

Resonance fluorescence from an asymmetric quantum dot dressed by a bichromatic electromagnetic field

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(Received 11 November 2016; published 23 January 2017)

We present the theory of resonance fluorescence from an asymmetric quantum dot driven by a two-component electromagnetic field with two different frequencies, polarizations, and amplitudes (bichromatic field) in the regime of strong light-matter coupling. It follows from the elaborated theory that the broken inversion symmetry of the driven quantum system and the bichromatic structure of the driving field result in unexpected features of the resonance fluorescence, including the infinite set of Mollow triplets, the quench of fluorescence peaks induced by the dressing field, and the oscillating behavior of the fluorescence intensity as a function of the dressing field amplitude. These quantum phenomena are of general physical nature and, therefore, can take place in various double-driven quantum systems with broken inversion symmetry.

DOI: [10.1103/PhysRevA.95.013834](https://doi.org/10.1103/PhysRevA.95.013834)

I. INTRODUCTION

Advances in nanotechnology, laser physics, and microwave techniques have created a basis for studies of the strong light-matter coupling in various quantum systems. Differently from the case of weak electromagnetic field, the interaction between electrons and a strong field cannot be treated as a perturbation. Therefore, the system “electron + strong field” is conventionally considered as a composite electron-field object which is called “electron dressed by field” (dressed electron) [1,2]. The field-induced modification of the physical properties of dressed electrons was studied in both atomic systems [1–3] and various condensed-matter structures, including bulk semiconductors [4–6], graphene [7–11], quantum wells [12–17], quantum rings [18–21], quantum dots [22–32], etc. Among these structures, quantum dots (QDs)—semiconductor three-dimensional (3D) structures of nanometer scale, which are referred to as “artificial atoms”—seem to be the most interesting for optical studies since they are basic elements of modern nanophotonics [33,34]. In contrast to natural atoms, most QDs are devoid of inversion symmetry and, therefore, are asymmetric. As an example, QDs based on gallium nitride heterostructures have a strong built-in electric field [35,36] and, therefore, acquire the giant anisotropy [37–41]. This motivates studies of various asymmetry-induced optical effects in QDs [29–32].

Most studies of dressed quantum systems have been performed for a monochromatic dressing field. However, there are a lot of interesting phenomena specific for quantum systems driven by a two-mode electromagnetic field with two different frequencies, polarizations, and amplitudes (bichromatic field). In symmetric quantum systems (atoms and superconducting qubits), the bichromatic coupling leads to features of photon correlations, squeezing, Autler-Townes effect, suppression of

spontaneous emission, multiphoton transitions, etc. [42–49]. Broken symmetry brings substantially new physics to bichromatically dressed quantum systems, including additional lines in optical spectra, multiple splitting of the dressed-state transitions, etc. [50–52]. Although these optical effects have been extensively studied for a long time, a consistent quantum theory of resonance fluorescence from bichromatically dressed asymmetric systems has not been elaborated. The present research is aimed to fill this gap at the border between quantum optics and physics of nanostructures. To solve the problem, we focused on the strong light-matter coupling regime when the interaction of an asymmetric QD with a bichromatic dressing field overcomes the spontaneous emission and nonradiative decay of QD excitations. In this case, the spectral lines of QD are well resolved and various radiation effects (particularly, the resonance fluorescence) can be analyzed using a concept of quasienergetic (dressed) electronic states. In the framework of this approach, such characteristics of dressed QDs as decay rates and line shapes can be calculated by solving the master equations in the representation of quasienergetic states. As a result, we found unexpected features of resonance fluorescence, which are discussed below.

The paper is organized as follows. In Sec. II, we derive quasienergetic electronic states for an asymmetric QD dressed by a bichromatic electromagnetic field and calculate matrix elements of optical dipole transitions between these states. In Sec. III, we apply the found quasienergetic spectrum of dressed electrons to elaborate the theory of resonance fluorescence from the QD. Section IV contains the discussion of the calculated spectra of resonance fluorescence and the conclusion.

II. MODEL OF ELECTRONIC STRUCTURE

Let us consider an asymmetric QD with broken inversion symmetry along the z axis, which is dressed by the bichromatic

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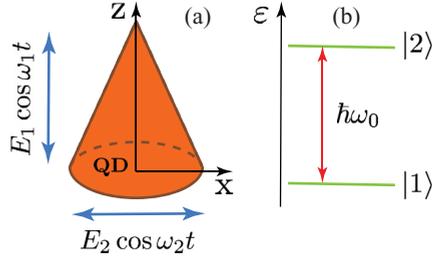


FIG. 1. Sketch of the system under consideration: (a) Asymmetric quantum dot (QD) with broken inversion symmetry along the z axis, which is exposed to a bichromatic electromagnetic field with the electric field amplitudes $E_{1,2}$ and the frequencies $\omega_{1,2}$. (b) Two-level model of the electronic energy spectrum of the QD, ϵ , with the interlevel distance $\hbar\omega_0$.

field

$$\mathbf{E}(t) = \mathbf{E}_1 \cos \omega_1 t + \mathbf{E}_2 \cos \omega_2 t, \quad (1)$$

where the electric field of the first mode, \mathbf{E}_1 , is directed along the z axis and the electric field of the second mode, \mathbf{E}_2 , is perpendicular to this axis [see Fig. 1(a)]. In what follows, we will assume that the second frequency ω_2 is near the electronic resonance frequency ω_0 , whereas the first frequency ω_1 is far from all resonance frequencies. As a consequence, the bichromatic field (1) effectively mixes only two electron states of the QD, $|1\rangle$ and $|2\rangle$, which are separated by the energy $\hbar\omega_0$ [see Fig. 1(b)]. Within the basis of these two states, the asymmetric QD can be described by the matrix Hamiltonian [29]

$$\hat{\mathcal{H}} = \begin{pmatrix} -\hbar\omega_0/2 - d_{11}E_1 \cos \omega_1 t & -d_{12}E_2 \cos \omega_2 t \\ -d_{21}E_2 \cos \omega_2 t & \hbar\omega_0/2 - d_{22}E_1 \cos \omega_1 t \end{pmatrix}, \quad (2)$$

where $d_{11} = \langle 1|ez|1\rangle$, $d_{22} = \langle 2|ez|2\rangle$, and $d_{12} = d_{21} = \langle 1|ex|2\rangle$ are the matrix elements of the operator of the electric dipole moment along the z, x axes, and e is the electron charge. To simplify calculations, the Hamiltonian (2) can be written as a sum, $\hat{\mathcal{H}} = \hat{\mathcal{H}}_0 + \hat{\mathcal{H}}'$, where

$$\hat{\mathcal{H}}_0 = \begin{pmatrix} -\hbar\omega_0/2 - d_{11}E_1 \cos \omega_1 t & 0 \\ 0 & \hbar\omega_0/2 - d_{22}E_1 \cos \omega_1 t \end{pmatrix} \quad (3)$$

is the diagonal part of the full Hamiltonian (2), and

$$\hat{\mathcal{H}}' = \begin{pmatrix} 0 & -d_{12}E_2 \cos \omega_2 t \\ -d_{21}E_2 \cos \omega_2 t & 0 \end{pmatrix} \quad (4)$$

is the nondiagonal part describing electron transitions between states $|1\rangle$ and $|2\rangle$ under influence of the field (1). Exact solutions of the nonstationary Schrödinger equation with the Hamiltonian (3),

$$i\hbar \frac{\partial \psi}{\partial t} = \hat{\mathcal{H}}_0 \psi,$$

can be written in the spinor form as

$$\psi^{(-)} = e^{i\omega_0 t/2} \exp \left[i \frac{d_{11}E_1}{\hbar\omega_1} \sin \omega_1 t \right] \begin{pmatrix} 1 \\ 0 \end{pmatrix} \quad (5)$$

and

$$\psi^{(+)} = e^{-i\omega_0 t/2} \exp \left[i \frac{d_{22}E_1}{\hbar\omega_1} \sin \omega_1 t \right] \begin{pmatrix} 0 \\ 1 \end{pmatrix}. \quad (6)$$

Since the two pseudospinors (5) and (6) are the complete basis of the considered electronic system at any time t , we can seek eigenstates of the full Hamiltonian (2) as an expansion

$$\tilde{\psi} = a^{(-)}(t)\psi^{(-)} + a^{(+)}(t)\psi^{(+)}, \quad (7)$$

where the time-dependent coefficients $a^{(\pm)}(t)$ obey the equation

$$i\hbar \dot{a}^{(\mp)}(t) = -a^{(\pm)}(t) d_{12} E_2 \cos \omega_2 t e^{\mp i\omega_0 t} \times \exp \left[\pm i \frac{(d_{22} - d_{11})E_1}{\hbar\omega_1} \sin \omega_1 t \right]. \quad (8)$$

Applying the Jacobi-Anger expansion,

$$e^{iz \sin \theta} = \sum_{n=-\infty}^{\infty} J_n(z) e^{in\theta},$$

we arrive from Eq. (8) at the equation

$$i\hbar \dot{a}^{(\mp)}(t) = -a^{(\pm)}(t) \frac{d_{12}E_2}{2} \sum_{n=-\infty}^{\infty} J_n \left[\frac{(d_{22} - d_{11})E_1}{\hbar\omega_1} \right] \times [e^{\mp i(\omega_0 + \omega_2 - n\omega_1)t} + e^{\mp i(\omega_0 - \omega_2 - n\omega_1)t}], \quad (9)$$

where $J_n(z)$ is the Bessel function of the first kind. Formally, the equation of quantum dynamics (9) describes a two-level quantum system subjected to a multimode field. It is well known that the main contribution to the solution of such equation arises from a mode which is nearest to the resonance. Correspondingly, near the resonance condition,

$$\omega_0 \pm \omega_2 = n\omega_1, \quad (10)$$

we can neglect all modes except the resonant one. In this approximation, Eq. (9) reads

$$i\dot{a}^{(\pm)}(t) = a^{(\mp)}(t) F_n e^{\pm i\varphi_n t}, \quad (11)$$

where $F_n = -(d_{12}E_2/2\hbar)J_n(\tilde{\omega}/\omega_1)$ are the Rabi frequencies of the considered system,

$$\tilde{\omega} = \frac{E_1(d_{22} - d_{11})}{\hbar} \quad (12)$$

is the effective frequency, and $\varphi_n = \omega_0 \pm \omega_2 - n\omega_1$ is the resonance detuning. It follows from Eq. (11) that the considered problem is reduced to the effective two-level system driven by the monochromatic field with the combined frequency φ_n . Using the well-known solution of Eq. (11) (see, e.g., Ref. [53]), we can write the sought wave functions (7) in the conventional form of quasienergetic (dressed) states as

$$\tilde{\psi}_1 = e^{-i\tilde{\epsilon}_1 t/\hbar} \left[\sqrt{\frac{1}{2} \left(1 + \frac{\varphi_n}{2\Omega_n} \right)} \Lambda_{11}(t) \begin{pmatrix} 1 \\ 0 \end{pmatrix} - \sqrt{\frac{1}{2} \left(1 - \frac{\varphi_n}{2\Omega_n} \right)} e^{i(\varphi_n - \omega_0)t} \Lambda_{22}(t) \begin{pmatrix} 0 \\ 1 \end{pmatrix} \right], \quad (13)$$

$$\begin{aligned} \tilde{\psi}_2 = e^{-i\tilde{\varepsilon}_2 t/\hbar} & \left[\sqrt{\frac{1}{2} \left(1 + \frac{\varphi_n}{2\Omega_n} \right)} \Lambda_{22}(t) \begin{pmatrix} 0 \\ 1 \end{pmatrix} \right. \\ & \left. + \sqrt{\frac{1}{2} \left(1 - \frac{\varphi_n}{2\Omega_n} \right)} e^{-i(\varphi_n - \omega_0)t} \Lambda_{11}(t) \begin{pmatrix} 1 \\ 0 \end{pmatrix} \right], \quad (14) \end{aligned}$$

where

$$\tilde{\varepsilon}_1 = -\frac{\hbar\omega_0}{2} - \hbar\Omega_n + \hbar\frac{\varphi_n}{2}, \quad \tilde{\varepsilon}_2 = \frac{\hbar\omega_0}{2} + \hbar\Omega_n - \frac{\hbar\varphi_n}{2} \quad (15)$$

are the corresponding quasienergies, and

$$\Lambda_{ii}(t) = \exp \left[i \frac{d_{ii} E_1}{\hbar\omega_1} \sin \omega_1 t \right], \quad \Omega_n = \sqrt{\frac{\varphi_n^2}{4} + F_n^2}.$$

As to optical transitions between the dressed states (13) and (14), they can be described by the matrix dipole elements

$$\tilde{d}_{ij}(t) = \langle \tilde{\psi}_i | e x | \tilde{\psi}_j \rangle = d_{ij}^{(+)} + d_{ij}^{(-)}, \quad (16)$$

where $d_{ij}^{(+)} = [d_{ji}^{(-)}]^*$,

$$\begin{aligned} d_{21}^{(+)} &= \frac{d_{21}}{2} \left[1 + \frac{\varphi_n}{2\Omega_n} \right] \Lambda^*(t) e^{i(\omega_0 - \varphi_n + 2\Omega_n)t}, \\ d_{21}^{(-)} &= -\frac{d_{12}}{2} \left[1 - \frac{\varphi_n}{2\Omega_n} \right] \Lambda(t) e^{-i(\omega_0 - \varphi_n - 2\Omega_n)t}, \\ d_{11}^{(+)} &= -d_{22}^{(+)} = -\frac{d_{12}}{2} \frac{F_n}{\Omega_n} \Lambda^*(t) e^{i(\omega_0 - \varphi_n)t}, \quad (17) \end{aligned}$$

and $\Lambda(t) = \Lambda_{22}(t)\Lambda_{11}^*(t)$.

III. RESONANCE FLUORESCENCE

The general theory to describe the resonance fluorescence in the representation of quasienergetic (dressed) states has been elaborated in Refs. [54–57]. Applying this known approach to the considered dressed QD, we have to write the Hamiltonian of interaction of the QD with the radiative field as

$$\hat{H}_{\text{int}} = -\tilde{d}_{ij}(t)[\mathbf{E}^+(t) + \mathbf{E}^-(t)], \quad (18)$$

where

$$\mathbf{E}^\pm(t) = \int d\omega g^\pm(\omega) \mathbf{e}_\omega^\pm a_\omega^\pm e^{\mp i\omega t}$$

is the positive (negative) frequency part of the radiative electric field, $\mathbf{E}(t) = \mathbf{E}^+(t) + \mathbf{E}^-(t)$, the parameters $g^\pm(\omega)$, \mathbf{e}_ω^\pm , a_ω^\pm describe the density, polarization, and amplitude of the corresponding electromagnetic modes, respectively, and $\tilde{d}_{ij}(t)$ are the dipole matrix elements of dressed QD (16) and (17). Within the conventional secular approximation and Markov approximation [1], the equations describing the quantum dynamics of the considered two-level system read

$$\begin{aligned} \frac{d\sigma_{11}(t)}{dt} &= -\Gamma_{11}[\sigma_{11}(t) - \sigma_{11}^S], \\ \frac{d\sigma_{12}(t)}{dt} &= -\Gamma_{12}\sigma_{12}(t), \quad \sigma_{11}(t) + \sigma_{22}(t) = 1, \end{aligned}$$

where $\sigma(t) = \text{Tr}\{\rho(t)\}$ is the reduced density operator which involves tracing over reservoir variables, $\sigma_{\alpha\beta} = \langle \tilde{\psi}_\alpha | \sigma | \tilde{\psi}_\beta \rangle$ are

the matrix elements of the density operator written in the basis of dressed states (13) and (14),

$$\sigma_{11}^S = \frac{w_{21}}{w_{12} + w_{21}}, \quad \sigma_{22}^S = \frac{w_{12}}{w_{12} + w_{21}} \quad (19)$$

are the steady-state populations of the dressed states (13) and (14),

$$\Gamma_{11} = w_{12} + w_{21}, \quad \Gamma_{12} = \frac{\Gamma_{11}}{2} - \text{Re}(M_{11,22} + M_{22,11}) \quad (20)$$

are the field-dependent decay rates for the dressed states (13) and (14),

$$w_{12} = \frac{\gamma}{4} \left(1 - \frac{\varphi_n}{2\Omega_n} \right)^2, \quad w_{21} = \frac{\gamma}{4} \left(1 + \frac{\varphi_n}{2\Omega_n} \right)^2 \quad (21)$$

are the probabilities of radiative transitions per unit time between the dressed states (13) and (14),

$$M_{\alpha,\alpha,\beta,\beta} = \int_0^\infty d\tau \tilde{d}_{\alpha\alpha}(t - \tau) \tilde{d}_{\beta\beta}(t) \langle E(t)E(t - \tau) \rangle, \quad (22)$$

where $\langle E(t)E(t - \tau) \rangle$ is the correlation function averaged over the initial state of electromagnetic field, and γ is the spontaneous emission rate. Substituting Eqs. (21) and (22) into Eqs. (19) and (20), we arrive at the width of the transitions,

$$\Gamma_{11} = \frac{\gamma}{2} \left(1 + \frac{\varphi_n^2}{4\Omega_n^2} \right), \quad \Gamma_{12} = \frac{\gamma}{4} \left(3 - \frac{\varphi_n^2}{4\Omega_n^2} \right), \quad (23)$$

and the difference between the populations of dressed electronic states (13) and (14),

$$\Delta_S = \sigma_{11}^S - \sigma_{22}^S = 2w_{21}/\Gamma_{11} - 1. \quad (24)$$

Taking into account the aforesaid, the spectrum of resonance fluorescence from the QD has the form [2]

$$S(\omega) \sim \frac{1}{\pi} \text{Re} \left\{ \int_0^\infty d\tau \langle D^{(+)}(t + \tau) D^{(-)}(t) \rangle e^{-i\omega\tau} \right\}, \quad (25)$$

where

$$D^{(\pm)}(t) = \sum_{\alpha,\beta} \sigma_{\alpha\beta}(t) d_{\alpha\beta}^{(\pm)}(t) \quad (26)$$

is the positive(negative)-frequency part of the polarization operator written in the basis of dressed states (13) and (14). Applying the quantum regression theorem [2] and taking into account Eqs. (16) and (17), we arrive at the correlation function

of the polarization operator (26) in the steady-state regime for long time t and arbitrary time τ ,

$$\begin{aligned} \langle D^{(+)}(t + \tau)D^{(-)}(t) \rangle &= [\Delta_S^2 + (1 - \Delta_S^2)e^{-\Gamma_{11}\tau}] \langle d_{11}^{(+)}(t + \tau)d_{11}^{(-)}(t) \rangle + \frac{1}{2}[(1 + \Delta_S)\langle d_{12}^{(+)}(t + \tau)d_{21}^{(-)}(t) \rangle \\ &\quad + (1 - \Delta_S)\langle d_{21}^{(+)}(t + \tau)d_{12}^{(-)}(t) \rangle] e^{-\Gamma_{12}\tau}, \end{aligned} \quad (27)$$

where

$$\begin{aligned} \langle d_{11}^{(+)}(t + \tau)d_{11}^{(-)}(t) \rangle &= \frac{d_{12}^2 F_n^2}{4 \Omega_n^2} \langle \Lambda^*(t + \tau)\Lambda(t) \rangle e^{i(\omega_0 - \varphi_n)\tau} \approx \frac{d_{12}^2 F_n^2}{4 \Omega_n^2} \Sigma(\tau) e^{-i\varphi_n\tau}, \\ \langle d_{12}^{(+)}(t + \tau)d_{21}^{(-)}(t) \rangle &= \frac{d_{12}^2}{4} \left(1 - \frac{\varphi_n}{\Omega_n}\right)^2 \langle \Lambda^*(t + \tau)\Lambda(t) \rangle e^{i(\omega_0 - \varphi_n - 2\Omega_n)\tau} \approx \frac{d_{12}^2}{4} \left(1 - \frac{\varphi_n}{\Omega_n}\right)^2 \Sigma(\tau) e^{-i(\varphi_n + 2\Omega_n)\tau}, \\ \langle d_{21}^{(+)}(t + \tau)d_{12}^{(-)}(t) \rangle &= \frac{d_{12}^2}{4} \left(1 + \frac{\varphi_n}{\Omega_n}\right)^2 \langle \Lambda^*(t + \tau)\Lambda(t) \rangle e^{i(\omega_0 - \varphi_n + 2\Omega_n)\tau} \approx \frac{d_{12}^2}{4} \left(1 + \frac{\varphi_n}{\Omega_n}\right)^2 \Sigma(\tau) e^{-i(\varphi_n - 2\Omega_n)\tau}, \end{aligned}$$

and $\Sigma(\tau) = \sum_m J_m^2(\tilde{\omega}/\omega_1) e^{i(\omega_0 - m\omega_1)\tau}$. Substituting Eq. (27) into Eq. (25), the spectrum of resonance fluorescence, $S(\omega) = S_1(\omega) + S_2(\omega)$, can be calculated as a sum of the two parts corresponding to the elastic and inelastic scattering of light [2], where the elastic term reads

$$S_1(\omega) \sim \frac{\Delta_S^2}{\pi} \text{Re} \left\{ \int_0^\infty d\tau d_{11}^{(+)}(t + \tau)d_{11}^{(-)}(t) e^{-i\omega\tau} \right\} = \left[\frac{\Delta_S d_{12} F_n}{2\Omega_n} \right]^2 \sum_m J_m^2\left(\frac{\tilde{\omega}}{\omega_1}\right) \delta(\omega - [n - m]\omega_1 - \omega_2), \quad (28)$$

and the inelastic term is

$$\begin{aligned} S_2(\omega) \sim \frac{d_{12}^2}{4\pi} \left\{ (1 - \Delta_S^2) \left(\frac{d_{12} E_2}{2\hbar\Omega_n}\right)^2 J_n^2\left(\frac{\tilde{\omega}}{\omega_1}\right) \sum_m \frac{J_m^2(\tilde{\omega}/\omega_1)\Gamma_{11}}{[\omega - (n - m)\omega_1 - \omega_2]^2 + \Gamma_{11}^2} + \frac{1}{2} \frac{(1 - \varphi_n^2/4\Omega_n^2)^2}{(1 + \varphi_n^2/4\Omega_n^2)} \right. \\ \left. \times \sum_m \frac{J_m^2(\tilde{\omega}/\omega_1)\Gamma_{12}}{[\omega - (n - m)\omega_1 - \omega_2 + 2\Omega_n]^2 + \Gamma_{12}^2} + \frac{1}{2} \frac{(1 - \varphi_n^2/4\Omega_n^2)^2}{(1 + \varphi_n^2/4\Omega_n^2)} \sum_m \frac{J_m^2(\tilde{\omega}/\omega_1)\Gamma_{12}}{[\omega - (n - m)\omega_1 - \omega_2 - 2\Omega_n]^2 + \Gamma_{12}^2} \right\}. \end{aligned} \quad (29)$$

In what follows, we will focus on the inelastic term (29) which is responsible for the spectral features of the resonance fluorescence.

IV. DISCUSSION AND CONCLUSION

Let us consider the effect of the two key factors of the considered system—the bichromatic structure of the dressing field and the asymmetry of the QD—on the resonance fluorescence. Mathematically, these two factors can be described by the effective frequency (12) appearing in various terms of Eq. (29). In order to explain it, we have to keep in mind that the ground and excited states of the asymmetric QD, $|1\rangle$ and $|2\rangle$, do not possess a certain spatial parity along the asymmetry axis z . Therefore, the diagonal matrix elements of the dipole moment operator in an asymmetric QD prove to be nonequivalent, $d_{22} \neq d_{11}$. As a consequence, the difference of diagonal dipole matrix elements, $d_{22} - d_{11}$, describes the asymmetry of the QD [29]. Since the effective frequency (12) is the product of this difference and the first field amplitude, E_1 , it can be considered as a quantitative measure of both the asymmetry of the QD and the bichromatic nature of the dressing field (1). In the following, we will discuss the dependence of the resonance fluorescence on this effective frequency, $\tilde{\omega} = (d_{22} - d_{11})E_1/\hbar$.

The calculated spectra of resonance fluorescence (29) are plotted in Figs. 2 and 3 for different effective frequencies $\tilde{\omega}$ near the resonance (10) with $n = 1$. If the QD is symmetric or the dressing field is monochromatic ($\tilde{\omega} = 0$), the terms

with $m \neq 0$ in Eq. (29) vanish since the Bessel functions of the first kind, $J_m(\tilde{\omega}/\omega_1)$, satisfy the condition $J_m(0) = \delta_{m0}$. The nonzero terms with $m = 0$ correspond physically to the well-known Mollow triplet in the fluorescence spectrum of a two-level system driven by a monochromatic field [2], which is plotted in Fig. 2(a). If the QD is asymmetric and the dressing field is bichromatic ($\tilde{\omega} \neq 0$), the nonzero terms with $m \neq 0$ in Eq. (29) result in the infinite set of Mollow triplets which can be numerated by the index $m = 0, \pm 1, \pm 2, \dots$ [see Fig. 2(b) and Figs. 3(a)–3(c)]. It is shown in Figs. 2 and 3 that the bichromatic dressing field generates side Mollow triplets ($m \neq 0$), shifts the main Mollow triplet ($m = 0$), and change the amplitudes of the Mollow triplets. According to Eq. (29), the amplitude of the m th Mollow triplet is proportional to the squared Bessel function, $J_m^2(\tilde{\omega}/\omega_1)$. This leads to the oscillating dependence of fluorescence peaks on the irradiation intensity, $I_1 = \varepsilon_0 E_1^2 c/4$ [see Fig. 3(d)]. It should be stressed that the zeros of the Bessel function, $J_m(\tilde{\omega}/\omega_1)$, correspond physically to the zero amplitude of the m th Mollow triplet. Thus, the dressing field can quench fluorescence peaks. Particularly, the absence of the main Mollow triplet ($m = 0$) in Fig. 3(c) is caused by the first zero of the Bessel function, $J_0(\tilde{\omega}/\omega_1)$.

Summarizing the aforesaid, we can conclude that the exciting of asymmetric quantum systems by a bichromatic field results in the following features of resonance fluorescence spectra: an infinite set of Mollow triplets, the quench of fluorescence peaks induced by the dressing field, and the

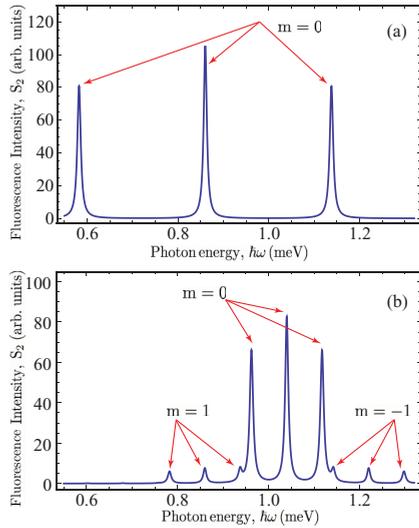


FIG. 2. The spectra of resonance fluorescence from an asymmetric quantum dot with the interlevel distance $\hbar\omega_0 = 1$ meV, dipole moments $d_{22} - d_{11} = d_{12} = 40$ D, and the decay rate $\gamma = 10^{-10}$ s $^{-1}$: (a) in the presence of a monochromatic dressing field with the amplitude $E_2 = 3 \times 10^3$ V/cm and the frequency $\omega_2 = 0.86 \omega_0$; (b) in the presence of a bichromatic dressing field with the amplitudes $E_1 = 1.3 \times 10^3$ V/cm, $E_2 = 3 \times 10^3$ V/cm and the frequencies $\omega_1 = 0.18 \omega_0$, $\omega_2 = 0.86 \omega_0$. The indices $m = 0, \pm 1$ indicate the different Mollow triplets.

oscillating behavior of the fluorescence intensity as a function of the dressing field amplitude. To explain the physics of these effects, we have to stress that the considered bichromatic

field (1) consists of an off-resonant dressing field \mathbf{E}_1 (which renormalizes electronic energy spectrum) and a near-resonant field \mathbf{E}_2 (which induces electron transitions between the two electronic levels). Although the dressing field \mathbf{E}_1 is off resonant, it is very strong. Therefore, there are noticeable multiphoton processes which can involve many photons of the field. Particularly, electron transitions between the electronic levels can be accompanied by absorption of both the near-resonant photon $\hbar\omega_2$ and many off-resonant photons $n\hbar\omega_1$, where n is the number of the photons. As a consequence, there is an infinite set of resonances (10) corresponding to the different numbers $n = 0, 1, 2, \dots$. Since each resonance is accompanied by its own Mollow triplet, an infinite set of Mollow triplets appears in the fluorescence spectrum. As to the field-induced quench of fluorescence peaks and the oscillating behavior of the fluorescence intensity as a function of the dressing field amplitude, these effects arise from the Bessel-function factors in Eq. (29). Physically, these factors describe the nonlinear renormalization of electronic properties with the strong dressing field \mathbf{E}_1 . It should be noted that the appearance of the Bessel functions in expressions describing dressed electrons is a characteristic feature of various quantum systems driven by a dressing field. Particularly, the similar Bessel-function factors describe renormalized electronic properties of dressed quantum wells [58,59] and graphene [60,61].

Since the considered quantum phenomena depend on electronic parameters, the elaborated theory paves the way to the nondestructive optical testing of various asymmetric structures. Applying the developed theory to experimental studies of asymmetric QDs, one should take into account that phonons

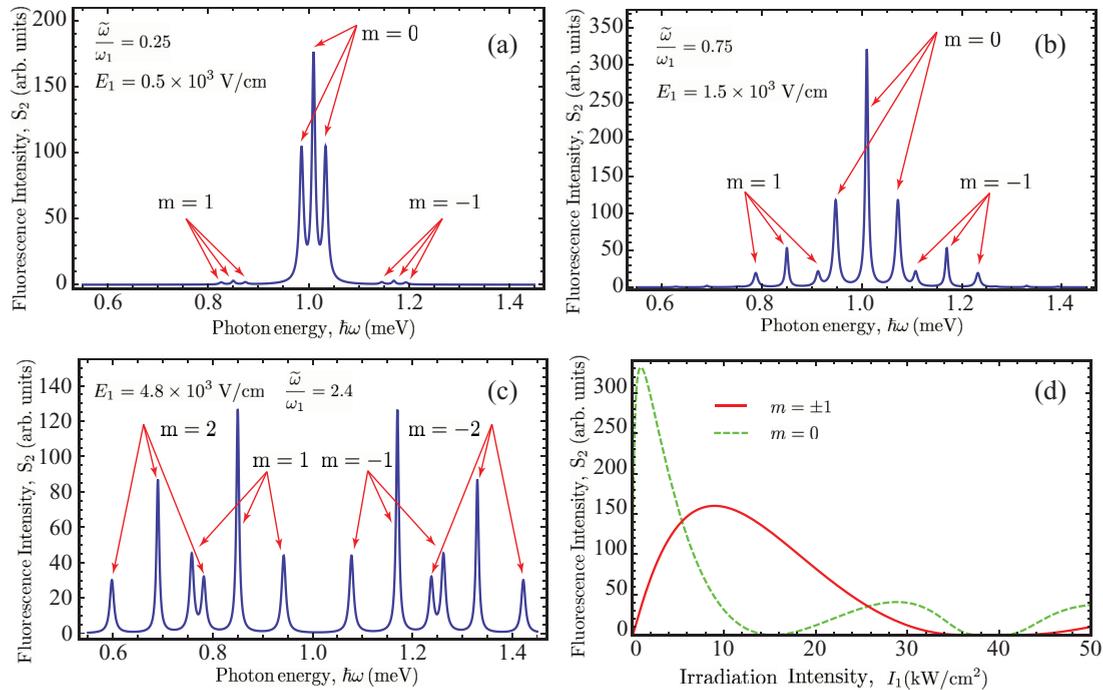


FIG. 3. The spectra of resonance fluorescence from an asymmetric quantum dot with the interlevel distance $\hbar\omega_0 = 1$ meV, dipole moments $d_{22} - d_{11} = d_{12} = 40$ D, and the decay rate $\gamma = 10^{-10}$ s $^{-1}$ in the presence of a bichromatic dressing field with the frequencies $\omega_1 = 0.16 \omega_0$ and $\omega_2 = 0.85 \omega_0$, the amplitude $E_2 = 2.2 \times 10^3$ V/cm, and different amplitudes E_1 : (a)–(c) structure of the Mollow triplets with the numbers $m = 0, \pm 1, \pm 2$; (d) dependence of the central resonant peaks of the Mollow triplets with the numbers $m = 0, \pm 1$ on the irradiation intensity $I_1 = \epsilon_0 E_1^2 c / 4$.

strongly affect optical transitions in semiconductor structures. To avoid the phonon-induced destruction of the discussed fine structure of the fluorescence spectra, measurements should be performed at low temperatures T , satisfying the condition $T \ll \Delta\varepsilon$, where $\Delta\varepsilon = \hbar|F_1|$ is the characteristic field-induced shift of electron energies (the dynamic Stark shift).

ACKNOWLEDGMENTS

The work was partially supported by the RISE Project No. CoExAN, FP7 ITN Project No. NOTEDEV, RFBR

Project No. 17-02-00053, the Rannis Projects No. 141241-051 and No. 163082-051, and the Russian Ministry of Education and Science (Project No. 3.4573.2017). G.Y.K. acknowledges the support of the Armenian State Committee of Science (Project No. 15T-1C052) and thanks the University of Iceland for hospitality. O.V.K. acknowledges the support from the Singaporean Ministry of Education under AcRF Tier 2 Grant No. MOE2015-T2-1-055. I.A.S. acknowledges the support of the Russian Government Programm 5-100.

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