Giant Kerr nonlinearity via tunneling induced double dark resonances in triangular quantum dot molecules

2015 Laser Phys. Lett. 12 125203

(http://iopscience.iop.org/1612-202X/12/12/125203)

IOP Publishing | Astro Ltd

Laser Physics Letters

Laser Phys. Lett. **12** (2015) 125203 (6pp)

doi:10.1088/1612-2011/12/12/125203

Giant Kerr nonlinearity via tunneling induced double dark resonances in triangular quantum dot molecules

Si-Cong Tian[1], Ren-Gang Wan[2], Cun-Zhu Tong[1], Xi-Hong Fu[1], Jun-Sheng Cao[1] and Yong-Qiang Ning[1]

[1] State Key Laboratory of Luminescence and Applications, Changchun Institute of Optics, Fine Mechanics and Physics, Chinese Academy of Sciences, Changchun 130033, People's Republic of China
[2] School of Physics and Information Technology, Shaanxi Normal University, Xi'an 710062, People's Republic of China

E-mail: tongcz@ciomp.ac.cn and tiansicong@ciomp.ac.cn

Received 24 October 2014
Accepted for publication 25 October 2015
Published 10 November 2015

Abstract

A scheme for giant Kerr nonlinearity via tunneling in triangular triple quantum dot molecules is proposed. In such a system, the linear absorption and the Kerr nonlinearity depend critically on the energy splitting of the excited states and the tunneling intensity. With proper parameters, giant Kerr nonlinearity accompanied by vanishing absorption can be realized. The enhancement of Kerr nonlinearity is attributed to the interacting double dark resonances induced by the tunneling between the quantum dots, requiring no extra coupling laser fields.

Keywords: Kerr nonlinearity, tunneling induced double dark resonances, triple quantum dots

(Some figures may appear in colour only in the online journal)

1. Introduction

Kerr nonlinearity, which corresponds to the refractive part of the third-order susceptibility in optical media, plays an important role in the field of nonlinear optics [1]. Recent studies have shown that Kerr nonlinearity can be used for many interesting applications, such as self focusing [2], optical solitons [3–5] and polarization phase gates [6–8]. It is desirable to have large nonlinear susceptibilities under the condition of low light levels. The electromagnetically induced transparency (EIT) scheme is capable of producing enhanced Kerr nonlinearity and suppressing the linear absorption [9]. Therefore, the enhanced Kerr nonlinearity in three-level EIT systems has been theoretically predicted [10] and experimentally measured [11]. Furthermore, in multiple-level atomic schemes (such as N- [12, 13], M- [14], four-level Λ [15, 16] and inverted-Y [17] systems), Kerr nonlinearity can be greatly enhanced. The large enhancement of nonlinear susceptibilities in EIT media led to the study of nonlinear optics at low light levels [18–20].

On the other hand, quantum dots (QDs) have higher nonlinear optical coefficients than atoms, therefore, they can be used for obtaining high Kerr nonlinearity [21, 22]. But in such systems, it is crucial to have additional coupling lasers to modify the linear and nonlinear optical properties. Quantum dot molecules (QDMs) are systems composed of two or more closely spaced and interacting QDs, and can be fabricated by using self-assembled dot growth technology [23]. The tunneling between the dots, which is controlled by an external electric field, can induce quantum interference and coherence [24–26]. Therefore QDMs have received much attention in EIT and slow light [27, 28], entanglement [29, 30], coherent population transfer [31], narrowing of fluorescence spectrum [32, 33], optical bistability [34], high order nonlinearity [35], phase gate [36] and linewidth narrowing of cavity [37, 38].

It is known that dark resonance is the basis for EIT. When a dark state is coherently coupled to another level by a coupling laser, double dark resonances can be attained [39, 40]. Intrigued by these studies, in the present paper we demonstrate a giant Kerr nonlinearity via tunneling–induced, interacting,

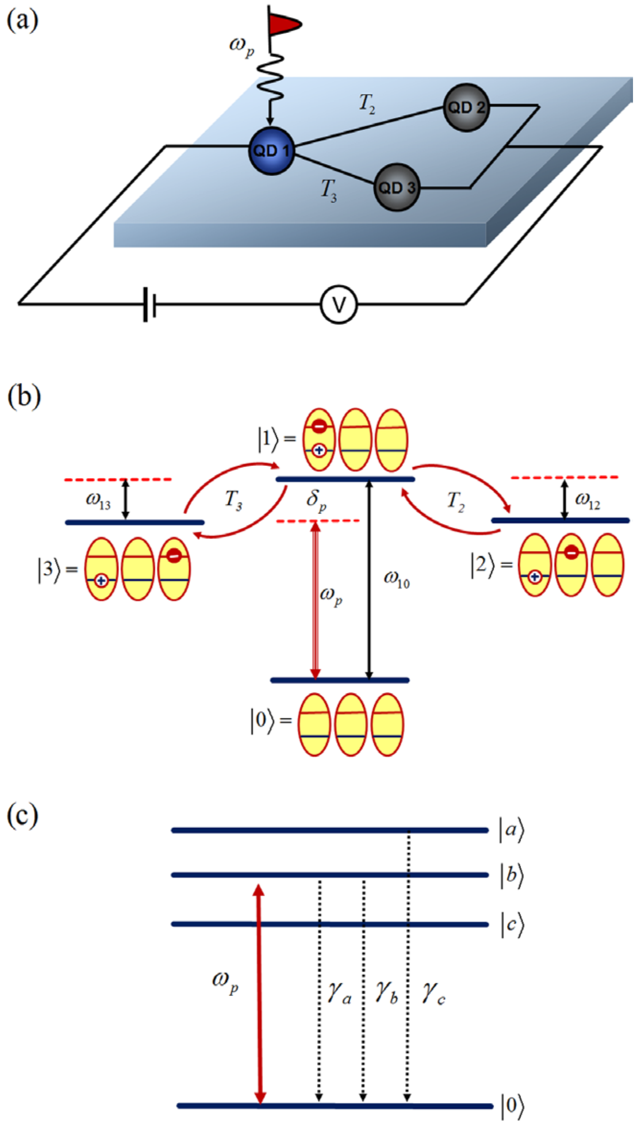

Figure 1. (a) Schematic of the setup of a triangular TQD. The probe field transmits the QD 1. V is a bias voltage. (b) The schematic of the level configuration of a triangular TQD. (c) Dressed states of a triangular TQD for two tunneling couplings.

double dark resonances in triangular, triple quantum, dot molecules (TQDs). Such TQDs can be fabricated experimentally [41–43]. The tunneling between the QDs can modify excitons and create a four-level tripod-type system [44]. With proper values of the energy splitting and the tunneling intensity, enhancement of Kerr nonlinearity accompanied by vanishing absorption is realized. The difference between our study and previous studies on Kerr nonlinearity is that, in our consideration, the interacting double dark resonances are induced by the tunneling effect, and no coupling laser fields are required.

2. Model and equations

The schematic of the setup of the triangular TQDs is shown in figure 1(a). QD 1 and QD 2, QD 1 and QD 3 are connected by two gate electrodes. Without a gate voltage, the conduction-band electron levels are out of resonance and the electron tunneling between the QDs is very weak. However, with a gate voltage, the conduction-band electron levels closely approach resonance, and the electron tunneling between the QDs is greatly enhanced. Because of the far off-resonant valence-band energy levels, hole tunneling can be neglected. Then the schematic of the level configuration of the triangular TQDs can be drawn as shown in figure 1(b). The ground state $|0\rangle$ has no excitations and the exciton state $|1\rangle$ has one electron-hole pair in QD 1. Under the tunneling coupling conditions, the electron can tunnel from QD 1 to QD 2 or from QD 1 to QD 3. Therefore, the indirect exciton state $|2\rangle$ has one hole in QD 1 and one electron in QD 2, and the indirect exciton state $|3\rangle$ has one hole in QD 1 and one electron in QD 3.

The Hamiltonian of the basis $\{|0\rangle, |1\rangle, |2\rangle, |3\rangle\}$ under the rotating-wave and the electric-dipole approximations can be written as (assumption of $\hbar = 1$):

$$H_I = \begin{pmatrix} 0 & -\Omega_p & 0 & 0 \\ -\Omega_p & \delta_p & -T_2 & -T_3 \\ 0 & -T_2 & \delta_p - \omega_{12} & 0 \\ 0 & -T_3 & 0 & \delta_p - \omega_{13} \end{pmatrix}. \tag{1}$$

Here, $\Omega_p = \mu_{01} E_p$ is the Rabi frequency of the transition $|0\rangle \to |1\rangle$. E_p denotes the electric-field amplitude of the laser, and $\mu_{01} = \boldsymbol{\mu}_{01} \cdot \mathbf{e}$ denotes the electric dipole moment for the excitonic transition between states $|0\rangle$ and $|1\rangle$, with $\mathbf{e}$ being the polarization vector. T_2 and T_3 are the tunneling couplings, which depend on the barrier characteristics and the external electric field. $\delta_p = \omega_{10} - \omega_p$ is the detuning of the probe field, with ω_{10} being the transition frequency between $|1\rangle$ and $|0\rangle$ states.

ω_{12} and ω_{13} are the energy splitting of the excited states and can be controlled by manipulation of the external electric field that changes the effective confinement potential.

The state vector at any time t is:

$$|\Psi_I(t)\rangle = a_0(t)|0\rangle + a_1(t)|1\rangle + a_2(t)|2\rangle + a_3(t)|3\rangle. \tag{2}$$

The evolution of the state vector obeys the Schrödinger equation:

$$\frac{d}{dt}|\Psi_I(t)\rangle = -\mathrm{i}H_I(t)|\Psi_I(t)\rangle. \tag{3}$$

Substituting equations (1) and (2) into equation (3) and then using the Weisskopf-Wigner theory [45, 46], the dynamic equations for atomic probability amplitudes in the interaction picture can be obtained:

$$\mathrm{i}\dot{a}_0 = -\Omega_p a_1, \tag{4a}$$

$$\mathrm{i}\dot{a}_1 = -\Omega_p a_0 - T_2 a_2 - T_3 a_3 + (\delta_p - \mathrm{i}\gamma_1) a_1, \tag{4b}$$

$$\mathrm{i}\dot{a}_2 = -T_2 a_1 + (\delta_p - \omega_{12} - \mathrm{i}\gamma_2) a_2, \tag{4c}$$

$$\mathrm{i}\dot{a}_3 = -T_3 a_1 + (\delta_p - \omega_{13} - \mathrm{i}\gamma_3) a_3, \tag{4d}$$

with $|a_0|^2 + |a_1|^2 + |a_2|^2 + |a_3|^2 = 1$. And here $\gamma_i = \frac{1}{2}\Gamma_{i0} + \gamma_{i0}^d \, (i = 1-3)$ is the typical effective decay rate, with Γ_{i0} being the radiative decay rate of populations from $|i\rangle \to |0\rangle$ and γ_{i0}^{d} being the pure dephasing rates.

It is well known that the polarization of the medium is:

$$P = \varepsilon_0 \chi_p E_p = \frac{\Gamma}{V}\mu_{01} a_1 a_0^*, \tag{5}$$

with Γ being the optical confinement factor, V being the volume of a single QD, and ε_0 being the dielectric constant [47]. From equation (5) the probe susceptibility can be obtained, which is given by:

$$\chi_p = \frac{\Gamma}{V}\frac{{\mu_{01}}^2}{\varepsilon_0 \hbar \Omega_p} a_1 a_0^* = \frac{\Gamma}{V}\frac{{\mu_{01}}^2}{\varepsilon_0 \hbar}\chi. \tag{6}$$

For steady state, we solve equations (4) under the weak field approximation ($|a_0|^2 = 1$), then:

$$\chi = \frac{1}{\Gamma_1 - \frac{T_2^2}{\Gamma_2} - \frac{T_3^2}{\Gamma_3}} \frac{1}{1 + \frac{\Omega_p^2}{|\Gamma_1 - \frac{T_2^2}{\Gamma_2} - \frac{T_3^2}{\Gamma_3}|^2}\left(1 + \frac{T_2^2}{|\Gamma_2|^2} + \frac{T_3^2}{|\Gamma_3|^2}\right)}, \tag{7}$$

where $\Gamma_1 = \delta_p - i\gamma_1$, $\Gamma_2 = \delta_p - \omega_{12} - i\gamma_2$, and $\Gamma_3 = \delta_p - \omega_{13} - i\gamma_3$. What we are interested in is the Kerr nonlinearity. So we expand the probe susceptibility, χ, into the second order of Ω_p using the Maclaurin formula:

$$\chi = \chi^{(1)} + \chi^{(3)}{\Omega_p}^2, \tag{8}$$

where $\chi^{(1)}$ and $\chi^{(3)}$ correspond to the first-order linear and third-order Kerr nonlinear parts of the susceptibility, respectively, and they are given by:

$$\chi^{(1)} = \frac{1}{\Gamma_1 - \frac{T_2^2}{\Gamma_2} - \frac{T_3^2}{\Gamma_3}}, \tag{9a}$$

$$\chi^{(3)} = -\frac{1}{\Gamma_1 - \frac{T_2^2}{\Gamma_2} - \frac{T_3^2}{\Gamma_3}} \frac{1}{|\Gamma_1 - \frac{T_2^2}{\Gamma_2} - \frac{T_3^2}{\Gamma_3}|^2}\left(1 + \frac{T_2^2}{|\Gamma_2|^2} + \frac{T_3^2}{|\Gamma_3|^2}\right). \tag{9b}$$

3. Results and discussions

In QDMs, the tunneling, T_2, and, T_3, depend on the barrier characteristics and the external electric field. Energy splitting ω_{12} and ω_{13} can be controlled by manipulation of the external electric field. In addition, in the low temperature regime, both the population decay rates and the dephasing rates should be considered. The realistic values of these parameters are according to [28] and Refs therein. And, for simplicity, all the parameters are scaled by the decay rate, γ_1.

We first consider the case in which the energy splitting is zero ($\omega_{12} = \omega_{13} = 0$) and display the linear absorption, Im[$\chi^{(1)}$], and the Kerr nonlinearity, Re[$\chi^{(3)}$], as a function of probe detuning with and without tunneling, T_3, in figure 2. In the absence of tunneling, T_3, the electron can tunnel from QD 1 to QD 2 , but can not tunnel from QD 1 to QD 3. Thus, one can obtain one transparency window as shown by the dotted line in figure 2(a). While in the presence of tunneling, T_3, the electrons

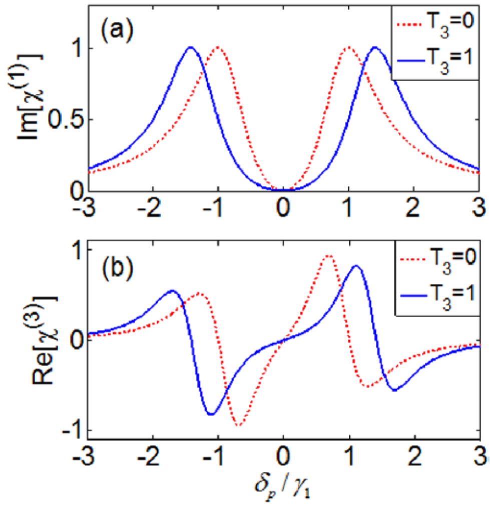

Figure 2. Variation of Re[$\chi^{(3)}$] and Im [$\chi^{(1)}$] as a function of the probe detuning δ_p with $T_3 = 0$ (dotted line) and $T_3 = 1$ (solid line), (a) Im [$\chi^{(1)}$], (b) Re[$\chi^{(3)}$]. Other parameters are $T_2 = 1$, $\omega_{12} = \omega_{13} = 0$, $\gamma_1 = 1$, $\gamma_2 = \gamma_3 = 10^{-3}$.

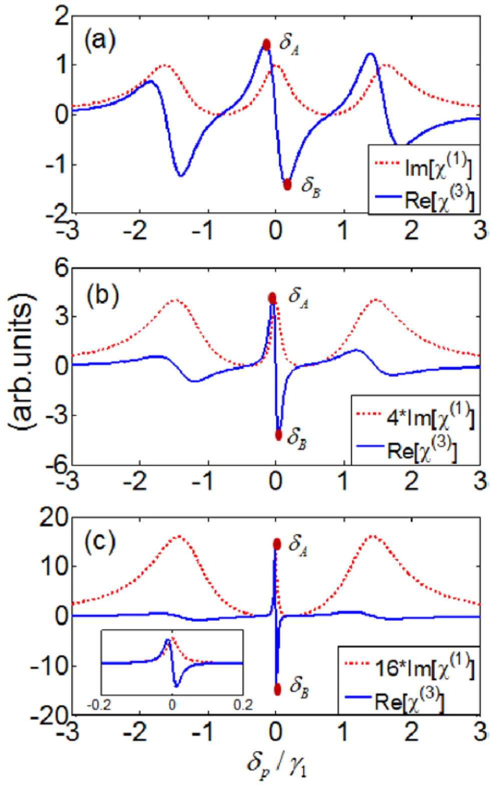

Figure 3. Variation of Re[$\chi^{(3)}$] (solid line) and Im [$\chi^{(1)}$] (dotted line) as a function of the probe detuning δ_p with different energy splitting $-\omega_{12} = \omega_{13} = \omega$, (a) $\omega = 0.8$, (b) $\omega = 0.4$, (c) $\omega = 0.2$. Other parameters are $T_2 = T_3 = 1$, $\gamma_1 = 1$, $\gamma_2 = \gamma_3 = 10^{-3}$.

can tunnel from QD 1 to QD 2 or from QD 1 to QD 3, which creates a four-level tripod structure. In this case, the transparency window becomes wider [solid line in figure 2(a)]. Then we focus on the Kerr nonlinearity, Re[$\chi^{(3)}$], which is shown in figure 2(b). As can be seen, within the transparency window, Re[$\chi^{(3)}$] with T_3 [solid line] is smaller than that of the one without T_3 [dotted line].

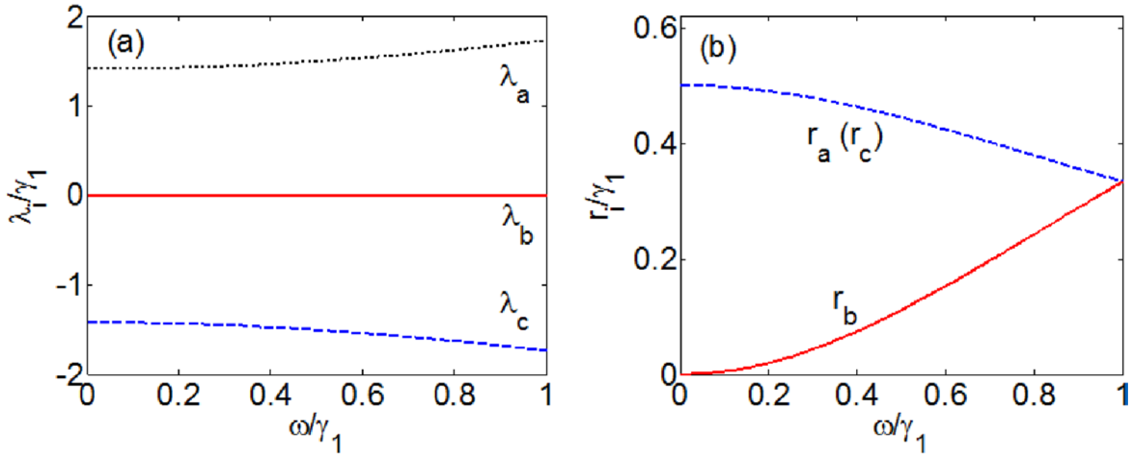

Figure 4. (a) The eigen energies λ_i $(i = a, b, c)$ as a function of the energy splitting ω. (b) The decay rates of the dressed state to the ground state γ_i $(i = a, b, c)$ as a function of the energy splitting ω. The parameters are the same as those in figure 3.

The physical interpretation can be seen clearly under the dressed state picture. To do that, the Hamiltonian corresponding to the system and the tunneling coupling needs to be diagonalized. Then the expressions of the dressed states [figure 1(c)] are:

$$|i\rangle = \cos\theta\cos\varphi|1\rangle + \cos\theta\sin\varphi|2\rangle + \sin\theta|3\rangle,\ (i = a, b, c) \tag{10}$$

where $\tan\varphi = -\frac{T_2}{\omega_{12}+\lambda_i}$ and $\tan\theta = \frac{(\omega_{12}+\lambda_i)T_3}{(\omega_{13}+\lambda_i)\sqrt{T_2^2+(\omega_{12}+\lambda_i)^2}}$. λ_i is the eigenvalues of the dressed level $|i\rangle$ $(i = a, b, c)$, giving the relative energy of the dressed sublevels $|i\rangle$ $(i = a, b, c)$. Therefore, the weak probe field couples the transition from state $|0\rangle$ to the dressed state $|i\rangle$ $(i = a, b, c)$.

Consider $\omega_{12} = \omega_{13} = 0$, and in the limit $T_3 \to 0$, equation (10) is simplified as:

$$|a\rangle = \frac{1}{\sqrt{2}}|1\rangle - \frac{1}{\sqrt{2}}|2\rangle, \tag{11a}$$

$$|b\rangle = |3\rangle \tag{11b}$$

$$|c\rangle = \frac{1}{\sqrt{2}}|1\rangle + \frac{1}{\sqrt{2}}|2\rangle, \tag{11c}$$

with the eigenvalues being $\lambda_a = T_2$, $\lambda_b = 0$ and $\lambda_c = -T_2$. So the dressed level $|b\rangle$ coincides with the bare state $|3\rangle$ and, hence, is decoupled from the system. And the two dressed levels $|a\rangle$ and $|c\rangle$ correspond to the usual Autler-Townes dressed components, and the energy splitting of them is $2T_2$. Because both dressed levels $|a\rangle$ and $|c\rangle$ have a finite overlap with the excited state $|1\rangle$, one dark resonance arises due to quantum interference in two probe transitions $|0\rangle \to |a\rangle$ and $|0\rangle \to |c\rangle$, which suppresses the linear absorption, $\mathrm{Im}[\chi^{(1)}]$ for $\delta_p = 0$.

While in the case of $T_3 \neq 0$, the eigenvalues are $\lambda_a = \sqrt{T_2{}^2 + T_3{}^2}$, $\lambda_b = 0$ and $\lambda_c = -\sqrt{T_2{}^2 + T_3{}^2}$, then equation (10) goes to:

$$|a\rangle = \frac{1}{\sqrt{2}}|1\rangle - \frac{T_2}{\sqrt{2T_2^2+2T_3^2}}|2\rangle + \frac{T_3}{\sqrt{2T_2^2+2T_3^2}}|3\rangle, \tag{12a}$$

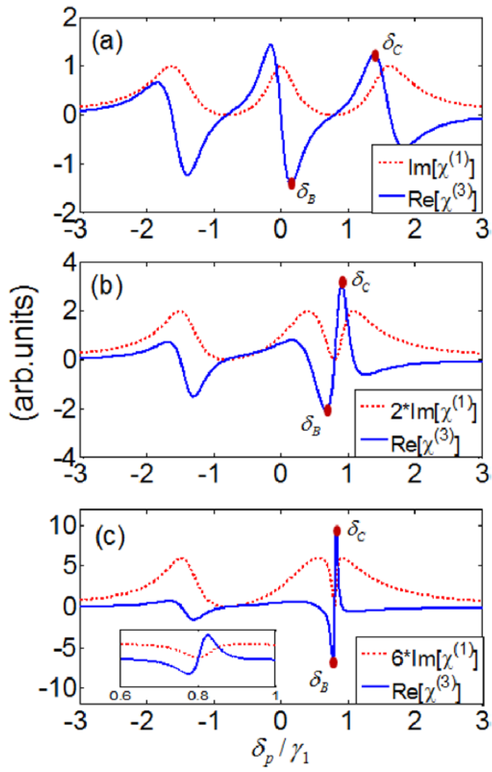

Figure 5. Variation of $\mathrm{Re}[\chi^{(3)}]$ (solid line) and Im $[\chi^{(1)}]$ (dotted line) as a function of the probe detuning δ_p with different tunneling intensity T_3, (a) $T_3 = 1$, (b) $T_3 = 0.4$, (c) $T_3 = 0.2$. Other parameters are $-\omega_{12} = \omega_{13} = 0.8$, $\gamma_1 = 1$, $\gamma_2 = \gamma_3 = 10^{-3}$.

$$|b\rangle = \frac{\sqrt{2T_2^2+T_3^2}}{\sqrt{2T_2^2+2T_3^2}}|2\rangle + \frac{T_3}{\sqrt{2T_2^2+2T_3^2}}|3\rangle, \tag{12b}$$

$$|c\rangle = \frac{1}{\sqrt{2}}|1\rangle + \frac{T_2}{\sqrt{2T_2^2+2T_3^2}}|2\rangle + \frac{T_3}{\sqrt{2T_2^2+2T_3^2}}|3\rangle. \tag{12c}$$

with the eigenvalues being $\lambda_a = \sqrt{T_2^2+T_3^2}$, $\lambda_b = 0$ and $\lambda_c = -\sqrt{T_2^2+T_3^2}$. As can be seen form equation (12b) that

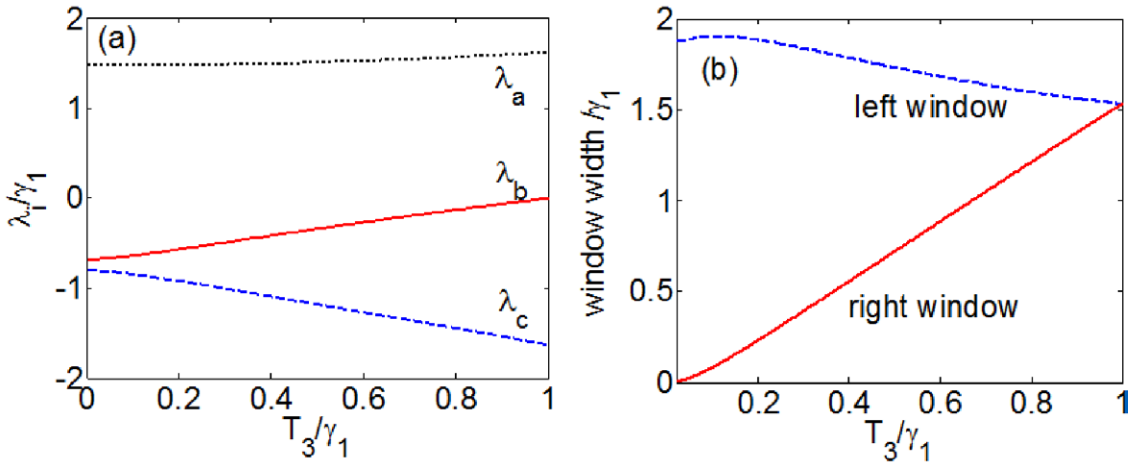

Figure 6. (a) The eigen energies λ_i ($i = a, b, c$) as a function of the tunneling intensity T_3. (b) The FWHM of the two transparency windows as a function of the tunneling intensity T_3. The parameters are the same as those in figure 5.

state $|b\rangle$ does not contain an admixture of $|1\rangle$ and thus has a zero dipole matrix element with ground state $|0\rangle$. There is still one dark resonance resulting from quantum interference in two probe transitions $|0\rangle \to |a\rangle$ and $|0\rangle \to |c\rangle$. But this time, the energy splitting of the two dressed states $|a\rangle$ and $|c\rangle$ is $2\sqrt{T_2^2 + T_3^2}$, resulting in broadening of the transparency window. Within the wider transparency window, the refractive part of the third-order susceptibility is reduced, that is to say, Kerr nonlinearity is reduced.

So, in the following, we will only investigate the case of $\omega_{12} \neq \omega_{13}$. In this situation, all the dressed levels $|i\rangle$ ($i = a, b, c$) [equation (10)] have a finite overlap with the excited state $|1\rangle$ and, thus, have a nonzero dipole matrix element with ground state $|0\rangle$. Therefore, there are two dark resonances resulting from quantum interference between the three dipole-allowed transitions $|0\rangle \to |a\rangle, |0\rangle \to |b\rangle$ and $|0\rangle \to |c\rangle$. The interaction of these two dark resonances will result in the emerging of the center absorption peak and the enhancement of Kerr nonlinearity, which will be discussed in more detail below.

In figure 3 we plot Im$[\chi^{(1)}]$ (dotted line) and Re$[\chi^{(3)}]$ (solid line) as a function of probe detuning for varied values of energy splitting, ($\omega_{12} = -\omega_{13} = \omega$). As can be seen, under two tunneling couplings there is one extra absorption peak showing up in the center of the absorption spectrum. With a decreasing value of ω, the linewidth of the center peak becomes narrowed [dotted line in figure 3]. Simultaneously the Kerr nonlinearity becomes enhanced [solid line in figure 3]. Compared with the results in figure 2, Re$[\chi^{(3)}]$ is greatly enhanced in the vicinity of the center absorption peak that resulted from the interacting double dark resonances.

The above results can be seen clearly in the dressed states picture. From equation (10), there are three dressed states $|i\rangle$ ($i = a, b, c$) and the eigenvalues, λ_i, represent the energy of the dressed states. If the probe field scans over the system, there will be three absorption peaks, and the maximal absorption occurs when the frequency of the probe field is chosen such that it is in resonance with one of the transitions $|0\rangle \leftrightarrow |i\rangle$. Therefore the position of the absorption peaks is determined by the eigenvalues, λ_i. We show in figure 4(a) the eigen energies (λ_i, $i = a, b, c$) as a function of the energy splitting, ω. With a decreasing value of ω, the eigen energy, λ_b, remains zero, and the eigen energies λ_a and λ_b become smaller, which means that the three absorption peaks become closer. On the other hand, the linewidth of these three absorption peaks corresponding to the decay rate of the dressed states, which can be calculated as:

$$\gamma_i = \cos^2\theta\cos^2\varphi|1\rangle + \cos^2\theta\sin^2\varphi|2\rangle + \sin^2\theta|3\rangle. \quad (i = a, b, c) \qquad (13)$$

With equation (13), we plot the decay rate of the three dressed states in figure 4(b). From the figure, one can see that as ω is decreasing, the decay rate of the state $|b\rangle$ is reduced, resulting in the narrowing of the center absorption peak. Therefore, within the narrowing center absorption peak, the Kerr nonlinearity (between detuning δ_A and δ_B) is enhanced.

In figure 3, though the Kerr nonlinearity Re$[\chi^{(3)}]$ is enhanced, it is accompanied by a strong linear absorption, which is not desirable for applications of low-intensity nonlinear optics. Fortunately, one can tune the tunneling intensity of T_3 to obtain the enhanced Kerr nonlinearity with vanishing absorption. In figure 5, we plot Im$[\chi^{(1)}]$ (dotted line) and Re$[\chi^{(3)}]$ (solid line) as a function of probe detuning for various values of tunneling, T_3.As can be seen, the center absorption peak is off center and the absorption spectrum becomes unsymmetrical under the condition of $T_3 \neq T_2$. As T_3 is decreasing, the transparency window on the right side becomes narrowed, and simultaneously, the Kerr nonlinearity, Re$[\chi^{(3)}]$, is enhanced and gradually enters the narrowed transparency window. That is to say, the Kerr nonlinearity is dramatically enhanced with suppressed linear absorption.

The results can also be explained in the dressed states picture. We show the eigen energies (λ_i, $i = a, b, c$) as a function of T_3 in figure 6(a). With a decreasing value of T_3, the energy difference between λ_a and λ_b is increased, while the energy difference between λ_c and λ_b is decreased. Therefore, left and center absorption peaks become farther away, while right and center absorption peaks become closer. The full width at half maximum (FWHM) of the transparency windows as a function of T_3 is shown in figure 6(b). As tunneling, T_3, is decreasing, the linewidth of the left transparency becomes wider while that of the right transparency becomes narrower. Within the narrower transparency window, the Kerr nonlinearity (between detuning δ_B and δ_C) can be enhanced.

From the results obtained in figures 3 and 5, it can be concluded that the Kerr nonlinearity can be controlled by the energy splitting and the tunneling intensity. The physical interpretation is that the interaction of double dark resonances results in the linewidth narrowing of the center absorption peak (between detuning δ_A and δ_B) or the linewidth narrowing of the transparency window (between detuning δ_B and δ_C), where the steep dispersion profile of the probe field makes it possible to enhance the Kerr nonlinearity. And in the later case, giant Kerr nonlinearity with vanishing absorption can be achieved.

4. Conclusions

In this paper, we demonstrate that it is possible to obtain giant enhancement of Kerr nonlinearity via tunneling effects in triangular TQDs. By properly choosing the energy splitting and the tunneling, the enhancement of Kerr nonlinearity can be accompanied by vanishing linear absorption. The results are interpreted in the dressed states picture. Analyses show that the interacting double dark resonances induced by the tunneling between the quantum dots can result in linewidth narrowing of the absorption peak or the transparency window, where the steep dispersion profile of the probe field makes it possible to enhance the Kerr nonlinearity. Potential applications of such semiconductor nanostructures are to enhance self phase modulation at low light levels such as optical solitons and self focusing.

Acknowledgments

This work is supported by the financial support from the National Natural Science Foundation of China (Grant Nos. 11304308 and 61176046), the National Basic Research Program of China (Grant Nos. 2013CB933300), the International Science Technology Cooperation Program of China (No. 2013DFR00730), and the Hundred Talents Program of Chinese Academy of Sciences, and Jilin Provincial Natural Science Foundation (Grant Nos. 20140101203JC).

References

[1] Boyd R W 1992 *Nonlinear Optics* (New York: Academic)
[2] Chi S and Guo Q 1995 *Opt. Lett.* **20** 1598
[3] Tikhonenko V, Christou J and Luther-Davies B 1996 *Phys. Rev. Lett.* **76** 2698
[4] Wu Y 2005 *Phys. Rev.* A **71** 053820
[5] Hang C, Huang G and Deng L 2006 *Phys. Rev.* E **73** 036607
[6] Rebić S, Vitali D, Ottaviani C, Tombesi P, Artoni M, Cataliotti F and Corbalán R 2004 *Phys. Rev.* A **70** 032317
[7] Joshi A and Xiao M 2005 *Phys. Rev.* A **72** 062319
[8] Kou J, Wan R G, Kang Z H, Wang H H, Jiang L, Zhang X J, Jiang Y and Gao J Y 2010 *J. Opt. Soc. Am.* B **27** 2035
[9] Harris S E 1997 *Phys. Today* **50** 36
[10] Harris S E, Field J E and Imamoğlu A 1990 *Phys. Rev. Lett.* **64** 1107
[11] Wang H, Goorskey D and Xiao M 2001 *Phys. Rev. Lett.* **87** 073601
[12] Schmidt H and Imamoglu A 1996 *Opt. Lett.* **21** 1936
[13] Kang H and Zhu Y 2003 *Phys. Rev. Lett.* **91** 093601
[14] Matsko A B, Novikova I, Welch G R and Zubairy M S 2003 *Opt. Lett.* **28** 96
[15] Niu Y, Gong S, Li R, Xu Z and Liang X 2005 *Opt. Lett.* **30** 3371
[16] Hui S, Yueping N, Shiqi J and Shangqing G 2008 *J. Phys. B: At. Mol. Opt. Phys.* **41** 065504
[17] Bai Y, Yang W and Yu X 2010 *Opt. Commun.* **283** 5062
[18] Hemmer P R, Katz D P, Donoghue J, Shahriar M S, Kumar P and Cronin-Golomb M 1995 *Opt. Lett.* **20** 982
[19] Imamoğlu A, Schmidt H, Woods G and Deutsch M 1997 *Phys. Rev. Lett.* **79** 1467
[20] Harris S E and Hau L V 1999 *Phys. Rev. Lett.* **82** 4611
[21] Al-Khursan A H, Al-Khakani M K and Al-Mossawi K H 2009 *Photonics Nanostruct.* **7** 153
[22] Ali K-N, Mohammad S, Mostafa S and Vahid F 2014 *J. Opt.* **16** 055004
[23] Wang L, Rastelli A, Kiravittaya S, Benyoucef M and Schmidt O G 2009 *Adv. Mater.* **21** 2601
[24] Müller K *et al* 2012 *Phys. Rev. Lett.* **108** 197402
[25] Weiss K M, Elzerman J M, Delley Y L, Miguel-Sanchez J and Imamoğlu A 2012 *Phys. Rev. Lett.* **109** 107401
[26] Sköld N, Boyer de la Giroday A, Bennett A J, Farrer I, Ritchie D A and Shields A J 2013 *Phys. Rev. Lett.* **110** 016804
[27] Yuan C H and Zhu K D 2006 *Appl. Phys. Lett.* **89** 052115
[28] Borges H S, Sanz L, Villas-Bôas J M, Diniz Neto O O and Alcalde A M 2012 *Phys. Rev.* B **85** 115425
[29] Lü X Y, Wu J, Zheng L L and Zhan Z M 2011 *Phys. Rev.* A **83** 042302
[30] Cheng M T, Ma X S, Luo Y Q, Wang P Z and Zhao G X 2011 *Appl. Phys. Lett.* **99** 223509
[31] Voutsinas E, Terzis A F and Paspalakis E 2014 C *Phy. Lett.* A **378** 219
[32] Tian S C, Tong C Z, Wang C L, Wang L J, Wu H, Xing E B, Ning Y Q and Wang L J 2014 *Opt. Commun.* **312** 296
[33] Tian S C, Wan R G, Tong C Z, Ning Y Q, Qin L and Liu Y 2014 *J. Phys. B: At. Mol. Opt. Phys.* **47** 155503
[34] Tian S C, Wan R G, Tong C Z and Ning Y Q 2014 *J. Opt. Soc. Am.* B **31** 2681
[35] Peng Y, Yang A, Chen B, Xu Y and Hu X 2014 *J. Opt. Soc. Am.* B **31** 2188
[36] Chen L B and Wen Y 2014 *Laser Phys. Lett.* **11** 105201
[37] Peng Y D, Yang A H, Li D H, Zhang H, Niu Y P and Gong S Q 2014 *Laser Phys. Lett.* **11** 065201
[38] Tian S-C, Wan R-G, Li L-H, Tong C-Z and Ning Y-Q 2015 *Opt. Commun.* **334** 94–100
[39] Lukin M D, Yelin S F, Fleischhauer M and Scully M O 1999 *Phys. Rev.* A **60** 3225
[40] Chen Y C, Liao Y A, Chiu H Y, Su J J and Yu I A 2001 *Phys. Rev.* A **64** 053806
[41] Songmuang R, Kiravittaya S and Schmidt O G 2003 *Appl. Phys. Lett.* **82** 2892
[42] Lee J H, Wang Z M, Strom N W, Mazur Y I and Salamo G J 2006 *Appl. Phys. Lett.* **89** 202101
[43] Chang-Yu H, Yun-Pil S, Marek K and Pawel H 2012 *Rep. Prog. Phys.* **75** 114501
[44] Tian S C, Tong C Z, Wan R G, Ning Y Q and Wang L J 2013 arXiv:1310.4599
[45] Agarwal G S 1974 *Quantum Optics* (Berlin: Springer)
[46] Radmore S M and Radmore P M 1997 *Methods in Theoretical Quantum Optics* (Oxford: Oxford University Press)
[47] Kim J, Chuang S L, Ku P C and Chang-Hasnain C J 2004 *J. Phys.: Condens. Matter* **16** S3727