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Enhancement of absorptive optical nonlinearity of asymmetrically coupled quantum well structures*

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The third-order nonlinear susceptibility $\chi^{(3)}$ of asymmetrically coupled quantum well (ACQW) structure is calculated by employing a V-type three-level model. The efficiency of absorptive optical nonlinearity, which we define as the ratio of the third-order nonlinear susceptibility $\chi^{(3)}$ to the linear absorptive coefficient α_0, is analysed for different electronic coherence oscillation frequencies. We show that the efficiency is enhanced with the increase of the electronic coherence oscillation frequency between the two wells but is degenerated as that of a single quantum well when the electronic coherence oscillation frequency is zero. Compared with single quantum well, the ACQW designed has the property of nonlinear absorption and dispersion depending strongly on the external electric field along the growth direction. We predict that ACQW structures can provide both high efficiency limiters and controllable optical Kerr switch in future communication systems.

Keywords: asymmetrically coupled quantum well structures, third-order nonlinearity, effect of absorptive optical nonlinearity

PACC: 7320, 7340G, 7210, 7865

1. Introduction

Quantum interferences[1,2] via quantum-mechanical paths have become a subject of active study in the atomic system, which include electromagnetically induced transparency (EIT),[3,4] lasers without inversion (LWI),[5,6] and the highly efficient scheme for nonlinear optics.[3,7] In practical application, since appropriate parameters in this system cannot be controlled, a few schemes using intersubband transitions based on the asymmetrically coupled quantum well (ACQW) structures have been proposed.[8−10]

Currently, the ACQW structures for use in optoelectronic devices[11] are of considerable interest because the coupling effects between the two wells can provide a large quantum-confined Stark shift of the transition energy,[11,12] a large refractive index change and an enhanced absorption change.[13,14] From this point of view, high speed optical communication devices associated with second-order optical nonlinear susceptibility $\chi^{(2)}$ are easily obtained from the ACQW structures because the centrosymmetry is broken and the fast response is obtained from electronic coherence oscillation, [15,16] which includes second harmonic-wave generation (SHG), sum (and difference) frequency generation (SFG and DFG) and linear electric-optic (Pockels) effect. Other possible applications of ACQW structures in the optical communication system[17,18] are the limiters and the switch associated with the third-order nonlinear susceptibility $\chi^{(3)}$. The former corresponds to the absorptive saturation of materials and the latter depends on the optical Kerr effect of materials. In general, an important merit is the efficiency of absorptive optical nonlinearity, which is ratio of the third-order nonlinear susceptibility $\chi^{(3)}$ to the linear absorptive coefficient α_0 for a given input optical intensity, for revealing the nonlinear behavior of material.[19] Over the past, to our knowledge, studies in this field have been concentrated mainly on bulk materials, multiple quantum well structures and materials of quantum dots but scarcely on ACQW struc-

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tures.

In this paper, we will analyse the third-order nonlinear behaviours of ACQW structures employing a V-type three-level configuration. Section 2 gives the model and equations of motion. Section 2 gives the linear susceptibility $\chi^{(1)}$ and the third-order nonlinear susceptibility $\chi^{(3)}$ via perturbation chains. Section 4 is the analysis for the efficiency of absorptive optical nonlinearity associated with optical Kerr effect. Finally, Section 5 presents the conclusion.

2. Model and equations of motion

In general, ACQW structures consist of a wide well (WW), a narrow well (NW) and a thin barrier. When the barrier width L is less than the de Broglie wavelength of the tunnelling electron $\lambda = 2\pi\hbar/\sqrt{2m(U-K)}$, where K is the quantum-confinement energy of electron with quantum number $n = 1$ in the WW and U is the heterojunction barrier height depending strongly on negative bias voltage, the significant overlap of the lowest electronic functions between the two wells should be considered. In terms of the tight-bond approximation,[21] the system is considered to be a V-type three-level configuration, which consists of new stationary states (delocalized states) $|+\rangle = (|1\rangle + |2\rangle)/\sqrt{2}$, $|-\rangle = (|1\rangle - |2\rangle)/\sqrt{2}$ and heavy-hole state $|0\rangle$ (ground state). Here $|1\rangle$ and $|2\rangle$ are the lowest electronic states in the WW and in the NW, respectively. The corresponding level schematic diagram is shown in Fig.1. Assume that the two heavy-hole states in the valence-band are uncoupled due to large effective masses. For the electronic-states only coupled in the conduction-band, the energy separation follows anticrossing rules as a function of a

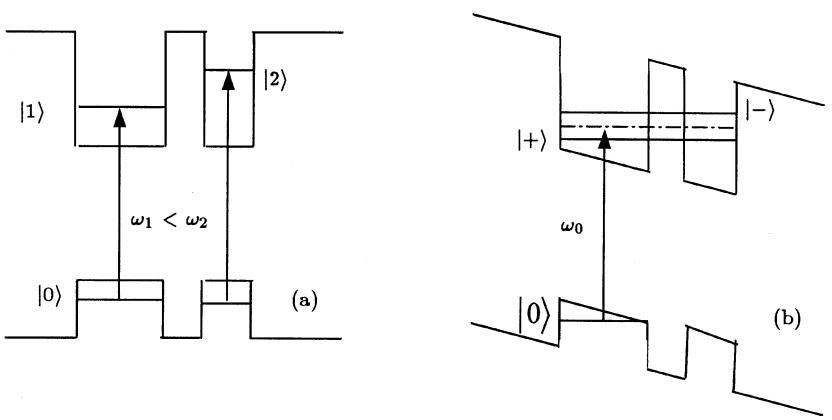


Fig.1. Level diagrams of asymmetrically couple quantum well structure: (a) uncoupled and (b) coupled. $|+\rangle$ and $|-\rangle$ are new stationary states given in text. ω_0 is the central transition frequency between the ground state $|0\rangle$ and the upper two states $|+\rangle$ and $|-\rangle$.

dc electric field along the growth direction of ACQW structures. At resonance, the corresponding angular frequency splitting $\omega_c = (E_- - E_+)/\hbar$ satisfies[20]

$$\omega_c = (\hbar k/mL)\mathrm{e}^{-kL}, \tag{1}$$

which denotes the electronic coherent oscillation between the two wells as a result of Fröhlich interaction.[13] Employing the V-type three-level model mentioned above, the density matrix element equations of this system interacting with external optical field $E(t)$ can be written as

$$\begin{aligned}\frac{\partial\rho_{00}}{\partial t} =& -\gamma_{00}\rho_{00} - \frac{\mathrm{i}}{\hbar}[\mu_{0+}\rho_{+0}E(t) - \mu_{0+}\rho_{+0}E^*(t)] \\ & - \frac{\mathrm{i}}{\hbar}[\mu_{0-}\rho_{-0}E(t) - \mu_{-0}\rho_{0-}E^*(t)] \\ & + \gamma_{00}^1 f_{00} + \gamma_{00}^{\mathrm{h}} f_{00}^{\mathrm{h}} + \gamma_{00}^{\mathrm{s}} f_{00}^{\mathrm{eq}}, \end{aligned} \tag{2a}$$

$$\begin{aligned}\frac{\partial\rho_{++}}{\partial t} =& -\gamma_{++}\rho_{++} + \frac{\mathrm{i}}{\hbar}[\mu_{+0}\rho_{0+}E(t) - \mu_{+0}\rho_{0+}E^*(t)] \\ & + \frac{\mathrm{i}}{\hbar}[\mu_{-+}\rho_{+-}E(t) - \mu_{+-}\rho_{-+}E^*(t)] \\ & + \gamma_{++}^1 f_{++} + \gamma_{++}^{\mathrm{h}} f_{++}^{\mathrm{h}} + \gamma_{++}^{\mathrm{s}} f_{++}^{\mathrm{eq}}, \end{aligned} \tag{2b}$$

$$\begin{aligned}\frac{\partial\rho_{--}}{\partial t} =& -\gamma_{--}\rho_{--} + \frac{\mathrm{i}}{\hbar}[\mu_{-0}\rho_{0-}E(t) - \mu_{-0}\rho_{0-}E^*(t)] \\ & + \frac{\mathrm{i}}{\hbar}[\mu_{-+}\rho_{+-}E(t) - \mu_{+-}\rho_{-+}E^*(t)] \\ & + \gamma_{--}^1 f_{--} + \gamma_{--}^{\mathrm{h}} f_{--}^{\mathrm{h}} + \gamma_{--}^{\mathrm{s}} f_{--}^{\mathrm{eq}}, \end{aligned} \tag{2c}$$

$$\begin{aligned}\frac{\partial\rho_{-+}}{\partial t} =& -(\gamma_{-+} + \mathrm{i}\omega_c)\rho_{-+} \\ & - \frac{\mathrm{i}}{\hbar}[\mu_{-+}\rho_{++}E(t) - \mu_{+-}\rho_{--}E^*(t)] \\ & - \frac{\mathrm{i}}{\hbar}[\mu_{-0}\rho_{0-}E(t) - \mu_{0+}\rho_{+0}E^*(t)], \end{aligned} \tag{2d}$$

$$\frac{\partial\rho_{+0}}{\partial t} = -[\gamma_{+0}+\mathrm{i}(\omega_0-\omega_c/2)]\rho_{+0} - \frac{\mathrm{i}}{\hbar}[\mu_{0+}\rho_{++}E(t)-\mu_{+0}\rho_{00}E^*(t)] + \frac{\mathrm{i}}{\hbar}[\mu_{+-}\rho_{+0}E(t)-\mu_{+0}\rho_{+-}E^*(t)], \quad (2e)$$

$$\frac{\partial\rho_{-0}}{\partial t} = -[\gamma_{-0}+\mathrm{i}(\omega_0+\omega_c/2)]\rho_{-0} - \frac{\mathrm{i}}{\hbar}[\mu_{0-}\rho_{--}E(t)-\mu_{-0}\rho_{00}E^*(t)] + \frac{\mathrm{i}}{\hbar}[\mu_{+-}\rho_{-0}E(t)-\mu_{-0}\rho_{-+}E^*(t)], \quad (2f)$$

$$\rho_{+-} = \rho_{-+}^*, \quad (2g)$$

$$\rho_{0-} = \rho_{-0}^*, \quad (2h)$$

$$\rho_{0+} = \rho_{+0}^*, \quad (2i)$$

where γ_{ii} and γ_{ij} represent, respectively, the population relaxation rate and dephasing rate of the transition dipole moment $\mu_{ij}(i,j = 0,\pm$ and $i \neq j)$. The trace of the density matrix is conserved, namely $\rho_{00}+\rho_{--}+\rho_{++}=1$. Considering $\gamma_{ii} \ll \gamma_{ij}$, we will omit the terms associated with γ_{ii} in the following.

3. Linear and nonlinear susceptibility

For an electric field $E(t)$ given by

$$E(t) = \frac{1}{2}[E(\omega)\mathrm{e}^{-\mathrm{i}\omega t}+\mathrm{c.c.}], \quad (3)$$

we expand density matrix in powers as

$$\rho_{ij} = \rho_{ij}^{(0)} + \rho_{ij}^{(1)} + \rho_{ij}^{(2)} + \cdots. \quad (4)$$

Assuming that the initial condition $\rho_{00}^{(0)}(-\infty)=1$ and other matrix elements are zero, the first-order transition dipole matrix elements in the $\{|0\rangle, |\pm\rangle\}$ representation can be obtained, respectively, as

$$\rho_{+0}^{(1)}(t) = \frac{\mathrm{i}\mu_{+0}}{\hbar}\int_{-\infty}^{t}\mathrm{d}t'\exp\{-[\gamma_{+0}+\mathrm{i}(\omega_0-\omega_c/2)(t-t')]\}E(t'), \quad (5a)$$

$$\rho_{-0}^{(1)}(t) = \frac{\mathrm{i}\mu_{-0}}{\hbar}\int_{-\infty}^{t}\mathrm{d}t'\exp\{-[\gamma_{-0}+\mathrm{i}(\omega_0+\omega_c/2)(t-t')]\}E(t'). \quad (5b)$$

Furthermore, by perturbation chains $\rho_{0+}^{(1)} \to (\rho_{++}^{(2)}, \rho_{--}^{(2)}) \to \rho_{+0}^{(3)}$ and $\rho_{0-}^{(1)} \to (\rho_{--}^{(2)}, \rho_{++}^{(2)}) \to \rho_{-0}^{(3)}$, the third-order transition dipole density matrix elements $\rho_{+0}^{(3)}$ and $\rho_{-0}^{(3)}$ are expressed, respectively, as

$$\begin{aligned}\rho_{+0}^{(3)}(t) =& (\mathrm{i}\mu_{+0}/\hbar)^3\int_{-\infty}^{t}\mathrm{d}t_3\int_{-\infty}^{t_3}\mathrm{d}t_2\int_{-\infty}^{t_2}\mathrm{d}t_1\exp[-\gamma_{+0}(t-t_3+t_2-t_1)]\\ &\times\{\{\exp[-\gamma_{00}(t_3-t_2)-\mathrm{i}(\omega_0-\omega_c/2)(t-t_3+t_2-t_1)]E(t_3)E^*(t_2)E(t_1)\\ &-\exp[-\gamma_{++}(t_3-t_2)-\mathrm{i}(\omega_0-\omega_c/2)(t-t_3-t_2+t_1)]\}E(t_3)E(t_2)E^*(t_1)\\ &+(\mathrm{i}/\hbar)^3|\mu_{-0}|^2\mu_{+0}\int_{-\infty}^{t}\mathrm{d}t_3\int_{-\infty}^{t_3}\mathrm{d}t_2\int_{-\infty}^{t_2}\mathrm{d}t_1\exp[-\gamma_{+0}(t-t_3)-\gamma_{00}(t_3-t_2)-\gamma_{-0}(t_2-t_1)]\\ &\times\{\exp[-\mathrm{i}(\omega_0+\omega_c/2)(t-t_3+t_2-t_1)]E(t_3)E(t_2)E^*(t_1)\\ &+\exp[-\mathrm{i}(\omega_0+\omega_c/2)(t-t_3-t_2+t_1)]E(t_3)E^*(t_2)E(t_1)\},\end{aligned} \quad (6a)$$

and

$$\begin{aligned}\rho_{-0}^{(3)}(t) =& (\mathrm{i}\mu_{-0}/\hbar)^3\int_{-\infty}^{t}\mathrm{d}t_3\int_{-\infty}^{t_3}\mathrm{d}t_2\int_{-\infty}^{t_2}\mathrm{d}t_1\exp[-\gamma_{-0}(t-t_3+t_2-t_1)]\\ &\times\{\{\exp[-\gamma_{00}(t_3-t_2)-\mathrm{i}(\omega_0+\omega_c/2)(t-t_3+t_2-t_1)]E(t_3)E^*(t_2)E(t_1)\\ &-\exp[-\gamma_{--}(t_3-t_2)-\mathrm{i}(\omega_0+\omega_c/2)(t-t_3-t_2+t_1)]\}E(t_3)E(t_2)E^*(t_1)\\ &+(\mathrm{i}/\hbar)^3|\mu_{+0}|^2\mu_{-0}\int_{-\infty}^{t}\mathrm{d}t_3\int_{-\infty}^{t_3}\mathrm{d}t_2\int_{-\infty}^{t_2}\mathrm{d}t_1\exp[-\gamma_{-0}(t-t_3)-\gamma_{00}(t_3-t_2)-\gamma_{+0}(t_2-t_1)]\\ &\times\{\exp[-\mathrm{i}(\omega_0-\omega_c/2)(t-t_3+t_2-t_1)]E(t_3)E(t_2)E^*(t_1)\\ &+\exp[-\mathrm{i}(\omega_0-\omega_c/2)(t-t_3-t_2+t_1)]E(t_3)E^*(t_2)E(t_1)\}.\end{aligned} \quad (6b)$$

Employ the polarization formula

$$P(t) = (1/V)\mathrm{tr}\mu\rho(t). \tag{7}$$

Here V is the active layer volume, μ is the dipole operator and ρ the density matrix operator, whose linear and nonlinear off-diagonal elements are determined by (5) and (6), respectively. Furthermore, defining the linear and the third-order nonlinear susceptibilities, respectively, as

$$P^{(1)}(\omega) = \frac{1}{2}[\varepsilon_0\chi^{(1)}(\omega)E(\omega) + \text{c.c.}], \tag{8a}$$

$$P^{(3)}(\omega_s) = \frac{1}{2}[\chi^{(3)}(\omega_s; -\omega, -\omega, \omega)|E(\omega)|^2E(\omega) + \text{c.c.}], \tag{8b}$$

we have

$$\chi^{(1)}(\omega) = \frac{\mathrm{i}}{\varepsilon_0\hbar V}\left[\frac{|\mu_{+0}|^2}{\gamma_{+0} + \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)} + \frac{|\mu_{-0}|^2}{\gamma_{-0} + \mathrm{i}(\omega_0 + \omega_\mathrm{c}/2 - \omega)}\right], \tag{9a}$$

$$\begin{aligned}
\chi^{(3)}(\omega_s; -\omega, -\omega, \omega) = &-\frac{\mathrm{i}|\mu_{-0}|^4}{\varepsilon_0\hbar^3 V}\frac{1}{\gamma_{-0} + \mathrm{i}(\omega_0 + \omega_\mathrm{c}/2 - \omega)} \\
&\times\left[\frac{1}{\gamma_{-0} - \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)} - \frac{1}{\gamma_{-0} + \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)}\right. \\
&\left.+\frac{1}{\gamma_{-0} + \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)} - \frac{1}{\gamma_{-0} + \mathrm{i}(\omega_0 + \omega_\mathrm{c}/2 + \omega)}\right] \\
&-\frac{\mathrm{i}|\mu_{+0}|^4}{\varepsilon_0\hbar^3 V}\frac{1}{\gamma_{+0} + \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)} \\
&\times\left[\frac{1}{\gamma_{+0} - \mathrm{i}(\omega_0 + \omega_\mathrm{c}/2 - \omega)} - \frac{1}{\gamma_{+0} - \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)}\right. \\
&\left.+\frac{1}{\gamma_{+0} + \mathrm{i}(\omega_0 + \omega_\mathrm{c}/2 - \omega)} - \frac{1}{\gamma_{+0} + \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)}\right] \\
&-\frac{\mathrm{i}|\mu_{+0}|^2|\mu_{-0}|^2}{\varepsilon_0\hbar^3 V}\frac{1}{\gamma_{+0} + \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)} \\
&\times\left[\frac{1}{\gamma_{-0} - \mathrm{i}(\omega_0 + \omega_\mathrm{c}/2 - \omega)} + \frac{1}{\gamma_{-0} + \mathrm{i}(\omega_0 + \omega_\mathrm{c}/2 - \omega)}\right] \\
&-\frac{\mathrm{i}|\mu_{+0}|^2|\mu_{-0}|^2}{\varepsilon_0\hbar^3 V}\frac{1}{\gamma_{-0} + \mathrm{i}(\omega_0 + \omega_\mathrm{c}/2 - \omega)} \\
&\times\left[\frac{1}{\gamma_{+0} - \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)} + \frac{1}{\gamma_{+0} + \mathrm{i}(\omega_0 - \omega_\mathrm{c}/2 - \omega)}\right].
\end{aligned} \tag{9b}$$

Here ε_0 is the dielectric constant in vacuum, the terms related to $|\mu_{-0}|^4$ and $|\mu_{+0}|^4$, respectively, are the nonlinear absorption and dispersion as a result of $|0\rangle \to |-\rangle$ and $|0\rangle \to |+\rangle$ transitions, and the terms related to $|\mu_{-0}|^2$ and $|\mu_{+0}|^2$ correspond to nonlinear absorption and dispersion caused by transition paths $|-\rangle \to |0\rangle$ and $|+\rangle \to |0\rangle$. In order to explore the features of linear and nonlinear absorptions from (9), we must calculate the values, or at least the relative values of μ_{-0} and μ_{+0}. Making change of presentation, we arrive immediately at

$$\mu_{+0} = \frac{1}{\sqrt{2}}(\mu_{10} + \mu_{20}), \tag{10a}$$

$$\mu_{-0} = \frac{1}{\sqrt{2}}(\mu_{10} - \mu_{20}), \tag{10b}$$

where μ_{10} and μ_{20} are the transition dipole matrix elements for the isolated WW and for the isolated NW, respectively. Since the two heavy-hole levels in the two isolated wells are non-coupled, the squared ratio of μ_{-0} to μ_{+0} is obtained as

$$|\mu_{-0}/\mu_{+0}|^2 = |(L_\mathrm{W} - L_\mathrm{N})/(L_\mathrm{W} + L_\mathrm{N})|^2 \tag{11}$$

from the envelope function approximation. Here we assume the z coordinate of electron is equal to that of heavy-hole in the same isolated well.[20] This result shows that the relative intensity of linear and nonlinear absorption for $|0\rangle \to |-\rangle$ and $|0\rangle \to |+\rangle$ depends strongly on the relative width between the two wells. As seen easily from (11), the absorption of $|0\rangle \to |+\rangle$ is generally much stronger than that of $|0\rangle \to |-\rangle$. For a special ACQW structure with $L_\mathrm{N} = L_\mathrm{W}$, we have $|\mu_{-0}/\mu_{+0}|^2 = 0$. This means the coherent absorption for a symmetrically coupled quantum well structure is only determined by $|0\rangle \to |+\rangle$ transition.

Clearly, the electronic coherent oscillation vanishes because the ACQW structures become symmetrical in this case. From the quantum mechanic point of view, the intersubband transition $|-\rangle \rightarrow |+\rangle$ is forbidden. The absorptive optical nonlinearity is similar to that of a single quantum well.

4. Analysis and discussion

In this section, we will analyse the efficiency of absorptive optical nonlinearities for the ACQW structures as a result of external optical field power. Denoting the incident power density by I, the nonlinear change of α related to the imaginary part of $\chi^{(3)}$ is

$$\alpha = \alpha_0 + \alpha_2 I, \tag{12a}$$

$$\alpha_2 = \frac{3\omega}{4n_0 c}\mathrm{Im}\chi^{(3)}, \tag{12b}$$

In the same way, the associated change of the refractive index n is related to the real part of $\chi^{(3)}$ as

$$n = n_0 + n_2 I, \tag{13a}$$

$$n_2 = \frac{3}{8n_0}\mathrm{Re}\chi^{(3)}. \tag{13b}$$

Here

$$\alpha_0(\omega) = \frac{\omega}{n_0 c}\mathrm{Im}[\chi^{(1)}(\omega)], \tag{14a}$$

$$n_0(\omega) = 1 + \frac{1}{2}\mathrm{Re}[\chi^{(1)}(\omega)], \tag{14b}$$

n_0 and α_0 represent the linear absorption coefficient and the refractive index of ACQW structures, respectively, c is the velocity of light in vacuum.

Although the transition between $|0\rangle$ and $|+\rangle$ is only resonantly excited by a narrow band laser pulse, the electrons in the $|+\rangle$ state can transfer to the $|-\rangle$ state and return to the $|+\rangle$ state as a result of LO-assisted photon. This coherent oscillation of electrons in the conduction band will lead to a large change of optical nonlinearity compared with that of single quantum well. In order to study the nonlinear properties of ACQW structures, we defined a fundamental quantity $\chi^{(3)}/\alpha_0$, namely the ratio of the third-order nonlinear susceptibility $\chi^{(3)}$ to the linear absorptive coefficient α_0, as the measure of the efficiency of optical absorptive nonlinearity.[19]

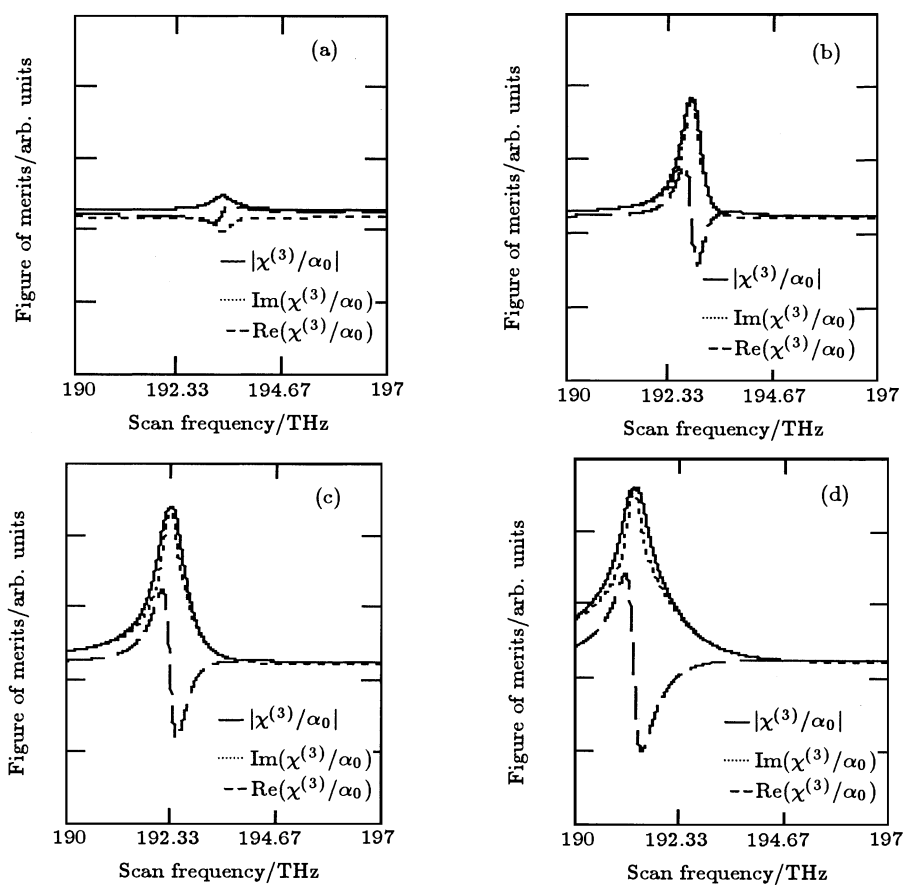


Fig.2. Figure of merit of the third-order optical nonlinearity $\chi^{(3)}/\alpha_0$ for oscillation frequency ω_c. The parameters used are γ_{+0} =0.20ps^{-1}, γ_{-0} =0.25ps^{-1} and $|\mu_{-0}/\mu_{+0}| = 1/3$. In the above figures, (a) ω_c = 0THz, (b) ω_c = 1THz, (c) ω_c = 2THz and (d) ω_c = 4THz.

For simplicity, an ACQW structure consisting of InGaAs–InGaAsP heterostructure tailored by energy-band engineering is only considered.[14,15] To this structure, the central transition wavelength between the upper two levels and the lower level exactly arrives at 1550nm (transition frequency 193.3THz) in the optical communication window. When the ACQW structures designed is applied along the growth direction, the frequency ω_c of electronic coherence oscillation is changed as the negative bias voltage increases. In Fig.2, we plot the figure of merit $\chi^{(3)}/\alpha_0$ determined by (9) and (14a) for different values of ω_c, where (a), (b), (c) and (d) correspond to ω_c =0, 1, 2 and 4THz, respectively. As can be seen from this figure, when the oscillation frequency is zero, the third-order nonlinear properties of ACQW structures are the same as that of single quantum well. In other words, in this case the quantum interference nonlinear effect of ACQW structures vanishes. However, when oscillation frequency is ω_c =1THz, the peak of nonlinear absorption is shifted to 192.8 THz, the efficiency of absorptive optical nonlinearity is remarkably enhanced, as shown in Fig.2(b). With the increase of ω_c, these changes are further enlarged, as illustrated in Fig.2(c) and (d). This means that electronic coherent oscillation can give rise to enchantment of absorptive optical nonlinearity. From the practical point of view, the ACQW structures can be applicable to highly efficient limiters and optical switch operating at the wavelength of 1550nm.

5. Conclusion

The third-order nonlinear susceptibility associated with optical Kerr effect is obtained from a V-type three-level configuration of ACQW structures. The efficiency of absorptive optical nonlinearity for the different electronic coherence oscillation frequencies is analysed. Our results show that the efficiency depends strongly on the electronic oscillation frequency. When the electronic oscillation frequency is equal to zero, the nonlinear optical properties of ACQW structures are similar to that of the single quantum well. However, when the oscillation frequency is nonzero, the efficiency of absorptive optical nonlinearity is dramatically changed and the nonlinear refractive index is remarkably enhanced with increase of electronic oscillation frequency. Clearly, the efficiency is only determined by negative bias voltage for the ACQW structures designed. In this sense, the large changes of ACQW structures associated with biased voltage can provide the highly efficient limiters and the controllable optical Kerr switch for optical communication systems.

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