symmetry breaking will require at least a weak induced grating and, consequently, a diffusion length of the excited carriers that does not exceed the grating period considerably. On the other hand, it is possible to increase the period length of the induced grating by a suitable reduction of the angles of incidence and thus restore the possibility for symmetry breaking. Additionally, we may conclude that the coherence of the two incoming beams with respect to each other is essential to obtain symmetry breaking. Solutions of (21) are shown in Fig. 6. Symmetry breaking occurs in the high transmission state beyond a certain input intensity (bifurcation point). Qualitatively similar curves are obtained for the Kerr case. As reported in [15,16], for the Kerr nonlinearity the asymmetric solutions proved to be stable, whereas the symmetric ones become unstable beyond the bifurcation point.

To discuss the influence of diffusion in more detail we have added the results for $\delta=0.6$. The bifurcation point shifts to a higher intensity and the separation between both asymmetric branches is reduced. If the diffusion strength increases further, bifurcation disappears completely. For the present detuning ($\Delta=6$) bifurcation occurs only provided that $\delta>0.5$. Note that this limit reduces if the detuning is increased.

Under pulsed beam excitation, the situation gets more involved. We find symmetry breaking with respect to the total intensities of the leaky waves for both types of nonlinearity. The respective response curves to a symmetric input pulse sequence are depicted in Fig. 7 for resonant nonlinearities. When comparing with the plane-wave results, the appearance of a stable, apparently symmetric branch (with equal total output intensities in the forward and the backward direction) seems to be surprising. But, when looking at the respective spatial field distributions (see Fig. 8), it is evident that the situation encountered for finite beams differs considerably from the symmetric plane-wave state, which is characterized by equal, guided wave intensities in both directions at any point along the distributed coupler. Both leaky waves exhibit an asymmetric intensity distribution and are mirror images of each other. Hence the state with equal total intensities proves to consist of two domains, in each of which a different asymmetric branch of the plane-wave output dominates. Since the longitudinal profiles are mirror im-

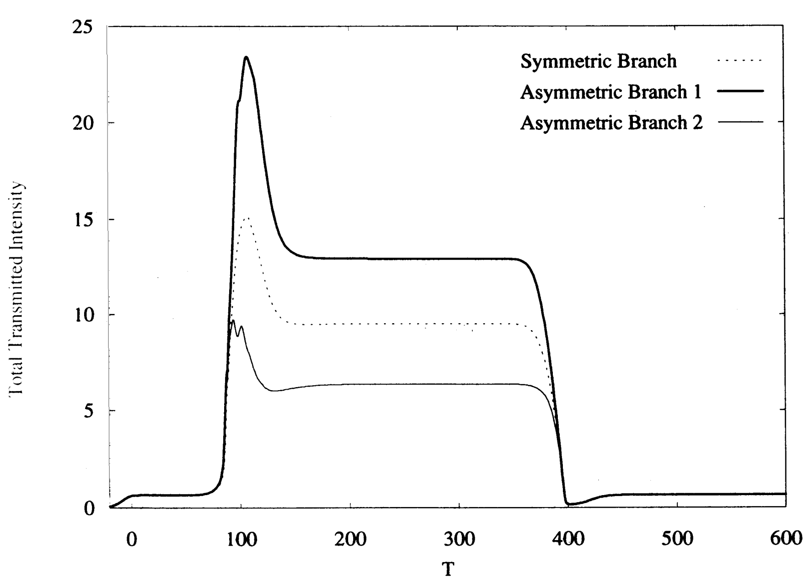

FIG. 7. Temporal evolution of the intensity of the forward propagating guided wave for equal input pulse sequences in both directions for large detuning and a resonant nonlinearity. The parameters are $\Delta=6$, $\chi_{\text{eff}}=-10$, $W=10$, $T=2$, $T_0=20$, and $A_{+\text{bias}}=A_{-\text{bias}}=1.14$.

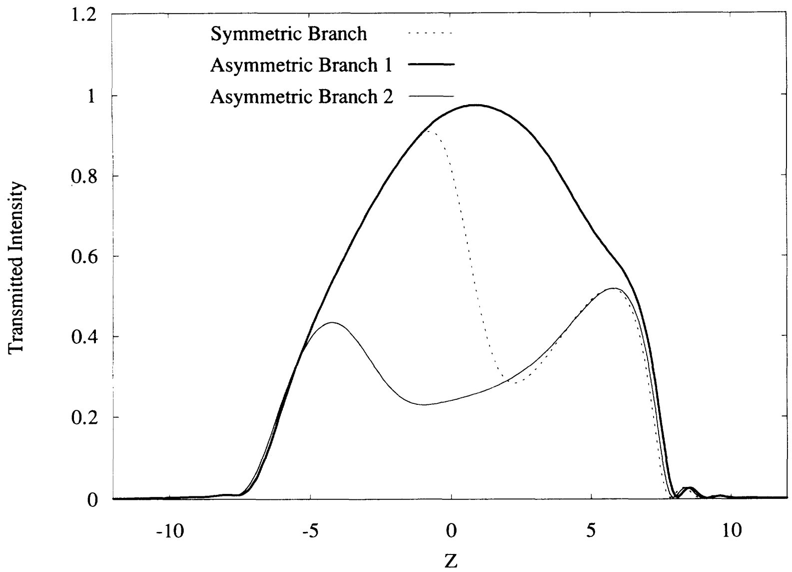

FIG. 8. Longitudinal intensity distribution of the guided fields in the forward direction corresponding to the different branches in Fig. 7.

ages, the total transmitted power is equal in both channels. We suppose a cooperative effect between the two beams to be the reason for the appearance of a two-domain structure: Each beam creates a region of optimum excitation for the other one. As a result, this kind of structure proves to be extremely stable.

On the "asymmetric" branches one domain dominates, resulting in different total intensities in the forward and backward directions. We may conclude that the stability of all states involved in Fig. 7 is consistent with the results of the plane-wave model [15,16] because they are related to the asymmetric branches occurring there. Note that this symmetry breaking can also be considered as another type of multistability, which is related to the formation of domains in the longitudinal pattern of the excited guided waves with different ratios of transmission in the forward and backward directions.

If we reduce the detuning (Fig. 9) without changing the width of the exciting beam, it turns out that the states containing only a single domain (i.e., asymmetric states) are no longer stable, though their decay is extremely slow compared to the linear-response times of the system. The ratio of the decay time with respect to the linear-response may exceed even two orders of magnitude.

On the other hand, it is possible to restore multistability also in the case of small detuning by reducing the width of the exciting beams. We demonstrate this effect for a beam width of $W=2$ (see Fig. 9). Obviously, narrower beams favor a one-domain structure of the excited fields. (But, nevertheless, the two-domain structure always remains stable.) This result differs in some respect from those found for usual optical bistability (see Sec. III B), where wide beams are preferred. On the other hand, it may present a considerable advantage for possible applications because it allows one to reduce the size of a device as well as the total power needed for excitation.

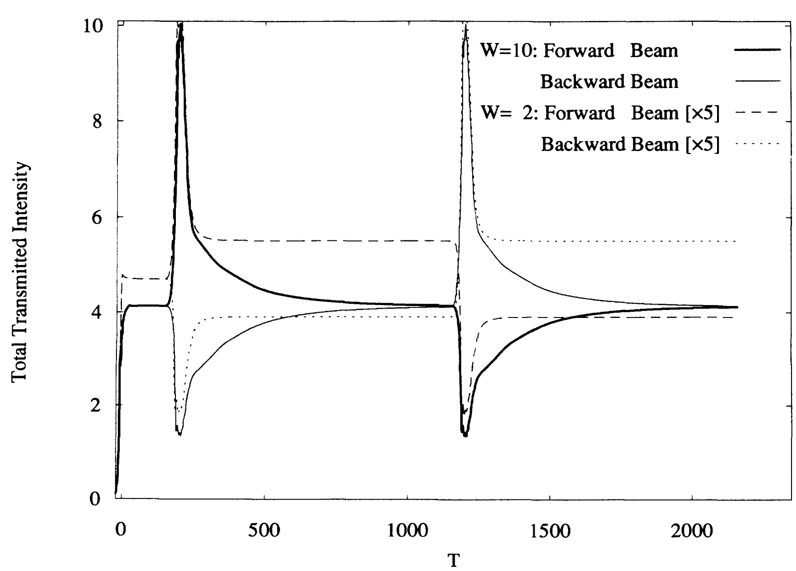

FIG. 9. Temporal evolution of the intensity of the forward propagating guided wave for equal input pulse sequences in both directions and small detuning. Carrier-induced nonlinearity was assumed. The beam widths considered were $W=10$ (full lines) and $W=2$ (broken lines), respectively. Other parameters are $\Delta=3$, $\chi_{\mathrm{eff}}=-10$, $T=2$, $T_0=20$, and $A_{+\mathrm{bias}}=A_{-\mathrm{bias}}=0.8$.

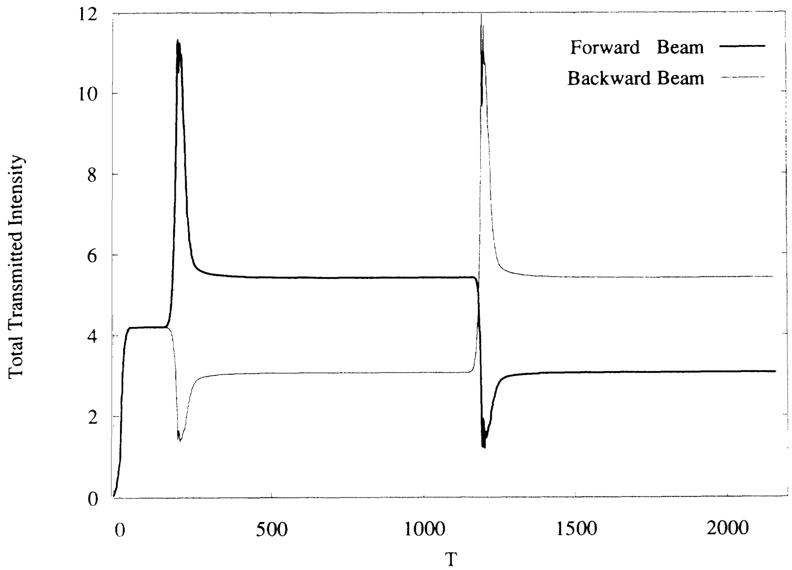

FIG. 10. Demonstration of the set-reset flip-flop operation with stationary signal input beams of equal intensity. Pulsed control beams into opposite directions are added subsequently and a resonant nonlinearity was used. The parameters are $\Delta=4$, $\chi_{\rm eff}=-10$, $T=2$, $W=10$, $T_0=20$, $A_{+\rm bias}=A_{-\rm bias}=0.75$, and $A_{\rm pulse}=0.4$.

D. Set-reset flip-flop operation

Our previous investigations have shown that one can switch the output intensity in one direction, say forward, between two well-defined states by adding or subtracting a weak control pulse in the backward direction, provided that an appropriate detuning has been chosen. The question arises whether it is possible to implement that switching operation by only adding and not subtracting control pulses. This would fit much better to experimental conditions because it turns out difficult to create a "negative" control pulse. This idea was first proposed by Haelterman [15,16], but studied for a plane-wave input only. His suggestion relies on the effect of symmetry breaking in the high-transmission state of the symmetrically excited coupler. One launches two stationary waves with equal amplitudes above the bifurcation point into the waveguide. Then one adds a weak control pulse in one direction, which drives the system to the asymmetric branch (high transmission in one direction, low transmission in the other one). Since the asymmetric branches are stable, the system sustains this state. After adding a weak control pulse in the other direction, the system switches to the second asymmetric branch with the reversed transmission characteristic (set-reset flip-flop operation). It is of particular interest whether this behavior survives if finite beams are used rather than plane waves. The results of our numerical simulations are depicted in Fig. 10 for a resonant nonlinearity, but appear to be similar for a Kerr nonlinearity. They show that flip-flop operation can also be achieved with finite beams.

Since the state with equal total power in both directions is stable, a certain minimum of power is needed in the switching pulses. The conditions to be met are the same as for symmetry breaking, discussed in Sec. III C. Especially the need for an induced grating and for coherence of both input beams has to be stressed.

IV. CONCLUSIONS

We have shown that optical bistability may occur in ATR-like configurations excited by finite beams if longitudinal feedback is provided by a counterpropagating input beam. It turned out that it suffices that this beam has an intensity one order of magnitude less than the primary one. Furthermore, we found that a strong signal beam may be modulated by a weak control pulse sequence. If one launches two beams with equal intensities into the waveguide, symmetry breaking may occur in the output channels beyond a certain bifurcation intensity. It has been shown that both the asymmetric and the symmetric branches are stable if we regard symmetry with respect to the total transmitted intensities. The longitudinal beam profiles, related to the symmetric branch characterized by equal total transmitted intensities in both directions, are asymmetric with respect to the local intensity. Hence their stability does not contrast to the plane-wave results. Eventually, symmetry breaking was exploited to achieve flip-flop operation between both transmission directions. All investigations were carried out for nonlinearities caused by virtual (Kerr nonlinearity) or real carrier excitation in direct semiconductors, where it turned out that only quantitative differences of the response characteristic exist.

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