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Photonic electromagnetically induced transparency and collapse of superradiant modes in Bragg multiple quantum wells

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We show that infrared dressing of the superradiant excitons in a multiple-quantum-well photonic band gap structure can lead to photonic electromagnetically induced transparency. This process, which occurs via coherent control of the contributions of resonant excitonic features and background refractive index contrast of the barriers and wells, allows the probe field to pass through the structure in the vicinity of the Bragg wavelength without being affected by the index or absorption periodicity. We also show that the exciton infrared dressing can spatially shift the light wave pattern inside the crystal and destroy the superradiant mode of the structure. When the latter happens, the structure acts as a nonresonant passive photonic band gap, abandoning its excitonic features.

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It is well known that in multiple-quantum-well (MQW) structures the influence of light-induced interwell coupling can be dramatic. In such structures the large spacing between the quantum wells (QWs) eliminates the direct Coulomb interaction between different QWs while their excitonic resonances are radiatively coupled.[1,2] This process generates collective excitation of the MQW structure, causing significant modification of its linear and nonlinear optical responses. In particular, when the interwell separation between the QWs is chosen to be half of the optical wavelength (Bragg structure), the strengths of all oscillators are concentrated in one superradiant mode, suppressing contributions of all subradiant modes.[2,3] Under these conditions, the radiative linewidth of the MQW structure can become much broader than those of single QWs, giving rise a resonant MQW photonic band gap (PBG).

The two main features that determine the optical responses of such PBGs are (i) the nonresonant (background) refractive index contrast of the wells and barriers, and (ii) the resonant excitonic features. In this paper we show how to use infrared (ir) dressing of the superradiant excitons in such structures to coherently control the interplay between these features. A prominent effect of this process is formation of a *photonic* electromagnetically induced transparency (PEIT). In such a process the transparency window occurs at the vicinity of the Bragg wavelength wherein the index and absorption periodicities associated with the MQW structure are coherently concealed from the probe field. In other words, at the center of this window the structure acts as a homogeneous semiconductor medium with some absorption while at other frequencies, within the original PBG, it shows strong signs of dielectric periodicity. The ir dressing can also lead to some other pertinent effects, i.e., collapse of superradiant mode and spatial shift of the light wave pattern inside the MQW structure. The former converts the structure into a nonresonant passive PBG in the absence of any excitonic feature. Such a PBG is generated by the background refractive index contrast of the barriers and wells and, therefore, is immune to the ir laser. The spatial shift of the light wave pattern allows us to control the position of the optical nodes inside the crystal, promising useful applications such as optical control of exciton decay.

Note that the ir laser field responsible for the dressing process is considered to be near resonant with the transitions between the superradiant excitons ($e1$-hh1) and the broad nonradiative ones ($e2$-hh1) associated with the second conduction subbands (Fig. 1). However, despite the apparent similarity of such a system with the Ξ (ladder) systems used to study EIT in atoms,[4] QWs,[5] or exciton-polariton stop bands,[6] the PEIT is not caused by the destructive quantum interference. Instead, it is generated by coherent adjustment of the contributions of the resonant and nonresonant features via optical Stark redshifting and broadening of the $e1$-hh1 excitons.

We are interested in the propagation of an electromagnetic field along the growth direction (z) of a Bragg MQW structure dressed by an ir laser field (control field). For this reason we consider each period of this structure to consist of a GaAs/$Al_{0.2}Ga_{0.8}As$ QW supporting two conduction subbands $e1$ and $e2$ (Fig. 1). In such a system we need to consider the mismatch between the background refractive indices of the barriers (n_b) and wells (n_w). Such a mismatch can form a nonresonant passive PBG by itself when the number of QWs (N) is sufficient. The control field propagates perpendicular to the probe field and is polarized along z. Since here we are only interested in the linear response of the system, propagation of the probe field along z can be described by[7,8]

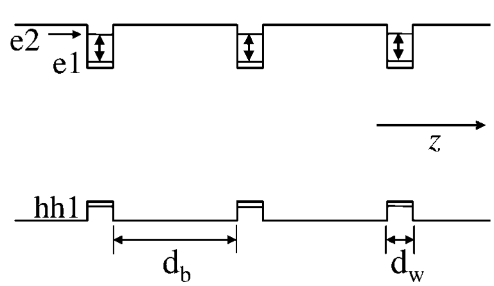

FIG. 1. Schematic diagram of the Bragg MQW structure interacting with an intense ir laser field. $e1$ and $e2$ refer to the conduction subbands and hh1 represents the heavy-hole subband. The two-sided arrows refer to the mixing processes caused by the ir laser.

$$\left[\frac{d^2}{dz^2}+\left(\frac{\omega_s}{c}\right)^2\varepsilon_b(z)\right]\varepsilon_0 E^{\omega_c,I_c}(z)=-\left(\frac{\omega_s}{c}\right)^2 P_{\rm exc}^{\omega_c,I_c}(z). \quad (1)$$

Here $\varepsilon_b = n_{\rm b}^2$ or $n_{\rm w}^2$, ω_s represents the interband probe or signal frequency, and ε_0 is the vacuum permittivity. $P_{\rm exc}^{\omega_c,I_c}$ is related to the exciton susceptibility $\chi_{\rm exc}^{\omega_c,I_c}(\omega_s)$. For given values of the control field intensity (I_c) and frequency (ω_c) it is given by

$$P_{\rm exc}^{\omega_c,I_c}=\chi_{\rm exc}^{\omega_c,I_c}(\omega_s)\varepsilon_0\int \Phi_{1s}^{\rm exc}(z)\Phi_{1s}^{\rm exc}(z')E^{\omega_c,I_c}(z')dz'. \quad (2)$$

Here $\Phi_{1s}^{\rm exc}(z)$ refer to the envelope functions of the $1s$ states of the $e1$-hh1 excitons.

To find out how the ir field influences the collective response of the MQW structure we adopt a transfer matrix method.[7] Based on this method propagation of the electromagnetic field along z through one period of the structure can be described by

$$T_{\omega_c,I_c}=T_{\rm b}^{1/2}T_{\rm bw}T_{\rm w}^{\omega_c,I_c}T_{\rm wb}T_{\rm b}^{1/2}. \quad (3)$$

Here the transfer matrix through one-half of the barrier is given by

$$T_{\rm b}^{1/2}=\begin{pmatrix} e^{i\varphi_b/2} & 0 \\ 0 & e^{-i\varphi_b/2}\end{pmatrix}, \quad (4)$$

where $\varphi_{\rm b}=\omega_s n_{\rm b} d_{\rm b}/c$ and $d_{\rm b}$ refers to the width of the barriers. $T_{\rm bw}=T_{\rm wb}^{-1}$, which describes scattering of the light by the contrast between $n_{\rm w}$ and $n_{\rm b}$, is given by

$$T_{\rm bw}=\frac{1}{1+\rho}\begin{pmatrix} 1 & \rho \\ \rho & 1\end{pmatrix}. \quad (5)$$

Here ρ is the Fresnel reflection coefficient. $T_{\rm w}^{\omega_c,I_c}$ is the transfer matrix associated with the QW region obtained from

$$T_{\rm w}^{\omega_c,I_c}=\begin{pmatrix} e^{i\varphi_w}(1-iS^{\omega_c,I_c}) & -iS^{\omega_c,I_c} \\ iS^{\omega_c,I_c} & e^{-i\varphi_w}(1+iS^{\omega_c,I_c})\end{pmatrix}. \quad (6)$$

Here $\varphi_{\rm w}=\omega_s n_{\rm w} d_{\rm w}/c$ and $d_{\rm w}$ refers to the width of the wells. $S^{\omega_c,I_c}=(\Gamma_0/\alpha_0)\chi_{\rm exc}^{\omega_c,I_c}$ wherein Γ_0 is the radiative decay rate of the $1s$ $e1$-hh1 exciton and $\alpha_0=\varepsilon_\infty\omega_{\rm LT}a_B^3\omega_{1s}^2/4c^2$. Here $\omega_{\rm LT}$ refers to the exciton longitudinal transverse splitting, a_B to bulk Bohr exciton radius, and ω_{1s} to the frequency of the $1s$ $e1$-hh1 excitons. Based on Ref. 3 for the GaAs/$Al_xGa_{1-x}As$ structure studied in this paper we can consider Γ_0 to be around 67 μeV.

To study the effect of the control field on $\chi_{\rm exc}^{\omega_c,I_c}$, we first consider the Hamiltonian of a single QW in the MQW structure when $I_c=0$, i.e., $H=H_0+V$. Here H_0 contains the single-particle energies of the photoexcited electrons and holes, and V describes their Coulomb interaction:

$$V=\sum_{i,j}\sum_{\mathbf{k},\mathbf{k}',\mathbf{q}\neq 0}V_{i,j}(\mathbf{q})a_{i,\mathbf{k}+\mathbf{q}}^\dagger a_{j,\mathbf{k}'-\mathbf{q}}^\dagger a_{j,\mathbf{k}'}a_{i,\mathbf{k}}. \quad (7)$$

In this equation $a_{i,\mathbf{k}}^\dagger$ ($a_{i,\mathbf{k}}$) are the creation (annihilation) operators for electrons in the conduction subband with index i, and $a_{j,\mathbf{k}}^\dagger$ ($a_{j,\mathbf{k}}$) are the corresponding operators for electrons in the valence subbands with index j. Note that, since we consider the probe field to be weak, the electron-electron and hole-hole interactions are ignored. In addition, since here we are dealing with near-resonant linear detection of the $e1$-hh1 excitons we have $i=e1$ and $e2$ and j=hh1. Coulomb mixing between different hole subbands is also considered negligible, as we focus on strain-free QWs with well width below 15 nm.[9]

Let us now consider interaction of such a system with the control field $[E_c(t)=E_c e^{i\omega_c t}]$. Here ω_c is taken to be near resonant with the $e1$-$e2$ transition. The equation of motion for the density matrix of the system ($\rho^{\mathbf{k}}$) can then be obtained from

$$\frac{\partial\rho^{\mathbf{k}}}{\partial t}=-\frac{i}{\hbar}[H+H_{\rm int},\rho^{\mathbf{k}}]+\left.\frac{\partial\rho^{\mathbf{k}}}{\partial t}\right|_{\rm relax}. \quad (8)$$

Here $H_{\rm int}$ refers to the interaction term of the control field with the single QWs. The second term in Eq. (8) represents the corresponding damping terms. The optical Bloch equations can then be transformed into an excitonic basis, acquiring the following form:[10]

$$\frac{d\mathbf{\Lambda}^{ij,\beta}}{dt}=\mathbf{L}^{ij,\beta}\mathbf{\Lambda}^{ij,\beta}+\mathbf{K}^{ij,\beta}. \quad (9)$$

Here the matrix $\mathbf{\Lambda}^{ij,\beta}$ contains the density matrix elements in the excitonic basis, while $\mathbf{L}^{ij,\beta}$ and $\mathbf{K}^{ij,\beta}$ consist of their coefficients. β is the energy index.

To calculate the interband susceptibility associated with the $e1$-hh1 excitons, we employ linear response theory.[11] For $\beta=1s$ we have

$$\chi_{\rm exc}^{\omega_c,I_c}(\omega_s)=\hat{A}\,\langle[P_{1s}^+(t'),\widetilde{P}_{1s}^-(t',z_1)]\rangle|_{z_1=i\omega_s}. \quad (10)$$

Here $P_{1s}^+(t)$ and $P_{1s}^-(t)$ are the positive and negative components of the system polarization and $\hat{A}$ is a normalization coefficient. $\widetilde{P}_{1s}^-(t',z)$ is the Laplace transform of $P_{1s}^-(t)=P_{1s}^-(t'+\tau)$ with respect to $\tau=t-t'$, where $\tau>0$. The two-time correlation functions in Eq. (10) are evaluated using the quantum regression theorem. To do this the density matrix elements are obtained from Eq. (9), assuming the ir field amplitude is slowly varying compared to the dephasing rates of the intersubband excitonic transitions. Having this in mind, we found

$$\langle[P_{1s}^+(\infty),\widetilde{P}_{1s}^-(\infty,z_1)]\rangle=|\mu_{ij}^{1s}|^2\{R_{8,3}^{ij,1s}(z_1)\Lambda_2^{ij,1s}(\infty)+R_{8,8}^{ij,1s}(z_1)\times[\Lambda_1^{ij,1s}(\infty)-\Lambda_7^{ij,1s}(\infty)]\}. \quad (11)$$

Here $i=e1$ and the 9×9 matrix $\mathbf{R}^{ij,1s}(z_1)$ is defined by

$$\mathbf{R}^{ij,1s}(z_1)=(z_1\mathbf{I}-\mathbf{L}^{ij,1s})^{-1}, \quad (12)$$

where $\mathbf{I}$ refers to the identity matrix and μ_{ij}^{1s} to the dipole moments associated with the $1s$ $e1$-hh1 exciton excitation.

To carry out the numerical calculations we assume that the well and barrier widths are 7 and 107.9 nm, respectively. In addition, at low temperature (around 10 K) one can consider $n_{\rm w}=3.6$ and $n_{\rm b}=3.46$.[3] The polarization dephasing rate associated with the nonradiative decay of $e1$-hh1 is also considered ~0.44 meV and that of $e2$-hh1 is ~4.1 meV. The latter is mostly caused by scattering of the excitons with the

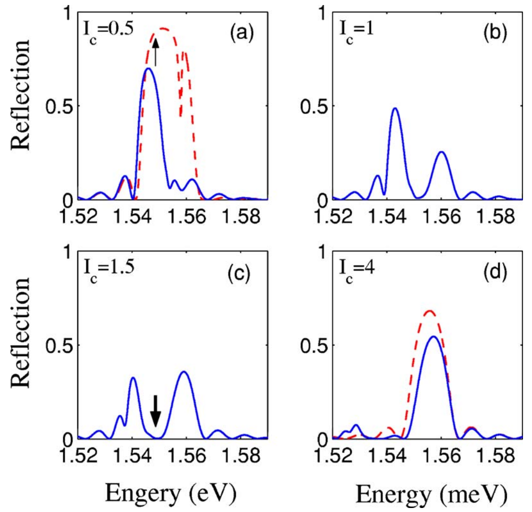

FIG. 2. (Color online) Reflection of the MQW system ($N=150$) for various intensities of the control field (I_c in MW/cm^2). The dashed line in (a) refers to the system's response when $I_c=0$ and in (d) to the nonresonance PBG response of the structure in the absence of the excitons.

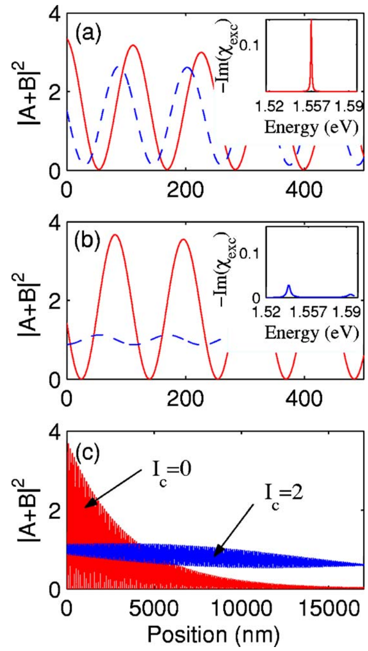

FIG. 3. (Color online) Total optical field inside the MQW structure for 1.557 (a) and 1.549 eV (b) and (c) when $I_c=0$ (solid line) and 2 MW/cm^2 (dashed line). Insets of (a) and (b) show variation of imaginary parts of the susceptibility with I_c. A and B refer to the backward and forward waves.

optical phonons. Figure 2(a) shows reflection of such a structure when $N=150$ and $I_c=0$ (dashed line). Such a reflection is the result of dominance of the superradiant mode. The dip occurs around the energies of the 1s states of these excitons, considered to be 1.557 eV. Assuming $\Delta_c=\hbar(\omega_c-\omega_{12})=-20$ meV ($\hbar\omega_{12}$ is the $e1$-$e2$ transition energy) and $I_c=0.5$ MW/cm^2, the MQW reflection mostly becomes narrow and redshifted [Fig. 2(a)]. For $I_c=1$ MW/cm^2, however, the amplitude of the redshifted peak is suppressed while another peak at higher energy appears, forming a doublet [Fig. 2(b)]. At higher control field intensity ($I_c=1.5$ MW/cm^2) the amplitude of the new peak becomes similar to that of the other peak [Fig. 2(c)]. The reflection intensity at the middle of this doublet is nearly zero (thick arrow), suggesting a switching off process with nearly 100% extension ratio. As I_c increases further the lower energy peak disappears while the amplitude of the other peak reaches a saturated value. For $I_c=4$ MW/cm^2 the reflection nearly freezes [Fig. 2(d), solid line] and becomes similar to that of the nonresonant PBG response of the structure purely generated by the difference between n_b and n_w in the absence of any excitonic effect [Fig. 2(d), dashed line].

To find the corresponding effects on the light propagation inside the MQW structure, let us consider the wave pattern of the probe field along z. As expected, when the photon energy of this field is 1.557 eV and $I_c=0$, the minima (nodes) are located at the QW positions [Fig. 3(a), solid line]. This can inhibit emission of the 1s e1-hh1 excitons, to a large extent, if the structure is weakly pumped slightly above its electronic band gap.[2] As shown in Fig. 3(a) (dashed line), however, for $I_c=2$ MW/cm^2 and $\Delta_c=-20$ meV the pattern is optically shifted in space, placing the minima in the barrier regions. At 1.549 eV, where the minimum of the doublet in Fig. 2(c) occurs (thick arrow), the effect of the control field is more dramatic. Here when $I_c=0$ the structure reflects back most of the light [Fig. 2(a), thin arrow]. As shown in Fig. 3(b) (solid line), this corresponds to a strong oscillation of the field inside the MQW structure. However, for $I_c=2$ MW/cm^2, although quite close to the Bragg wavelength, the oscillatory feature nearly disappears (dashed line). In other words, under these conditions at 1.549 eV the probe field does not experience the periodicity of the MQW structure. Away from this energy, where the peaks in Fig. 2(c) happen, the periodicity of the MQW structure remains visible to this field. Such a process represents a PEIT where at its central frequency the PBG structure acts a homogeneous medium. Figure 3(c) shows how transformation from a nearly complete band gap to the transparency state influences the overall probe field pattern inside the MQW structure. Note that the reduction of the amplitude associated with the transparency state can be attributed to the absorption of the probe field by the e1-hh1 excitons.

The features seen in Figs. 2 and 3 can be explained in terms of the interplay between the excitonic dressed effects and the superradiant properties of the MQW structure. In fact, for $I_c\neq 0$ the primary effect of the control field is to replace the bare exciton states with optically mixed states. To see this, note that in a dressed state picture with the ir field photon number N_p, these states can be as follows:[4]

$$|+\rangle_{1s}=\cos\theta|N_p,e1\text{-hh}1\rangle_{1s}+\sin\theta|(N_p-1),e2\text{-hh}1\rangle_{1s}, \quad (13)$$

$$|-\rangle_{1s}=\sin\theta|N_p,e1\text{-hh}1\rangle_{1s}-\cos\theta|(N_p-1),e2\text{-hh}1\rangle_{1s}. \quad (14)$$

Here $\tan 2\theta=-\Omega_c/\Delta_c$ where Ω_c is the Rabi frequency of the ir coupled QWs. As shown in the inset of Fig. 3(a), when

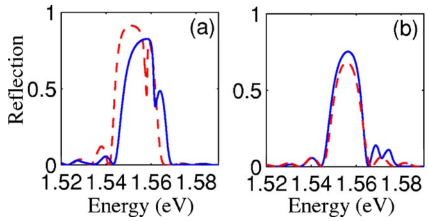

FIG. 4. (Color online) Reflection of the MQW ($N=150$) for various I_c and $\Delta_c=+20$ meV. The solid lines in (a) and (b) refer to $I_c=0.5$ and $2/\text{cm}$ MW/cm^2, respectively. Dashed line in (a) refers to the MQW response when $I_c=0$ and in (b) to the nonresonant PBG in the absence of the exciton effects.

$I_c=0$ the bare states are at 1.557 eV and are very narrow (~1 meV full width at half maximum). These states satisfy the Bragg condition needed for the dominance of superradiant excitons. For $I_c=2$ MW/cm^2, the bare states are substituted by the extensively broadened dressed states, i.e., $|+\rangle_{1s}$ and $|-\rangle_{1s}$ [Fig. 3(b) inset]. When $I_c=4$ MW/cm^2 or higher, the optical shift and broadening increase further, causing collapse of the superradiant mode. This leads to the disappearance of the excitonic effects, leaving alone the nonresonant PBG features. Under this condition, while there is almost no sign of excitons in the reflection [Fig. 2(d)], the transmission band gap is flanked by two absorption peaks associated with $|+\rangle_{1s}$ and $|-\rangle_{1s}$ (not shown).

The results presented in Fig. 2 were associated with $\Delta_c=-20$ meV. The picture changes dramatically if one considers $\Delta_c=+20$ meV (Fig. 4). Here, instead of formation of a doublet, one sees a direct adaptation of the reflection with that associated with the nonresonant PBG. This can be associated with the fact that here the $1s$ $e1$-hh1 states are blueshifted toward the Bragg wavelength of the nonresonant PBG. Therefore, for the range of I_c where the superradiant mode still exists [Fig. 4(b)], there is no competition between the resonant and nonresonant responses. For large values of I_c, similar to the case of Fig. 2, the superradiant mode collapses, leaving alone the nonresonant response.

In conclusion, we showed that ir dressing of superradiant excitons could lead to a photonic EIT where the absoprtion or index periodicity of the MQW structure was disguised from the probe field. This process, which is intrinsically different from the conventional EIT, is caused by coherent adjustment of the contributions of excitons and the background index contrast of the wells and barriers. We also showed that the dressing process could lead to total collapse of the superradiant mode and conversion of the structure into a nonresonant passive PBG structure. For experimental tests of these results one can consider the techniques used to observe MQW PBGs,[2] and investigate photoinduced intersubband transitions.[12]

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