Polariton local states in periodic Bragg multiple-quantum-well structures

Lev I. Deych

Department of Physics, Seton Hall University, South Orange, New Jersey 07079

A. Yamilov and A. A. Lisyansky

Department of Physics, Queens College of the City University of New York, Flushing, New York 11367

Received June 9, 2000

We study analytically the optical properties of several types of defect in Bragg multiple-quantum-well structures. We show that a single defect leads to two local polariton modes in the photonic bandgap. These modes lead to peculiarities in reflection and transmission spectra. Detailed recommendations for experimental observation of the effects studied here are given.

OCIS codes: 160.6000,240.5420, 230.1480.

It was demonstrated recently[1] that long multiple-quantum-well (MQW) systems can form optical lattices in which various quantum wells (QW's) are coherently coupled as a result of interaction with a retarded electromagnetic field. Light–matter interaction in such systems depends on the structure of the systems and can be significantly and controllably modified. A polariton formalism provides an adequate self-consistent way to describe the strong interaction of the QW excitons and the light in MQW systems.[2,3] These systems have become a subject of very active research in the past few years (see, for instance, Refs. 2–4 and references therein). Special attention has been paid to so-called Bragg structures in which interwell spacing a is exactly equal to the half-wavelength of light at the frequency of excitonic resonance, $\lambda_0/2 = a$.[1,5,6] Peculiarities of the Bragg structures follow from the fact that a photonic bandgap in the vicinity of the exciton frequency is degenerate. In other words, the bandgap is formed from two adjacent gaps with coinciding boundaries. Detuning the structure from the exact Bragg condition shifts those boundaries away from each other, giving rise to a conduction band between them.[7]

Should the periodicity in the arrangement of MQW's be locally altered, one would expect the appearence of defect local modes inside the photonic bandgaps. This phenomenon provides additional possibilities for controlling the optical properties of MQW's and therefore is of considerable interest. This idea was put forward in Ref. 8, in which a dispersion equation for frequencies of local modes with different polarizations was derived. In the case of transversely polarized excitation, the equation that describes MQW's is essentially equivalent to a model of a one-dimensional chain of dipoles that we used previously to discuss local polariton states in polar crystals.[9–11] In the context of MQW's, the local polariton states considered in Refs. 9–11 correspond to a mode localized in the growth direction of the MQW structure but extended in the in-plane directions.

In this Letter we study local defect polariton states in Bragg MQW structures and defect-induced changes in transmission and reflection spectra. Defect layers can differ from host layers in three ways: in exciton–light-coupling strength (Γ defect), in exciton resonance frequency (Ω defect), and in interwell spacing (a defect). Each of these types of defect can be reproduced experimentally, and we show below that each of them plays a distinctly different role in the optical properties of the system. We obtain closed analytical expressions for respective local frequencies as well as for reflection and transmission coefficients. On the basis of the results obtained, we give a practical recommendation for experimental observation of the studied effects in samples used in the research reported in Refs. 1 and 5.

The optical properties of QW's are usually described with the use of nonlocal susceptibility determined by energies and wave functions of a QW exciton.[2,12] For sufficiently thin QW's, a simplified approach is possible, in which the polarization density of the QW is presented in the form $P(\mathbf{r},z) = P_n(\mathbf{r})\delta(z - z_n)$, where $\mathbf{r}$ is an in-plane position vector, z_n represents a coordinate of the nth well, and P_n is the surface polarization density of the well. When light is incident in the direction of growth z of MQW's, $k_\parallel = 0$ and there are two independent degenerate transverse polarizations, T and L, that are not coupled to the longitudinal Z mode. In this case the dynamics of transverse modes can be described by

$$(\Omega_n^2 - \omega^2)P_n = (c/\pi)\Gamma_n E(z_n)\,, \tag{1}$$

$$\frac{\omega^2}{c^2}E(z) + \frac{\mathrm{d}^2E(z)}{\mathrm{d}z^2} = -4\pi\frac{\omega^2}{c^2}\sum_n P_n\delta(z - z_n)\,, \tag{2}$$

which coincide with equations used in Refs. 9–11 and 13 to describe one-dimensional chains of atoms. Here Ω_n and Γ_n are the excitonic frequency and the exciton–light coupling of the nth QW, respectively. In an infinite, pure system, all $\Gamma_n = \Gamma_0$, $\Omega_n = \Omega_0$, and $z_n = na = n\lambda_0/2$. The spectra of ideal MQW's have been

0146-9592/00/231705-03$15.00/0

studied in many papers.[2,3,7,13,14] In the specific case of Bragg structures, the exciton resonance frequency is at the center of the bandgap determined by the inequality $\omega_l = \Omega_0(1 - \sqrt{2\Gamma_0/\pi\Omega_0}) < \omega < \Omega_0(1 + \sqrt{2\Gamma_0/\pi\Omega_0}) = \omega_u$.[7] This bandgap is the frequency region where we look for new local states associated with defects in MQW's.

Ω and Γ defects introduce perturbations into the equation of motion that are localized at one site (diagonal disorder). Therefore they can be studied by the usual Green's function technique (see, for instance, Ref. 15). The resultant dispersion equations have the form

$$G_{\Omega,\Gamma} = \beta/2\sqrt{D}\,, \tag{3}$$

where $\beta = 4\Gamma_0\omega/(\omega^2 - \Omega_0^2)$ and $D = -1 + \beta^2/4 + \beta\cot(\omega a/c)$. For the Ω defect, the function is $G_\Omega = (\Omega_1^2 - \omega^2)/(\Omega_1^2 - \Omega_0^2)$; for the Γ defect, the respective function is $G_\Gamma = \Gamma_0/(\Gamma_1 - \Gamma_0)$. Ω_1 and Γ_1 denote the respective parameters of the defect layer. A similar equation for the Ω defect was studied in the long-wave approximation in Ref. 9. It was found that the equation has one real-valued solution for any $\Omega_1 > \Omega_0$. In the case of Bragg structures there are always two solutions for both types of defect, one below Ω_1 and one above. This is a manifestation of the degenerate nature of the bandgap in Bragg structures. Equation (3) can be solved approximately by use of the fact that $\Gamma_0 \ll \Omega_0$ in most cases. For the Ω defect, one solution demonstrates a radiative shift from the defect frequency Ω_1:

$$\omega_{\mathrm{def}}^{(1)} = \Omega_1 - \Gamma_0(\Omega_1 - \Omega_0)/[(\omega_u - \Omega_1)(\Omega_1 - \omega_l)]^{1/2}, \tag{4}$$

whereas the second solution splits off the upper or lower boundary, depending on the sign of $\Omega_1 - \Omega_0$:

$$\omega_{\mathrm{def}}^{(2)} = \omega_{u,l} \pm \pi^2(\omega_u - \omega_l)(\Omega_1 - \Omega_0)^2/4\Omega_0^2\,, \tag{5}$$

where one chooses ω_u and minus for $\Omega_1 > \Omega_0$, and ω_l and plus in the opposite case. In the case of the Γ defect, both solutions appear in the vicinity of the gap boundaries $\omega_{\mathrm{def}}^{(1,2)} = \omega_{u,l} \pm 2(\Gamma_1 - \Gamma_0)^2(\omega_u - \omega_l)$. These solutions exist only for $0 < \Gamma_1 < \Gamma_0$ and are close to the gap boundaries. One could expect, therefore, that the states at these frequencies will be vulnerable to even a weak dissipation and will not significantly affect the optical spectra of the system.

The a defect significantly differs from the two other types. An increase in the interwell distance between any two wells automatically changes the coordinates of an infinite number of wells: $z_n = na$ for $n \le n_d$ and $z_n = (b - a) + na$ for $n_d < n$, where b is the distance between the n_dth and the $(n_d + 1)$st wells. Therefore this defect is nonlocal and cannot be treated with the same methods as in the two cases described above. The best approach to this situation is to match solutions of semi-infinite chains for $n < n_d$ and $n > n_d + 1$ with a solution for $na < z < na + (b - a)$. Solutions for semi-infinite chains can be constructed by the transfer matrix approach, in which the state of the system is described by a two-dimensional vector v_n with components $E(z_n)$ and $(c/\omega)\mathrm{d}E(z_n)/\mathrm{d}z$. Propagation of this vector through the system is described by the transfer matrix $\hat{\tau}_n$:

$$\widehat{\tau_n} = \begin{bmatrix} \cos[(\omega/c)a_n] + \beta\sin[(\omega/c)a_n] & \sin[(\omega/c)a_n] \\ -\sin[(\omega/c)a_n] + \beta\cos[(\omega/c)a_n] & \cos[(\omega/c)a_n] \end{bmatrix}, \tag{6}$$

where $a_n = z_{n+1} - z_n$. As a result, one obtains the dispersion equation for the defect mode in terms of elements of the total transfer matrix $\hat{T}$, equal to the product of all site matrices $\hat{\tau}$:

$$(T_{11} + T_{22}) - i(T_{12} - T_{21}) = 0\,. \tag{7}$$

In the limit of an infinitely long system, the imaginary part of Eq. (7) vanishes, and one has a real-valued dispersion equation for the frequency of a stationary local mode. Using Eq. (6), one can write Eq. (7) for an infinite MQW system as

$$\cot[(\omega/c)b] = -\{\sin[(\omega/c)a] - \beta\lambda_-/2\}/\{\cos[(\omega/c)a] - \lambda_-\}, \tag{8}$$

where $\lambda_- = [\cot(\omega a/c) + \beta/2 - \sqrt{D}]\sin(\omega a/c)$ is one of the eigenvalues of the transfer matrix [Eq. (6)]. This equation also has two solutions, above and below Ω_0. Assuming that $\sqrt{\Gamma_0}\,b/\Omega_0 a \ll 1$, one of these solutions can be expressed as

$$\omega_{\mathrm{def}}^{(1)} = \Omega_0 - \frac{\omega_u - \omega_l}{2} \times \frac{(-1)^{(\xi+1/2)}\sin(\pi/2)\xi}{1 + (\omega_u - \omega_l)/2\Omega_0(b/a)(-1)^{(\xi+1/2)}\cos(\pi/2)\xi}, \tag{9}$$

where $\xi = b/a$ and [ ] denotes an integer part. We can obtain the second solution from Eq. (9) by replacing ξ with $\xi + 1$. Therefore, for $\Gamma_0 \ll \Omega_0$ and ξ not very large, both solutions are almost periodic functions of b/a with a period of 1.

The expression on the left-hand side of Eq. (7) coincides with the denominator of the transmission and reflection coefficients in a system of finite length, and, with the appropriate choice of transfer matrices τ, the equation $T_{11} + T_{22} = 0$ produces dispersion equations for local states of all three types of defect, including Eq. (3) for Ω and Γ defects. In the absence of homogeneous broadening of the exciton resonance, the defects would cause a resonance increase in transmission at the local mode frequency.[9] The resonance occurs when the defect is placed at the center of the system. Then the maximum transmission becomes independent of the system's length, and in the case of Ω and Γ defects it can be presented as

$$|t_{\max}|^2 = 1 - 4\left[\left(\frac{\omega_{\mathrm{def}} - \Omega_0}{\omega_u - \Omega_0}\right)^2 - \frac{1}{2}\right]^2. \tag{10}$$

In the absence of absorption, transmission reaches unity if the frequency of the local state is $\omega_{\mathrm{def}} = \Omega_0 \pm (\omega_u - \Omega_0)/\sqrt{2}$. This cannot happen for the Γ defect,

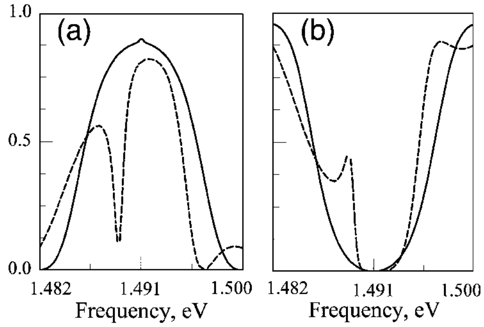

Fig. 1. (a) Reflection and (b) transmission coefficients of 200 QW's with the defect placed in the middle of the stack. Solid curves, the pure system; dashed curves, the system with the a defect ($b/a = 1.3$).

because frequencies of the respective local states always lie close to the band edge, but for the Ω defect it is quite possible to create the state with the required frequency.

For the third type of defect, the resonance transmission also takes place when the defect is in the center of the stack. $t_{\max}$, in this case, can be expressed as

$$|t_{\max}|^2 = 8\left\{\frac{\omega_{\rm def} - \Omega_0}{\omega_u - \Omega_0}\left[1 - \left(\frac{\omega_{\rm def} - \Omega_0}{\omega_u - \Omega_0}\right)^2\right]\right\}^2. \quad (11)$$

It becomes unity for two frequencies: $\omega_{\rm def}^{(1,2)} = \Omega_0 \pm (\omega_u - \Omega_0)/\sqrt{2}$ located symmetrically with respect to the center of the gap. As one can see from Eq. (9), these conditions can be satisfied for both defect frequencies at the same time when $b \simeq (\text{integer} + 1/2)a$.

In a real system, enhancement of the transmission coefficient is usually limited by homogeneous broadening. Two situations are possible when exciton damping is taken into account. Damping can suppress the resonance transmission, and the presence of the local states will be observed only as an enhancement of absorption at the local frequency. This case can be called a weak-coupling regime for the local state, when incident radiation is resonantly absorbed by a local exciton state. The opposite case, when the resonance transmission persists in the presence of damping, can be called a strong-coupling regime. In this case there is coherent coupling between excitons and the electromagnetic field, so the local state can appropriately be called a local polariton. Among the three types of defect considered here, the Γ defect is less likely to survive absorption because of the proximity of the respective frequencies to bandgap edges. For the Ω defect, one of the local frequencies appears far enough from the boundaries and can be less sensitive to absorption. However, the width of the transmission resonance is determined by its radiative shift from Ω_1, where transmission goes to zero. This shift is rather small, and small absorption can still suppress the resonance transmission. Therefore the best candidate to produce a local polariton state in the strong-coupling regime is the a defect.

To account quantitatively for homogeneous broadening we add an imaginary part to the exciton polarizability, $\beta = 4\Gamma_0\omega/(\omega^2 - \Omega_0^2 + 2i\gamma\omega)$. For numerical calculations we use parameters from Ref. 1: $\Omega_0 = 1.491$ eV, $\Gamma_0 = 27$ meV, and $\gamma = 0.28$ eV. The localization length at the center of the forbidden bandgap is in this case $\sim 80a$, whereas the length of the samples used reached $100a$. Figure 1 presents plots of reflection and transmission for a MQW system with an a defect for which the resonance transmission is the most pronounced. We can conclude that the interwell spacing defect gives rise to local polariton states in regular MQW InGaAs/GaAs MQW's. These states manifest themselves in strong resonant tunneling of light through a MQW system with 100 or more wells and can be observed in transmission experiments. These states can be also viewed as long-lived optical waveguide modes[8] in Bragg MQW structures. The a type of defect can be implemented experimentally and presents additional opportunities for controlling light–matter interaction and photonic engineering. It gives a unique opportunity for observing and studying local polaritons.

We are pleased to acknowledge valuable remarks of H. M. Gibbs and G. Khitrova. We are indebted to S. Schwarz for reading and commenting on the manuscript. This research was partially supported by NATO linkage grant N974573 and a Professional Staff Congress of the City University of New York Research Award. A. A. Lisyansky's e-mail address is lisyansky@qc.edu.

References

1. M. Hübner, J. P. Prineas, C. Ell, P. Brick, E. S. Lee, G. Khitrova, H. M. Gibbs, and S. W. Koch, Phys. Rev. Lett. **83,** 2841 (1999).
2. D. S. Citrin, Solid State Commun. **80,** 139 (1994).
3. L. C. Andreani, Phys. Lett. A **192,** 99 (1994); Phys. Status Solidi B **188,** 29 (1995).
4. T. Stroucken, A. Knorr, P. Thomas, and S. W. Koch, Phys. Rev. B **53,** 2026 (1996).
5. M. Hübner, J. Kuhl, T. Stroucken, A. Knorr, S. W. Koch, R. Hey, and K. Ploog, Phys. Rev. Lett. **76,** 4199 (1996).
6. M. P. Vladimirova, E. L. Ivchenko, and A. V. Kavokin, Semiconductors **32,** 90 (1998).
7. L. I. Deych and A. A. Lisyansky, Phys. Rev. B **62,** 4242 (2000).
8. D. S. Citrin, Appl. Phys. Lett. **66,** 994 (1995).
9. L. I. Deych and A. A. Lisyansky, Phys. Lett. A **243,** 156 (1998); L. I. Deych, A. Yamilov, and A. A. Lisyansky, Europhys. Lett. **46,** 524 (1999).
10. L. I. Deych, A. Yamilov, and A. A. Lisyansky, Phys. Rev. B **59,** 11339 (1999).
11. A. Yamilov, L. I. Deych, and A. A. Lisyansky, J. Opt. Soc. Am. B **17,** 1498 (2000).
12. L. C. Andreani, in *Confined Electrons and Photons*, E. Burstein and C. Weisbuch, eds. (Plenum, New York, 1995), p. 57.
13. I. H. Deutsch, R. J. C. Spreeuw, S. L. Rolston, and W. D. Phillips, Phys. Rev. A **52,** 1394 (1995).
14. E. L. Ivchenko, Phys. Solid State **33,** 1344 (1991).
15. I. M. Lifshitz, Nuovo Cimento Suppl. A1 **3,** 716 (1956).