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A Possible Configuration of Ultrafast Photonic Switching

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Various optical nonlinearities are examined from the viewpoint of applications to all-optical ultrafast optical switching elements (OSEs). It is pointed out that optical nonlinearities in the usual configurations, in which the pump and signal photon energies are both close to the band gap energy, do *not* satisfy some basic requirements for the nonlinear materials for the ultrafast OSEs. A general configuration, the two-photon-absorption-free inverse Raman process, is suggested as an possible configuration which satisfies all the basic requirements. The optical quantum-confined Stark effect (QCSE) in quantum-well structures (QWSs) corresponds to this configuration. It is pointed out that we can achieve practical ultrafast OSEs by using the optical QCSE in a high-quality GaAs/AlAs QWS with thin wells (~ 40 $\mathring{A}$).

1. Introduction and summary

There are two possibilities of applications of optical nonlinearities to all-optical signal processings[1]. One is parallel processors composed of optical switching elements (OSEs) with a slow responce time (μs typically). The other is ultrafast scalor processors (or parallel processors composed of these ultrafast scalor processors) composed of OSEs with an ultrafast (subpicosecond) responce time. This paper discusses a possible physical configuration of optical nonlinearity which can realize the latter application; the ultrafast OSEs. First, the necessary conditions for the nonlinear materials for the ultrafast OSEs are summarized.[1,2] Then, it is pointed out that any experimentally-observed optical nonlinearities in the usual configurations, in which the photon energies of the pump and signal pulses are both close to the band gap energy of the nonlinear materials, do *not* fulfill the necesary conditions. A general configuration which can fulfill all of those conditions is suggested. Finally, it is pointed out that the optical quantum confined Stark effect found by Frörich et al.[3] corresponds to the proposed configuration and does indeed satisfy all the conditions.

2. Nonlinear materials for the ultrafast OSEs

We consider an ultrafast OSE which utilizes a nonlinear refractive index change induced by an intense light field. We denote the photon energies of the pump and signal pulses $\hbar\omega_p$ and $\hbar\omega_s$, respectively, their intensities I_p and I_s, the absorption coefficient α_p and α_s, and the linear refractive index $n_p \approx n_s \equiv n$. Although in some applications ω and I of the pump and signal pulses are the same, we here treat a general case; *i.e.*, we assume that they are independently controlable. Let the third order susceptibility be $\chi^{(3)} \equiv \chi^{(3)}(\omega_s; -\omega_p, \omega_p, \omega_s)$, and the interaction length L. (For OSEs with an etalon cavity, L should be replaced by FL in the folowing equations, where F is the finess of the cavity.) The following properties are required for the nonlinear materials for the ultrafast OSEs.[1,2)]

1) Operates at room temperature.

2) The intrinsic responce time τ_i of the optical nonlinearity should be subpicosecond, *i.e.*, $\tau_i < 1ps$.

3) To get a sufficient modulation depth of the order of 100 %, $\chi^{(3)} I_p L / n_p n_s \lambda_s \geq 10^{-5}$, where $\chi^{(3)}$ and I_p are measured in esu and MW/cm^2, respectively.

4) To get a sufficient transmittance of the the order of 100 %, $\alpha_s L << 1$. (The cascadability condition)

5) A unit processing should be accomplished in a subpicosecond time scale, *i.e.*, $nL/c < 1ps$.

Note that condition 1 is a strong restriction in the sense that all of the conditions 2-5 can be easily fulfilled at low temperatures ($\sim 4.2K$) in, say, the $CuCl$ crystal near the exciton resonance. Note also that the last condition 5 is sometimes argued unnecessary because the processing for the next optical pulse(s) can be started before the processing for the previous pulse(s) is finished and consequently the *net* speed per unit processing can become very high. However, the processor in that case would be equivalent to a parallel processor because each scalor processing takes long time. We thus include condition 5 into the necessary conditions for ultrafast OSEs.

From condition 2, we have to avoid real excitations and utilize virtual excitations. Although τ_i of the order of ps is obtained for *real* excitations in organic molecules or intersubband *real* excitations in quantum-well structures (QWSs), such an fast responce is obtained *only* when the time intervals of pump pulses are long enough to prevent heating of the materials. That is, the short τ_i in these materials is due to fast phonon emission processes, so that the materials would heat up rapidly if the pump pulses arrive at a *practical* high bit rate. We therefore confine ourselves to virtual excitations.

3. Usual configuration

It seems widely accepted that low-dimensional systems (LDSs), such as QWSs, have an advantage over bulk crystals as nonlinear materials, because one can expect enhancements of $\chi^{(3)}$ due to quantum size effects. It should be noted, however, that under completely-off-resonance conditions (*i.e.*, when all of $\hbar\omega_p$, $\hbar\omega_s$, and the sum or difference of them are far off from any electronic excitation energies of the material) $\chi^{(3)}$ of *any* LDSs is of the same order in magnitude as that of the bulk crystals.[4] We thus have to utilize resonance enhancements of $\chi^{(3)}$ in LDSs, while keeping the photon energies slightly off from the electronic excitation energies in order to avoid real excitations.

In the usual configurations, both $\hbar\omega_p$ and $\hbar\omega_s$ are taken slightly below the band gap energy E_G in order to rely upon the resonance enhancement of $\chi^{(3)}$ due to excitons. In that case, the optical Stark effect[5,6] gives the largest $\chi^{(3)}$, which is $\sim 10^{-4}/\Delta_p\Delta_s^2\ esu$ for GaAs/AlGaAs multiple QWSs, where Δ_p and Δ_s are the pump and signal detuning energies measured in meV. Let us discuss whether all the conditions 1-5 can be fulfilled in this case. To do this, we have to consider the reduction of intrinsic speed $1/\tau_i$ caused by the real carrier generation, which *does* occur in some degree in actual systems. We denote the width and intervals of optical pulses τ_P and τ_I, respectively, and assume that $\tau_I \approx \tau_P$ in order to get a maximum processing rate. The real carrier density δN_{eh} excited by a single pump pulse is given by

$$\delta N_{eh} \approx I_p \alpha_p \tau_P / \hbar\omega_p. \tag{1}$$

Since the pulses come one after another with the time interval τ_I ($\sim 1\ ps$) which is much shorter than the recombination lifetime τ_R ($\sim 1\ ns$, typically), the total real carrier density N_{eh} is much larger than δN_{eh}. That is, N_{eh} as a function of time fluctuates around its average,

$$\overline{N_{eh}} \approx I_p \alpha_p \tau_R / 4\hbar\omega_p, \tag{2}$$

and the magnitude of the fluctuation is of the order of δN_{eh}. For changes in n caused by the real carriers to be much smaller than that due to virtually excited carriers,

$$Changes\ in\ n\ caused\ by\ \delta N_{eh} << 16\pi^2\chi^{(3)} I_p / n_p n_s c \tag{3}$$

should be satisfied. For the optical Stark effect, this is reduced to $\delta N_{eh} << N_{eh}^v$, where N_{eh}^v denotes the density of virtually excited carriers.[6] Furthermore, for the excitons to be stable,

$$\overline{N_{eh}} < N_{PSF}, \tag{4}$$

should also be satisfied,[6,7] where N_{PSF} denotes the phase-space filling density.[6,7] Similar equations can also be obtained for the real carriers generated by signal pulses. In that case, we have to note that each signal pulse must carry sufficient number of photons (say, 100) to reduce quantum noises.[1] As a result, I_s cannot be made arbitrary small, so that we have to take account of the real carrier generation by the signal pulses as well as that by the pump pulses.

To examine whether the above equations are satisfied in actual systems, not only the linear absorption but also the two-photon absorption (TPA) must be considered. That is, α_p is given by

$$\alpha_p = \alpha_p^{(0)} + \alpha_p^{(2)} I_p. \quad (5)$$

From experimental results for $\alpha^{(0)}$ and $\alpha^{(2)}$ of GaAs-AlGaAs QWSs, we can see that the above equations can *not* be satisfied simultaneously in the case of the optical Stark effect. That is, we have to take Δ_p large enough to be off from low-energy tail[8] (due to phonon-assisted absorption) of $\alpha_p^{(0)}$. In that case the resonance enhancement of $\chi^{(3)}$ becomes relatively weak, so that we have to increase I_p in compensation. Large I_p, however, results in large $\alpha_p^{(2)} I_p$, so that the real carrier generation becomes significant.[9] For example, for $I_p = 100MW/cm^2$, Eq.(4) requires that $\alpha_p < 0.1cm^{-1}$, whereas $\alpha_p^{(2)} I_p \sim 3cm^{-1}$. Note that this conclusion is quite general for optical nonlinearities in the usual configuration because the pump photons always suffer from $\alpha_p^{(2)}$ of the order of that of the bulk crystal.[10,11]

4. A possible configuration for ultrafast OSEs

We now suggest a physical configuration of ultrafast photonic switching which enables conditions 1-5 to be satisfied simultaneously. As seen in the preceeding section, we have to avoid the TPA of pump photons, while utilizing a resonance enhancement of $\chi^{(3)}$. This is possible in the following TPA-free inverse Raman configuration. That is, among many electronic excitations, we propose to utilize the three discrete levels; the ground state $|g\rangle$, and two excited states $|A\rangle$ and $|B\rangle$, which satisfy the following relations.

a) One-photon (OP) transition $|g\rangle \rightarrow |A\rangle$ is allowed and $\hbar\omega_s$ is taken close to its resonance; $\hbar\omega_s = E_A - \Delta_s$.

b) OP transition $|A\rangle \rightarrow |B\rangle$ is allowed and $\hbar\omega_p$ is taken close to its resonance; $\hbar\omega_p = E_B - E_A - \Delta_p$.

c) OP transitions $|g\rangle \rightarrow |any\ states\rangle$ by the pump light are forbidden by *some* selection rule(s).

d) Two-photon (TP) transitions $|g\rangle \rightarrow |any\ states\rangle$ by the pump light are forbidden by *some* selection rule(s).

Here, "some selection rule(s)" means any of the energy conservation, the parity selection rule, the angular momentum conservation, and so on. In this configuration, we can obtain large $\chi^{(3)}$ due to the triple resonance, without sufferring from any absorption of the pump photons (except for simultaneous absorption of the pump and signal photons discussed below). The most resonant term of $\chi^{(3)}$ is given by

$$\chi_{ijji}^{(3)} = \frac{|\langle B|P_j|A\rangle|^2 \, |\langle A|P_i|g\rangle|^2}{(\Delta_s - i\Gamma_A)^2(\Delta_s + \Delta_p - i\Gamma_B)}. \quad (6)$$

This $\chi^{(3)}$ is enhanced over the bulk crystal by many orders of magnitude, as demonstrated in an example given below. We can therefore achieve a triply-resonant large $\chi^{(3)}$ without generating real excitations, by appropriately taking Δ_p and Δ_s.

Many experiments on inverse Raman scattering processes which coincide with the above configuration were already performed. Most of them, however, were performed as a tool of nonlinear spectroscopy. The first basic experiment of photonic switching on the order of 1 *ps* in the above configuration has been reported only very recently by Kuwata et al.,[12] using the exciton and biexciton levels in $CuCl$ as states $|A\rangle$ and $|B\rangle$, respectively. Unfortunately, however, $CuCl$ does not satisfy condition 1, because both the exciton and biexciton in $CuCl$ are unstable at room temperature. Therefore, our task is to find systems posessing *discete* levels $|A\rangle$ and $|B\rangle$ (they needs to be discrete states in order to utilize resonance) which are not only in the TPA-free inverse Raman configuration but also stable at room temperture. We will point out in the following section that among many optical nonlinearities so far reported the optical QCSE (quantum confined Stark effect) found by Frörich et al.[3] is the only one which satisfies all of conditions 1-5.

5. The optical QCSE

In the optical QCSE,[3] the levels $|A\rangle$ and $|B\rangle$ correspond to the (c1h1,1S) exciton and (c2h1, 1S) exciton in an *undoped* QWS, respectively, where a (cnhm, 1S) exciton denotes the exciton associated with the nth electron and mth hole subbands. Owing to quantum confinement, they are stable even at room temperature.[11,13] It is clear that both OPA and TPA of the pump light are energetically forbidden.

Frörich et al.[3] tuned the energy difference between the above two excitons $E_{1S}^{c2h1} - E_{1S}^{c1h1}$ to $\hbar\omega_p$ of a CO_2 laser ($\lambda_p \simeq 11\ \mu m$), and they measured changes in the transmission spectra of the probe light which is polarized parallel to the QW layers. From the viewpoint of practical applications, such a long-wavelength pump light is difficult to generate and, what is worse, it would be absorbed by direct excitations of phonons. We here propose to use the pump light of a shorter wavelength, $\lambda_p \leq 3\ \mu m$. If the well thickness L_z is reduced to 40 Å, $E_{1S}^{c2h1} - E_{1S}^{c1h1} \simeq 450\ meV$, which corresponds to $\lambda_p \simeq 2.7\ \mu m$. The pump light of wavelength in this range can be generated by, say, the difference-frequency generation from YAG and semiconductor lasers.

For simplicity, we assume that the signal light is confined in a waveguide structure, and that both the pump and signal lights are polarized perpendicular to the QW layers, for which we have to consider light hole excitons only. For *40Å-GaAs–80Å-AlAs* MQWSs, $\chi^{(3)}$ is calculated as

$$\chi^{(3)} = \frac{4 \times 10^{-4}}{(\Delta_s - i\Gamma_{1S}^{c1lh1})^2(\Delta_s + \Delta_p - i\Gamma_{1S}^{c2lh1})}\ esu, \tag{7}$$

where Δ's and Γ's are measured in meV. From condition 5 and $n \sim 3.5$, we take $L = 80$ Å. Among various possibilities for the choice of Δ's, we will demonstarate one example. That is, we take $\Delta_s = 30\ meV$ and $\Delta_p = 3meV$. It should be noted that we have to reduce inhomogenious broadenings less than a few meV, which can be achieved by elaborate growth techniques.[14] Otherwise, $Im\ \chi^{(3)}$ becomes significantly large, which results in a large TPA of simultaneous absorption of signal and pump pulses. For extremly good samples thus obtained, we expect that Δ's are dominated by the phonon scatterings, and thus $\Gamma_{1S}^{c1lh1} \simeq \Gamma_{1S}^{c2lh1} \simeq 2\ meV$ at room temperature.[11,13] We then obtain $Re\chi^{(3)} \sim 1.5 \times 10^{-8} esu$, and $Im\chi^{(3)} \sim 3 \times 10^{-9} esu$. From condition 3 we find $I_p \sim 80\ MW/cm^2$, which is available by focusing a 7 W pulse into a $(3\ \mu m)^2$ area. As for absorption, from an experimantal result[8] for $\alpha_s^{(0)} \simeq 30\ cm^{-1} \times 2/3$ (the factor 2/3 takes account of the difference in the QW density) and the magnitude of *TPA of pairs of pump and signal photons* estimated from the above value of $Im\chi^{(3)}$, we find that about 35 % of signal photons are absorbed during their passage through the waveguide. From Eqs.(1)-(4), the resulting real carrier density can be shown to be irrelavant if we assume 100 photons per signal pulse and $\tau_P \simeq 1ps$. Unfortunately, however, condition 4 is satisfied only weakly, i.e., $\alpha_s L \sim 0.35$. Nontherless, we have achieved much improvement over the usual configuration, for which Eq.(4) would be completely violated owing to *TPA of pairs of pump photons*.

Finally, we would like to stress that more detailed studies of applications of the optical QCSE may open new applications, and that it is also important to search for other optical nonlinearities in the proposed configuration.

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