논문 96-5-5-11

# Embodiment of all-optical switching phenomena on a GaAs waveguide

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## Abstract

Based on the transmission of coupled gap solitons in nonlinear periodic media, we present an all-optical switching scheme which has a novel architecture and principle. The proposed switch with an extremely small switching element can be realized on a semiconductor waveguide. We here investigate the switching performance with a GaAs waveguide in order to give criteria for the experimental realization of the all-optical switching phenomena. We also suggest a variation of an index-matching scheme to solve the technical problem such as the input-energy coupling into a periodic waveguide.


## I. Introduction

Recent work[1] has shown that a nonlinear periodic dielectric medium ( NPDM ) with $\chi^{(3)}$ nonlinearity can transmit an input laser beam with a certain intensity although its frequency is located within the stop gap. Such optical fields transmitted in the NPDM is referred to as gap solitons because their envelopes have hyperbolic secant shapes with solitary-propagating properties. The existence of stationary solitons in a NPDM was first investigated by Chen and Mills[1],[2] by use of a computer simulation in layer-by-layer calculation. Later, de Sterke et al.[3],[4] and Christodoulides et al.[5] showed that gap solitons can propagate in a NPDM with a group velocity much slower than the usual group velocity determined by the medium's refractive-index dispersion. In particular, de Sterke et al.[6],[7] derived a nonlinear Schrodinger equation ( NLSE ) for the gap solitons by using an envelope-function approach. On the associated dipersion curve obtained form Kronig-Penny model, a stationary gap soliton will be positioned at the Brillouin zone edge where the group velocity vanishes. However, a moving gap soliton has somehow higher peak intensity, leading to larger shift of the stop gap than a stationary gap soliton does. Consequently, the associated dispersion relation will be governed by a nonzero-slope point out of the Brillouin zone, thereby allowing the pulse to attain wave-momentum, giving rise to a traveling wave. Meanwhile, Christodoulides et al.[5],[8] and Wabnitz[8] have used a coupled-mode formalism to analyze the electrodynamics in a NPDM. The electric fields in a periodic structure are decomposed into a forward-propagating mode and a backward-propagating mode. The periodic index, however, is modified by the intensity-dependent index due to the third-order nonlinearity. As a result, a set of coupled-wave equations will describe the propagation of the two contradirectional

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modes, which are coupled by a cross-phase modulation as well as by a distributed feedback mechanism. It was found[5],[8] by means of the massive Thirring model in a quantum field theory that they have solitary-wave solutions with square-root hyperbolic-secant envelopes.

We are now interested in the propagation of two orthogonally-polarized solitary modes in a NPDM with $\chi^{(3)}$ nonlinearity. In the previous paper[9],[10], we have proposed the all-optical switch based on the transmission of coupled gap solitons in a NPDM. Specifically, we showed that two orthogonally polarized pulses can copropagate as a coupled gap soliton through a NPDM, while each pulse alone will be strongly reflected because its amplitude is less than that needed to propagate a gap soliton in a single polarization. Based on the results, we presented a new all-optical switching scheme. In the present paper, we investigate the switching performance with a GaAs waveguide as the NPDM in order to give criteria for the experimental realization of the all-optical switching phenomena. We also consider a technical problem in a specific periodic waveguide. To be specific, it would be difficult to couple the input energy into the periodic waveguide because of the reflections at the interfaces between the Bragg grating region and the surrounding region. As a solution for such a problem, we suggest a variation of an index-matching scheme, which was originally proposed for out-of-gap operation in a linear periodic structure by Haus[11].

The body of the paper is comprised of the following sections. In Sec. II, we study the spatio-temporal evolutions of two orthogonally polarized modes in a NPDM with $\chi^{(3)}$nonlinearity. Employing a coupled-mode approach, we find the four coupled differential equations describing the propagation of the four orthogonally polarized modes, and then we analyze the linear and the nonlinear interactions between the four modes. By means of an envelope-function approach, we derive the coupled NLSEs governing the nonlinear dynamics of the electric fields in a NPDM, and then we show that the coupled NLSEs have coupled-gap solitary-wave solutions in a simple case. In Sec. III, we present an all-optical switching scheme based on the transmission of coupled gap solitons. In particular, we investigate the feasibility of the proposed switch when the switching element is embodied by a corrugated GaAs waveguide. The technical problem such as the input-energy coupling into a periodic waveguide is also treated. Finally we discuss the results of the present paper in Sec. IV.

## II. Basic equations for optical waves with two orthogonally polarized modes in a nonlinear periodic dielectric medium

We here study the propagation of two pulses with orthogonal polarizations in a NPDM with $\chi^{(3)}$ nonlinearity. Before proceeding to the main derivations, we take several assumptions given as follow. First, the plane wave approximation and the slowly varying envelope approximation underpin the entire theoretical development. The transverse profiles of the electric fields is assumed to remain constant during its propagation in a lossless waveguide. Consequently, the general expression for total electric field is

$$\begin{aligned}\mathbf{E}\ (x,y,z,t) = \frac{1}{2}\{[E_{fx}(z,t)e^{i(kz-\omega t)} \\ + E_{bx}(z,t)e^{-i(kz+\omega t)}]F_1(x,y)\,\hat{x} \\ + [E_{fy}(z,t)e^{i(lz-\omega t)} + E_{by}(z,t)e^{-(lz+\omega t)}] \\ F_2(x,y)\,\hat{y} + C.C\}\end{aligned} \tag{1}$$

where $E_{f_x}$ is the envelope of $x$-polarized, forward-propagating field, $E_{b_y}$ is the envelope of $y$-polarized, backward-propagating field, and so on. The influence of transverse variation of field will be approximately included in the effective area

$A_{eff}=[\int\int_{-\infty}^{\infty}|F(x,y)|^2dxdy]^2/\int\int_{-\infty}^{\infty}|F(x,y)|^4dxdy$ and then in the nonlinear coefficients $\gamma=3\pi\omega_0\chi^{(3)}/4\pi cA_{eff}$. Second, the off-diagonal elements of linear dielectric tensor will be neglected, and $x$-and $y$-axis in waveguide coincide the ordinary and the extraordinary axis of birefringent medium, respectively. Thus there is no linear coupling between copropagating orthogonally-polarized components. Third, it is required that the periodic perturbation of refractive index be small. Suppose that $n_x(z)$ and $n_y(z)$ are the linear periodic refractive index along $x$-axis and $y$-axis respectively, and they have the same modulation depths. Then they can be expressed as

$$n_x(z) = n_{0x}+n_1\cos(\frac{2\pi}{\Lambda}z)$$
$$n_y(z) = n_{0y}+n_1\cos(\frac{2\pi}{\Lambda}z), \quad (2)$$

where the modulation depth $n_1$ of refractive index is assumed to be so small that

$$n_x^2 \approx n_{0x}^2+2n_{0x}n_1\cos(\frac{2\pi}{\Lambda}z)$$

and $n_y^2 \approx n_{0y}^2+2n_{0y}n_1\cos(\frac{2\pi}{\Lambda}z)$.

Using a coupled mode approach, we find a set of coupled equations given as follows:

$$\frac{\partial E_{fx}}{\partial z}+\frac{1}{v_{gx}}\frac{\partial E_{fx}}{\partial t}$$
$$=i\chi_xE_{bx}e^{-2i\Delta kz}+i\gamma(|E_{fx}|^2+2|E_{bx}|^2+\frac{2}{3}|E_{fy}|^2+\frac{2}{3}|E_{by}|^2)E_{fx}+i\frac{2}{3}\gamma E_{fy}E_{by}E_{bx}^*$$

$$\frac{\partial E_{bx}}{\partial z}-\frac{1}{v_{gx}}\frac{\partial E_{bx}}{\partial t}$$
$$=-i\chi_xE_{fx}e^{2i\Delta kz}+i\gamma(|E_{bx}|^2+2|E_{fx}|^2+\frac{2}{3}|E_{fy}|^2+\frac{2}{3}|E_{by}|^2)E_{bx}+i\frac{2}{3}\gamma E_{fy}E_{by}E_{fx}^*$$

$$\frac{\partial E_{fy}}{\partial z}+\frac{1}{v_{gy}}\frac{\partial E_{fy}}{\partial t}$$
$$=i\chi_yE_{by}e^{-2i\Delta lz}+i\gamma(|E_{fy}|^2+2|E_{by}|^2+\frac{2}{3}|E_{fx}|^2+\frac{2}{3}|E_{bx}|^2)E_{fy}+i\frac{2}{3}\gamma E_{fx}E_{bx}E_{by}^*$$

$$\frac{\partial E_{by}}{\partial z}-\frac{1}{v_{gy}}\frac{\partial E_{by}}{\partial t}$$
$$=-i\chi_yE_{fy}e^{2i\Delta lz}+i\gamma(|E_{by}|^2+2|E_{fy}|^2+\frac{2}{3}|E_{fx}|^2+\frac{2}{3}|E_{bx}|^2)E_{by}+i\frac{2}{3}\gamma E_{fx}E_{bx}E_{fy}^* \quad (3)$$

where

$$v_{gx}=\frac{c}{n_{0x}},\ v_{gy}=\frac{c}{n_{0y}},\ \chi_x=\chi_y=\frac{\pi n_1}{\lambda_0},$$
$$\Delta k=k-\frac{\pi}{\Lambda},\ \text{and}\ \Delta l=l-\frac{\pi}{\Lambda}.$$

The first terms in the right side of Eq.(3) indicate that there exist the linear couplings between the two parallel polarized components due to the feedback mechanism in the periodic structure. All four components, on the other hand, are nonlinearly coupled with each other through the combined actions of nonlinearity and birefringence, giving rise to nonlinear phase shifts, degenerate four-wave mixings (DFWM), and energy fluctuations. For example, the second terms in Eq.(3) are responsible for self-phase modulation (SPM) while the third, the fourth, and the fifth terms are all responsible for cross-phase modulation (XPM) of which the coefficients 2 and 2/3 correspond to the interaction between the parallel polarized components and to the interaction between the orthogonally polarized fields, respectively. The last terms in Eq.(3) account for the phase-matched DFWM. Fig. 1 shows how each of four components interact nonlinearly each other in a NPDM. Note that all energy fluctuations are neglected in Eq.(3).

In order to study the dynamic behavior of the coupled gap solitons, we use an envelope-function approach to derive the coupled NLSEs in a NPDM. Following de Sterke et al.[6],[7] the coupled NLSEs are derived by expressing the electric field $E(z,t)$ in terms of different length scales,

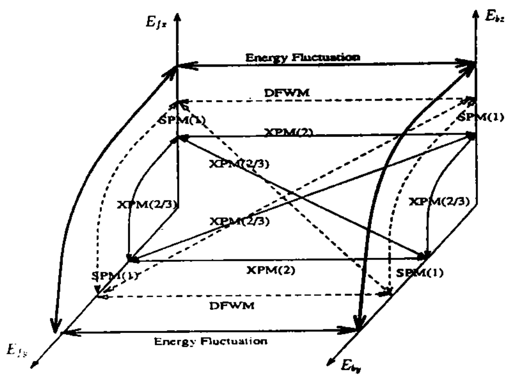


Fig. 1. Illustrative diagram of nonlinear interactions between the four modes in a nonlinear periodic medium.

$z_n = p^n z$ and time scales, $t_n = p^n t$:

$$\mathbf{E}(z,t) = p(e_{1x}\hat{x} + e_{1y}\hat{y}) + p^2(e_{2x}\hat{x} + e_{2y}\hat{y}) + \cdots .$$

Let $a_x$ and $a_y$ be the slowly varying envelopes of two pulses with orthogonal polarizations. Substituting $\mathbf{E}(z,t)$ into the Maxwell equations and then calculating order by order for $p^n$, one can obtain the coupled nonlinear differential equations by combining each results for $n=1$, 2, and 3. The coupled nonlinear differential equations governing temporal evolutions of two envelopes are given as follows:

$$\begin{aligned} & i\frac{\partial a_x}{\partial t} + \frac{1}{2}\omega_{mx}''\frac{\partial^2 a_x}{\partial z^2} + (\alpha_{mxx}|a_x|^2 \\ & + \alpha_{mxy}|a_y|^2)a_x + \alpha_{mk}a_y^2 a_x^* e^{2i(\omega_{mx}-\omega_{my})t} = 0 \\ & i\frac{\partial a_y}{\partial t} + \frac{1}{2}\omega_{my}''\frac{\partial^2 a_y}{\partial z^2} + (\alpha_{myy}|a_y|^2 \\ & + \alpha_{myx}|a_x|^2)a_y + \alpha_{mk}a_x^2 a_y^* e^{2i(\omega_{my}-\omega_{mx})t} = 0, \end{aligned} \tag{4}$$

where

$$\alpha_{mxx} = 6\pi\omega_{mx} L\int_0^L \chi^{(3)}(z_0)|\varphi_{mx}(z_0)|^4 dz_0,$$

$$\alpha_{mxy} = 4\pi\omega_{mx} L\int_0^L \chi^{(3)}(z_0)|\varphi_{mx}(z_0)|^2|\varphi_{my}(z_0)|^2 dz_0 ,$$

$$\alpha_{myx} = 4\pi\omega_{my} L\int_0^L \chi^{(3)}(z_0)|\varphi_{mx}(z_0)|^2|\varphi_{my}(z_0)|^2 dz_0 ,$$

$$\alpha_{mk} = [2\pi(\omega_{mx}-2\omega_{my})^2/\omega_{mx}] \; L\int_0^L \chi^{(3)}(z_0)\varphi_{mx}^2(z_0)\varphi_{my}^{*2}(z_0)dz_0 ,$$

$\omega_{mx}''$ is the curvature of the photonic band, and $\varphi_{mx}$ and $\varphi_{my}$ are the fast varying Bloch functions along $x$- and $y$-axis, respectively. The cross-coupling coefficients $\chi^{(3)}_{abcd}$ are related to the self-coupling coefficient $\chi^{(3)}_{aaaa}$ by $\chi^{(3)}_{xxyy} = \chi^{(3)}_{xyxy} = \chi^{(3)}_{xyyx} = \frac{1}{3}\chi^{(3)}_{xxxx}$. In the present paper we deal mainly with the case in which the last terms (called the analogous four-wave-mixing term) in Eq. (4) can be neglected. It is valid if $x$ and $y$-polarized fields are in an anisotropic medium. For such cases, isolated $x$- and $y$-polarized solitons will propagate with different group velocities. The coupled $x$- and $y$-polarized solitons, however, will exhibit dragging phenomena. For the both cases, Eq.(4) will be reduced to the coupled NLSEs. As discussed below, in the case where the terms of concern is not negligible, simple solutions can be found if $\omega_{mx} = \omega_{my}$. Suppose that the medium and pulse characteristics are identical for TE and TM mode with a slightly different frequency, then $\varphi_{mx} \simeq \varphi_{my}$, $\omega_m \equiv \omega_{mx} \simeq \omega_{my}$, $\omega_{mx}'' \simeq \omega_{my}''$, and thus $\alpha_m \equiv \alpha_{mxx} \simeq \alpha_{myy} \simeq \frac{3}{2}\alpha_{mxy}$. In this case the envelope functions of gap solitons can be further simplified, giving

$$a = a_x = a_y = \sqrt{3/5}\,\mathrm{A}\ \mathrm{sech}(\mathrm{B}\,z)e^{-i\delta t} ,$$

where $\mathrm{A} = \sqrt{-2\delta/\alpha_m}$ and $\mathrm{B} = \sqrt{-2\delta/\omega_m''}$.
In the general case of moving solitons, the solutions give

$$a = a_x = a_y = \sqrt{3/5}\,\mathrm{A}e^{i\mathrm{B}_2 z}e^{-i(\delta+\Delta)t}\,\mathrm{sech}\,\mathrm{B}_1(z - v_g t) , \tag{5}$$

where

$$\mathrm{A} = \sqrt{-2\delta/\alpha_m},\quad \mathrm{B}_1 = \sqrt{-2\delta/\omega_m''},$$
$$\mathrm{B}_2 = \sqrt{2\Delta/\omega_m''},\quad v_g = \sqrt{2\Delta\omega_m''}$$

and the center frequency of soliton is $\omega_c = \omega_m + \delta + \Delta$. The factor $\delta+\Delta$ corresponds to the frequency detuning of the pulse center frequency from the edge frequency of the stop band. We note that our solution can be straight-forwardly extended to the case where the analogous four-wave-mixing (AFWM) term is not negligible provided $\omega_{mx} = \omega_{my}$. In this case there is also a coupled soliton solution where $a_x$ is proportional to $a_y$. In fact, with $a_x$ proportional to $a_y$, the AFWM term would simply give rise to an effective change of coefficients of the bracketed terms in Eq.(4). For example, the case where $a_x = a_y$ discussed above, would have a solution of

$a = a_x = a_y = (1/\sqrt{2})\,A\ \mathrm{sech}(Bz)e^{-i\delta t}$instead.

## III. All-optical switching phenomena in GaAs waveguides

We now present a new all-optical switching scheme based on the previous results. The scheme and principle are different from the intensity-dependent switch studied in the literature[12]. In our proposed switching scheme, we makes use of the fact from Eq.(5) that two orthogonal $\sqrt{\frac{3}{5}}$A sech pulses can copropagate as coupled gap solitons due to cross-coupling $\chi^{(3)}$ nonlinearity although one $\sqrt{\frac{3}{5}}$A sech pulse will be strongly reflected. A single $\sqrt{\frac{3}{5}}$Asech pulse will not propagate through the medium because in the formalism here, the condition required to propagate a gap-soliton pulse in a single polarization is given by a pulse with amplitude A, i.e. a Asech pulse. Thus a $\sqrt{\frac{3}{5}}$Asech pulse with an amplitude smaller than A will be strongly reflected. This proposed light-by-light switch is shown in Fig.2. The main part of the device consists simply of a very small switching element —GaAs waveguide, and two polarization beam splitters (PBS's). In the pulse generation part, a $\sqrt{\frac{6}{5}}$Asech pulse is divided by a PBS into two $\sqrt{\frac{3}{5}}$Asech pulses which are orthogonally polarized with respect to each other. The two synchronized pulses—the signal S and control C are then combined at another PBS, and coupled into the GaAs waveguide. S will be transmitted through the waveguide only with the presence of C due to the characteristic of coupled gap solitons. S then is separated from the coupled state by another PBS at the waveguide output. Compared with the other optical switch, the proposed switch has some potential advantages. For example, the control pulse in the proposed switch will switch directly the signal pulse, which means it can be exploited as a true *logic gate*. Moreover, it is very stable for the external conditions such as sound and heat waves, etc. Remind that in a Mach- Zehnder interferometer the accurate "$\pi$-phase shift" is required for a good switching performance. Furthermore, the proposed switch with an extremely small switching element can be realized on a semiconductor waveguide, and thus it can be easily integrated into optical circuits.

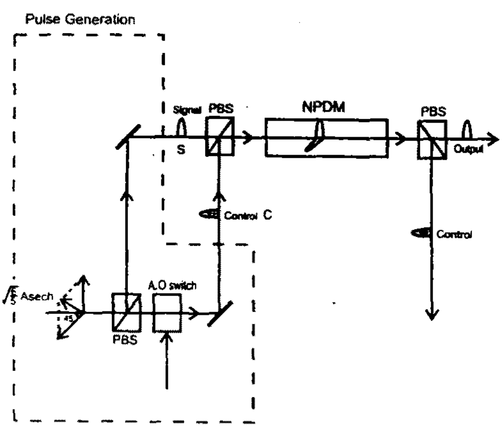


Fig. 2. Schematic of light-by-light switching based on the transmission of coupled gap solitons. A.O, acousto-optical.

Several kinds of pulse and structure parameters will control a switching operation in the proposed scheme. A key relation that connects the pulse parameters and the structure parameters is $\omega_c = \omega_m + \delta + \Delta$. First of all, we need $\omega_c$ (center frequency of pulse) to be located around $\omega_m$ (band-edge frequency) inside the stop gap. Since $\omega_o$ (midgap frequency) and $\Delta\omega$ (bandwidth of stop gap) are determined by $\Lambda$ (spatial period of structure) and $\Delta n$ (index-modulation depth), respectively, one can adjust $\omega_m = \omega_o + \Delta\omega/2$ to a tuning range of the source laser by an appropriate design of $\Lambda$ and $\Delta n$. Let us define a total detuning by $D \equiv |\delta+\Delta| = |\omega_m - \omega_c|$, which indicates a separation of $\omega_c$ from the band-edge. D will be, of course, controlled by tuning of $\omega_c$ in a tunable laser. As D decreases with the fixed peak intensity, $\tau_p$ (temporal width) should be shortened and the gap solitons will propagate in NPDM with a higher group velocity. For good switching performance, we require that $\omega_c$ be far enough from the edge so that the entire pulse spectrum is within the stop band. Let us call this requirement the detuning condition. From the above relations, we find the detuning condition given as follows:

$$
\begin{aligned}
D &= |\delta+\Delta| \geq 2\pi(pulse-spectrum)/2 \\
&= 2\pi\sqrt{\delta(D+\delta)} \\
&\Rightarrow D \geq -0.976\delta.
\end{aligned}
\qquad (6)
$$

As will be seen later, such a detuning condition will restrict significantly the allowed range of parameters. Meanwhile, $\sqrt{\frac{3}{5}}A$ (peak amplitude) influences $\tau_p$ (temporal width) and also $\sigma_p$ (spatial width). First, $\sigma_p$ is inversely proportional to $A$ (; $\sigma_p = \sqrt{\omega_m''/\alpha_m}/A$ from the relations of the soliton parameters given in the previous section). Basically, $L$ (the length of NPDM) should be designed such that it exceeds at least $\sigma_p$. In addition, for a good "off-condition" in the switch, we require enough length of NPDM for a sole $\sqrt{\frac{3}{5}}A$ sech pulse to decay out sufficiently at the output end of NPDM. The waveguide length needed for a coverage of the spatial width, however, will provide a good off-state. Second, the dependence of $\tau_p$ on $I(=|A|^2)$ described by $\tau_p = 1/2\sqrt{\frac{\alpha_m}{2}I(\frac{\alpha_m}{2}I-D)}$ is very strong especially around the stationary soliton condition. Note that the spectrum of the gap soliton will be restricted by $\Delta\omega$ and furthermore the detuning condition. Peculiarly, temporal width as well as the spatial width will contract with increasing $I$ (peak intensity) because of a dependence of $v_g$ (group velocity) on $I$. The $v_g$ is governed by the relation $v_g = \sqrt{\omega''_m(\alpha_m I - 2D)}$ and will mainly influence a switching latency. The $v_g$ will be usually far below the nondispersive group velocity where the soliton characteristics vanish. Nevertheless, the increase of the switching latency due to the slow group velocity is not crucial in an extremely short waveguide used in the proposed scheme.

We have done some numerical calculations to look at the feasibility of the proposed switch. A good candidate for the NPDM is a GaAs waveguide, which was demonstrated by Ho et. al. to have a strong $\chi^{(3)}$ nonlinearity at 1.55-1.65 $\mu m$ [13] From their experimental data, we take the nonlinear refractive index of GaAs at 1.5 $\mu m$ to be $n^{(2)} = 3.6\times10^{-14} cm^2/W$.

Consider an asymmetric corrugated GaAs waveguide with a channel index of $n_f = 3.26291$ and a substrate index of $n_s = 3.14391$ as shown in Fig.3. According to computer calculation, $\Delta h = 0.11\mu m$ (the depth of corrugations) will provide a modulation depth of $\Delta n/\bar{n} = 0.1\%$ in the above waveguide with the channel height of

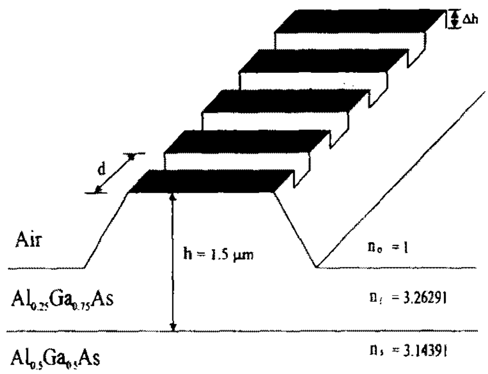


Fig. 3. Asymmetric ridge corrugated GaAs waveguide for the switching element.

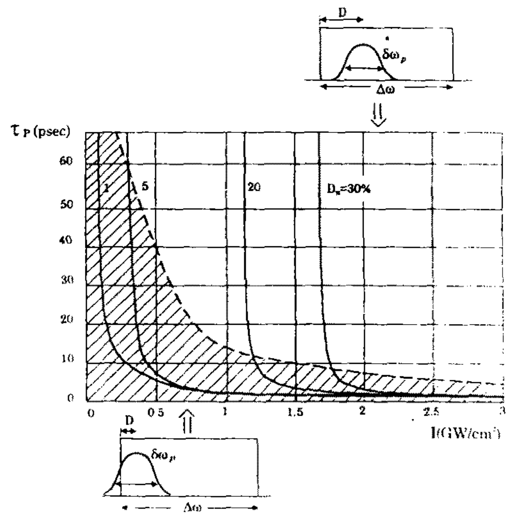


Fig. 4. Dependence of a temporal width on the intensity for each fixed normalized detuning at the waveguide of $\Delta n/\bar{n}=0.1\%$; the dashed line indicates a limitation of pulse bandwidth in conjunction with the center frequency and intensity as illustrated in the miniature figures.

$h=1.5\mu m$. For such a case, the average index $\bar{n}$ becomes 3.23632 so that a period of $d=0.23\mu m$ gives a Bragg wavelength of $\lambda_0=1.5\mu m$. The stop gap bandwidth $\Delta\omega$ is determined by the Bragg structure's refractive-index modulation depth ($\Delta n$). For $\Delta n/\bar{n}=0.1\%$, we have $\Delta\omega=1.25664\times10^{12}/\sec$. In this case, the pulse width is restricted to be larger than 5*psec* if the entire pulse spectrum is to be within the stop-gap bandwidth. Using the above parameters, we investigate the pulse width ($\tau_p$) of the coupled gap soliton as a function of its peak intensity ($I$) for the waveguide of $\Delta n/\bar{n}=0.1\%$. The $I$ is defined as the pulse intensity inside the Bragg structure. The dependence of $\tau_p$ on $I$ would be different for different frequency detuning from the edge of the stop band. Let us define a normalized frequency detuning factor $D_N$ by $D_N\equiv(|\delta+\Delta|)/\Delta\omega=(|\omega_c-\omega_m|)/\Delta\omega$. The numerical results for the dependence of $\tau_p$ on $I$ is shown in Fig.4 where a family of curves are generated with different $D_N$ values.

The region where the detuning condition is satisfied is the unshaded region in the upper right-hand corner of the dashed line. As discussed the above, the detuning condition is satisfied when the pulse spectrum $\delta\omega_p$ is within the stop band $\Delta\omega$ as illustrated in the insert at the upper right corner of Fig.4. We can observe a drastic variation of the temporal width $\tau_p$ around the condition of the stationary coupled-gap solitons in contrast to a slow variation of $\tau_p$ in the other region. In addition to the dependence of pulse width with pulse intensity and pulse frequency detuning, it turns out that the group velocity $v_g$ of the pulse is also a strong function of the pulse parameters. The dependence of $v_g$ on $D_N$ for various pulse widths is plotted as a family of curves in Fig.5. We assume that the medium is nondispersive with the usual group velocity given by $c/\bar{n}$. The value of $v_g$ is normalized by $c/\bar{n}$. Again, the unshaded region in the figure indicates the region where the detuning condition is satisfied.

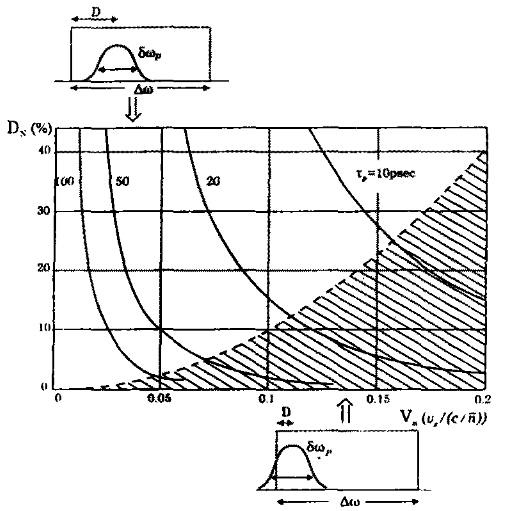

Fig. 5.Dependence of a normalized velocity $(v_g/\frac{c}{\bar{n}})$ on the normalized detuning for each temporal width at the waveguide of $\Delta n/\bar{n}=0.1\%$; the dashed line indicates a limitation of pulse bandwidth in conjunction with the center frequency and intensity as illustrated in the miniature figures.

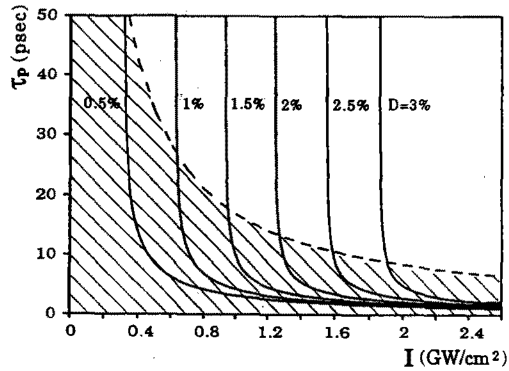

Fig. 6. Dependence of a temporal width on the intensity for each fixed normalized detuning at the waveguide of $\Delta n/\bar{n}=1\%$.

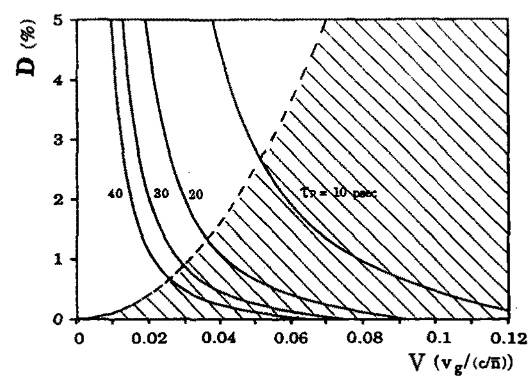

Fig. 7. Dependence of a normalized velocity $(v_g/\frac{c}{\bar{n}})$ on the normalized detuning for each temporal width at the waveguide of $\Delta n/\bar{n}=1\%$.

As an example, we see from Fig.4 that if we have a laser with 20 *psec* pulses, then a coupled gap soliton of 20 *psec* width can be generated with an intensity of larger than 1.2 $GW/cm^2$ and a normalized frequency detuning of larger 20% from the edge of the stop band. From Fig.5, we then see that the group velocity of the coupled gap soliton can range from 0.088 to 0.072 of $c/\bar{n}$ by adjusting the normalized frequency detuning from 20% to 30%. We also investigate the characteristics for another waveguide of $\Delta n/\bar{n}=1\%$ as shown in Fig.6 and Fig.7. As a modulation depth $(\Delta n)$ is increased, the allowed ranges for $\tau_p$ and $\omega_c$ become substantially enlarged with the stop gap. However, there is a limitation for increasing a modulation depth since it is technically difficult to make a deep corrugation on a GaAs waveguide with a finite channel thickness.

An important factor governing the switch is the switching contrast ratio defined as $Q=(Power_{(off-state)}/Power_{(on-state)})$. When the control pulse is off, the single $\sqrt{\frac{3}{5}}A$ sech pulse will decay exponentially in NPDM with a decay length given approximately by $1/B_1$. Using this, we can obtain an estimate for the value of $Q$. For example, if the medium length is 1mm, then $Q$ can be as small as $Q \leq 10^{-7}$ for $I=0.8\sim1.5GW/cm^2$

in the waveguide with a modulation depth of $\Delta n/\bar{n}=1\%$. This 1mm length corresponds to 4348 periods of waveguide corrugations in the above waveguide. This is more than enough periods for the formation of gap solitons. In addition to the above parameters, there are inflexible parameters such as $\chi^{(3)}$ and $\omega''_m$. In particular, high $\chi^{(3)}$ material is preferred as NPDM since in a higher nonlinear material, smaller intensity of laser beam can give rise to the same nonlinear effect. Conclusively, what kind of pulse will be used as coupled gap solitons depends on the laser intensity and the type of waveguide material.

Much higher input power than $I$ would be normally required for the generation of gap soliton because of the reflections at the interfaces between the Bragg grating region and the surrounding region. However, this power loss due to reflections can be greatly reduced using the index-matching methods suggested recently by de Sterke[14] and Haus[11]. In particular, Haus showed that a launched pulse can penetrate into a distributed-feedback (DFB) structure without significant power loss provided that appropriate matching sections are symmetrically placed before and behind the DFB structure. Such matching sections can be constructed by using the Smith chart, which have been devised to calculate a reflection coefficient in the transmission line with a distributed impedance. Even if the matching scheme is based on a linear coupled-mode theory, one can make an ingenious link between our result and the matching method. On account of nonlinear shift of the dispersion curve, a moving gap soliton would operate as if its center frequency is within the allowed band. There are two kinds of mechanisms that can make a nonlinear shift of the photonic band. First, the stop gap will shrink if a nonlinear refractive-index change due to a positive $\chi^{(3)}$ nonlinearity modifies those of only alternative layers (or grating) in the periodic structure. In this case, the size of the stop gap will change due to a change in index-modulation depth while the mid-gap frequency (or the Bragg frequency) is fixed. One, however, can generally expect a nonlinear shift of the stop gap based on another mechanism—the identical nonlinear index change in the both layers (or gratings). In this case, the stop gap will shift without a change in the gap size. Consider only the latter case in the present paper. We now transform Haus' scheme into an index-matching scheme in a nonlinear periodic structure. Let $\delta_{CM}=\frac{\bar{n}}{c}(\omega_c-\omega_0)$, which is the separation of the center frequency of the input frequency from the mid-gap frequency. After the nonlinear shift of the associated dispersion curve down by $\delta$, $\delta_{CM}$ will become $\delta_{non-CM}=\frac{\bar{n}}{c}(\omega_c-\omega_0+\delta)$. Accordingly, an eigen-mode propagation constant $\beta$ should be redefined as $\beta_{non}=\sqrt{\delta^2_{non-CM}-x^2}$ where $x=\pi\Delta n/\lambda_0$ is a coupling constant. Note, however, that in a strict sense, $\beta$ is not equal to $k$ in the previous section. The $k$ in coupled-mode approach indicates an eigen-propagation constant while $\beta$ is involved in the propagation of an whole envelope. An envelope of a gap soliton is made up of an infinite set of partial waves, which bring about another partial waves in a opposite-direction mode due to reflection at the each grating. Substituting the above $\delta_{non-CM}$ and $\beta_{non}$ into the Haus' results, one can find the condition for the length of the matching section,

$$l=\frac{1}{\beta_{non}}\sin^{-1}\{[\beta_{non}(\delta_{non-CM}-\beta_{non})]/[x\sqrt{x^2-(\delta_{non-CM}-\beta_{non})^2}]\} \quad (7)$$

and the condition for the phase shift in the gap section, $\phi_0=\frac{\pi}{2}-\tan^{-1}[\frac{\delta_{non-CM}}{\beta_{non}}\tan(\beta_{non}l)]$. (8)

Fig.8 shows the index-matching scheme where the matching sections are symmetrically placed

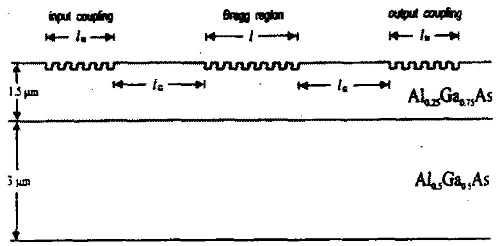


Fig. 8. Index-matching scheme for input-energy coupling into Bragg structure.

before and behind the Bragg structure. If we choose a soliton of $D_N$=3% and $I=1.89 GW/cm^2$ as an example, we know that $\tau_p=5 psec$ from Fig.6 ($\Delta n/\bar{n}=1\%$). From the above relations, we can also easily find that $\kappa=$ $6.87\times10^4$ $1/m$, $\delta_{non-CM}=6.88\times10^4$ $1/m$, and $\beta_{non}=3.71\times10^3$ $1/m$. From Eq.(7) and Eq.(8), we know that the length of the matching section $l_M$ should be $43.3\mu m$ and the phase shift of 0.32 *Rad* is required in the gap section. The gap length of $l_G=0.205$ *mm* will provide the above phase shift for the peak intensity. Note that the nonlinear phase shift due to a peak intensity is given by $\phi_0=2\pi I l n^{(2)}/\lambda$. For a given temporal width, $\Delta\delta/\kappa\simeq0.1$, which guarantees a reflectivity of less than 0.02 over the entire pulse spectrum. As an another example, we choose coupled-gap solitons of $\tau_p=60 psec$ and $D_N=5\%$ for the waveguide with a modulation depth of $\Delta n/\bar{n}=0.1\%$ from Fig.4. The condition corresponds to $I=0.287$ $GW/cm^2$, yielding $\delta=6.302\times10^{10}$ 1/sec and thus $\delta_{non-CM}=6.78\times10^3$ 1/sec. For such a case, we find $\beta_{non}=162.2$ $1/m$ with $\kappa=6.778\times10^3$ $1/m$. Therefore, for an input-energy coupling of coupled-gap solitons, the length of matching sections $l_M$ should be $0.668mm$, and the length of gaps $l_G$ should be $3mm$, providing a nonlinear phase shift of $0.216Rad$.

## Ⅳ. Conclusion and discussion

In this paper we have provided the basic theory governing the operation of the all-optical switch. Specifically, we have studied the propagation of two pulses with two orthogonal polarizations in a NPDM with $\chi^{(3)}$ nonlinearity. First, we investigated the dynamics of the nonlinear pulses in a microscopic viewpoint. Using a coupledmode approach, we see how each of four components in the orthogonally polarized and counterpropagating modes interact nonlinearly each other. Secondly, we found the coupled NLSE's governing the temporal and spatial evolutions of the two orthogonally polarized modes, and then their solitary-wave solutions, referred to as coupled gap solitons. In that procedure, we have used an envelope-function approach, which can provide the overview regarding the behavior of the slowly-varying envelopes in each modes. In fact, the coupled-gap solitary-wave solutions are in good agreement with the nonlinear shifts of parabolic dispersion curves. One can explain the wave-momentum $\Delta k$ and the slow group velocity $v_g=\sqrt{2\Delta\omega''_m}$ in accordance with the nonlinear shift of dispersion relation. Such an important concept underlies the whole analysis throughout this paper.

In conclusion, we have proposed a compact all-optical switch and also investigated its feasibility. Its implementation is simple because the main part is consisted simply of two PBSs and a switching element—GaAs waveguide, which is compact in size. Moreover the proposed device can potentially exhibit good switching performances, including good switching speed with reasonable pulse width, and good on-off intensity contrast ratio. From a practical viewpoint, however, one should solve some technical difficulties to realize the proposed switch. A problem is a large power loss occurring at the input and output ends of NPDM but it can be compensated by higher power laser. Although we have suggested the specific design for an efficient coupling of input beam, we need to confirm experimentally such an index-matching method.

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