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Interaction of quantum well excitons with evanescent plane electromagnetic waves

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Abstract
A formalism based on non-local dielectric response theory and Green function techniques has been developed to describe the interaction of quantum well excitons with an evanescent optical wave of a planar waveguide. Reflection spectra of a system in which a quantum well placed behind a dielectric interface at which light experiences total internal reflection have been calculated. It is shown that the spectral feature corresponding to the exciton resonance becomes much more pronounced if the angle of incidence is close to the critical angle of total internal reflection. The concept of a generalized Snell law has been applied to provide simplification of the formalism.

1. Introduction

The waveguiding regime of light propagation is very widely studied in the context of optical fibres, laser resonators, and other different kinds of conventional optoelectronic devices. However, a new generation of optoelectronic devices is now being discussed: namely, polariton devices that exploit the coupling of light with excitons, which are the elementary excitations of semiconductor crystals. Exciton–polaritons combine the properties of electromagnetic modes (high group velocities, large coherence lengths) and excitons (finite effective masses, dipole moments, coupling by exchange interaction). Also, their bosonic nature suggests that it should be possible to form a Bose condensate which would result in spontaneous emission of coherent, monochromatic light, which is currently referred to as the polariton laser effect [1–3]. One of the main obstacles on the way towards realization of polariton devices is the finite radiative lifetime of these half-light–half-matter quasiparticles. Polaritons with small in-plane wavevectors have a lifetime of only a few picoseconds, which makes any manipulations with them difficult. On the other hand, polaritons having in-plane wavevectors $k_x > \omega/c$, where ω is their angular frequency, propagate in the waveguiding regime and have a much longer (theoretically, infinite) lifetime. Clearly, these long-lived polaritons are more suitable for device applications.

Further, they can be excited using a prism or a diffraction grating, and detected in the same way as light. It is essential to model the propagation of waveguided polariton modes correctly and to calculate the in-plane and normal electric field profiles induced in the structure by them.

This paper presents a semi-classical formalism describing the optical properties of the waveguided polariton modes in the framework of the generalized transfer matrix technique. We focus especially on the states situated at the edge of the light-cone (i.e. having $k_x \approx \omega/c$) and show that light coupling with the quantum well exciton resonance is greatly enhanced for these modes, and the reflectivity of quantum wells measured via a prism at the light-cone edge exhibits extremely strong resonant features. Popov *et al* [4] have recently published a theoretical paper on the exciton–polaritons at the light-cone edge. However, their work did not focus on reflection from a quantum well in the waveguiding regime, which is the main subject of interest in this paper. The aim of the present work is to investigate the interaction of a quantum well exciton with an evanescent electromagnetic field in a quantum well placed behind an interface at which light experiences total internal reflection or situated in the cladding layer of a planar optical waveguide.

2. Total internal reflection at a dielectric interface

Consider an interface normal to the z-direction between two dielectric media (labelled 1 and 2) having refractive indices n_1 and n_2, respectively, where $n_1 > n_2$. Let a plane light wave of angular frequency ω be incident on medium 2 from medium 1 at an angle of incidence θ in the x–z plane. When θ exceeds the critical value θ_c defined by the equation

$$\sin\theta_c = \frac{n_2}{n_1}, \tag{1}$$

total internal reflection takes place with the amplitude of the electromagnetic wave decaying exponentially in medium 2 and the reflection coefficient becoming equal to unity. Also, at this critical angle of incidence, the component of the wavevector $\mathbf{k}$ perpendicular to the interface becomes imaginary in medium 2, and there the spatial variation of the electric field E is described by

$$E \propto \exp(-\chi z)\exp(\mathrm{i}k_x x), \tag{2}$$

where

$$\begin{aligned} k_z &= \mathrm{i}\chi, \\ k_x^2 - \chi^2 &= k^2, \\ k &= n_2\frac{\omega}{c}, \end{aligned} \tag{3}$$

and c is the velocity of light in a vacuum.

Formally, such an electromagnetic field can be considered as a plane wave propagating in the direction defined by the complex angle θ_2

$$\theta_2 = -\mathrm{i}\ln\big(\mathrm{i}b + \sqrt{1-b^2}\big), \tag{4}$$

where $b = (n_1/n_2)\sin\theta_1$.

Equation (4) is a generalized form of Snell's law, and it gives a proper mathematical description of the electromagnetic field in the case of total internal reflection. It is also valid in absorbing and left-handed materials [5]. Bearing in mind equation (4), one can write

$$k_z = \mathrm{i}\chi = k\cos\theta_2, \tag{5}$$

where the cosine of the angle of refraction is seen to be purely imaginary.

3. Reflections and transmission at a quantum well close to the excitonic resonance

3.1. Solution of the inhomogeneous wave equation

The electromagnetic field in the vicinity of the quantum well (QW) in the spectral region near the QW exciton resonance, ω_0, is described by the wave equation

$$\nabla \times \nabla \times \mathbf{E} - \varepsilon_b k_0^2 \mathbf{E} = 4\pi k_0^2 \mathbf{P}_{\mathrm{exc}}, \tag{6}$$

where the right-hand term accounts for the excitonic contribution to the dielectric polarization $\mathbf{P}_{\mathrm{exc}}$ and ε_b is the background dielectric constant in the QW and surrounding barrier material (assumed to be equal for simplicity). According to the theory of non-local dielectric response [6], the polarization $\mathbf{P}_{\mathrm{exc}}$ can be written as

$$\mathbf{P}_{\mathrm{exc}}(z) = \frac{1}{4\pi} \int \tilde{T}(\omega, z, z') E(z')\, \mathrm{d}z', \tag{7}$$

where the non-local dielectric susceptibility is given by $\tilde{T}(\omega, z, z') = T(\omega)\Phi(z)\Phi(z')$, the function $\Phi(\mathbf{r})$ is proportional to the envelope of the exciton wavefunction taken with equal electron and hole coordinates $\Psi(\mathbf{r}, \mathbf{r}')$: $\Phi(\mathbf{r}) = \frac{1}{\sqrt{\mathbf{S}}}\Psi(\mathbf{r}, \mathbf{r})$ and

$$T(\omega) = \frac{\varepsilon_b \omega_{\mathrm{LT}} \pi a_{\mathrm{B}}^3}{\omega_0 - \omega - \mathrm{i}\Gamma}. \tag{8}$$

Here, ω_{LT} and a_{B} are respectively the longitudinal–transverse splitting and the Bohr radius of the exciton in the bulk material, and Γ is the non-radiative exciton damping.

In the presence of a QW, the wave equation can be solved within non-local dielectric response theory by a Green function method [1, 6]. The electric field $\mathbf{E}(\mathbf{r})$ can be represented in the form

$$\mathbf{E}(\mathbf{r}) = \mathbf{E}^0 + \tilde{\mathbf{E}}, \tag{9}$$

where $\mathbf{E}^0$ is the solution to the homogeneous counterpart of equation (6) and $\tilde{\mathbf{E}}$ can be expressed in terms of the Green function $G_{\alpha\beta}$ as

$$\tilde{E}_\alpha(\mathbf{r}) = k_0^2 T \Lambda_\beta \int G_{\alpha\beta}(x, z - z')\Phi(z')\, \mathrm{d}z', \tag{10}$$

where $\alpha, \beta = x, y, z$, $\Lambda_\beta = \int \Phi(z) E_\beta(z)\, \mathrm{d}z$ and

$$G_{\alpha\beta}(\mathbf{r}) = \left(\delta_{\alpha\beta} + \frac{1}{k^2} \frac{\partial^2}{\partial r_\alpha\, \partial r_\beta}\right) \bar{G}(x, z - z') \tag{11}$$

is the Green function.

3.2. Propagating waves

Normally we are interested in the interaction of propagating waves with the QW. For a propagating wave, the homogeneous solution $\mathbf{E}^0$ is a plane wave having the form $E_\alpha^0 \exp(\mathrm{i}k_z z + \mathrm{i}k_x x)$, and the Green function is given by

$$\bar{G}(x, z - z') = \frac{\mathrm{i}}{2k_z} \exp(\mathrm{i}k_z |z - z'| + \mathrm{i}k_x x). \tag{12}$$

Multiplying the left- and right-hand sides of equation (10) by $\Phi(z)$ and integrating over z, we can obtain the electric field [6]. Denoting

$$\Gamma_0 = \frac{k}{2} k \omega_{\mathrm{LT}} \pi a_{\mathrm{B}}^3 \left(\int \Phi(z) \cos k_z z\, \mathrm{d}z\right)^2, \tag{13}$$

the reflection and transmission coefficients of the QW for TE-polarized light (with electric field $\mathbf{E} = (0, E_y, 0)$) can be expressed as

$$r_{QW} = \frac{i\tilde{\Gamma}_0}{\omega_0 - \omega - i(\Gamma + \tilde{\Gamma}_0)} \tag{14}$$

$$t_{QW} = 1 + r_{QW}, \tag{15}$$

where

$$\tilde{\Gamma}_0 = \Gamma_0/\cos\theta \tag{16}$$

is the exciton radiative broadening, and

$$\tilde{\omega}_0 = \omega_0 + \frac{k^2}{2k_z}\omega_{LT}\pi a_B^3 \int\int \Phi(z)\Phi(z') \sin k_z|z - z'|\, dz\, dz' \tag{17}$$

is the renormalization of the resonance frequency due to the polariton effect.

Similar expressions can be found for the reflection and transmission coefficients in TM polarization (with electric field $\mathbf{E} = (E_x, 0, E_z)$):

$$r_{QW} = \frac{i\tilde{\Gamma}_0'}{\tilde{\omega}_0' - \omega - i(\Gamma + \tilde{\Gamma}_0')} - \frac{i\tilde{\Gamma}_0''}{\tilde{\omega}_0'' - \omega - i(\Gamma + \tilde{\Gamma}_0'')} \tag{18}$$

$$t_{QW} = 1 + \frac{i\tilde{\Gamma}_0'}{\tilde{\omega}_0' - \omega - i(\Gamma + \tilde{\Gamma}_0')} + \frac{i\tilde{\Gamma}_0''}{\tilde{\omega}_0'' - \omega - i(\Gamma + \tilde{\Gamma}_0'')}, \tag{19}$$

where

$$\tilde{\Gamma}_0' = \Gamma_0 \cos\theta, \tag{20}$$

$$\tilde{\omega}_0' = \omega_0 + \frac{k_z}{2}\omega_{LT}\pi a_B^3 \int\int \Phi(z)\Phi(z') \sin k_z|z - z'|\, dz'\, dz, \tag{21}$$

$$\tilde{\Gamma}_0'' = \Gamma_0 \sin^2\theta/\cos\theta, \tag{22}$$

$$\tilde{\omega}_0'' = \omega_0 + \frac{k_x^2}{2k_z}\omega_{LT}\pi a_B^3 \int\int \Phi(z)\Phi(z') \sin k_z|z - z'|\, dz'\, dz + \omega_{LT}\pi a_B^3 \int [\Phi(z)]^2\, dz. \tag{23}$$

The exciton polaritons' eigenfrequencies are given by the poles of r_{QW} (equations (14) and (18)).

3.3. Evanescent wave solutions

In the case of a QW placed behind a dielectric interface at which total internal reflection occurs, it is possible to have an evanescent wave interact with a QW. Although now there are only evanescent fields, it is still convenient to define reflection and transmission coefficients in a formal mathematical fashion, even though they may be regarded in the physical sense as a 'figure of speech'. For the purpose of developing the theory we consider the somewhat artificial case where an 'incident' evanescent wave field exists in the absence of a dielectric interface. The reflection coefficient is then introduced as the ratio of the amplitude of the wave decaying to $z = +\infty$ (incident field) to the amplitude of the wave decaying to $z = -\infty$ (reflected field) far to the left from the QW.

The basic equations used for solving the wave equation (6) are similar to before, but the Green function is now

$$\bar{G}(x, z - z') = \frac{1}{2\chi} \exp[-\chi|z| + ik_x x], \tag{24}$$

and the solution to the homogeneous wave equation is an evanescent plane wave having form $E^0_\alpha \exp(-\chi z)\exp(\mathrm{i}k_x x)$.

For the TE-polarization case, the electric field obeys the equation

$$E_y(z) = E_{0,y}\mathrm{e}^{-\chi z} + k_0^2 T\Lambda_y \int \Phi(z')\bar{G}(z-z')\,\mathrm{d}z', \tag{25}$$

and for the TM-polarization case, the non-zero components of electric field satisfy the equations

$$E_x(z) = E_{0,x}\mathrm{e}^{-\chi z} - k_0^2 T\Lambda_x \frac{\chi^2}{k^2}\int \Phi(z')\bar{G}(z-z')\,\mathrm{d}z' + \mathrm{i}\frac{k_x}{k^2}k_0^2 T\Lambda_z \times \int \Phi(z')\frac{\partial}{\partial z}\bar{G}(z-z')\,\mathrm{d}z' \tag{26}$$

$$E_z(z) = E_{0,z}\mathrm{e}^{-\chi z} + \mathrm{i}\frac{k_x}{k^2}k_0^2 T\Lambda_x \int \Phi(z')\frac{\partial}{\partial z}\bar{G}(z-z')\,\mathrm{d}z' + \frac{k_x^2}{k^2}k_0^2 T\Lambda_z \times \int \Phi(z')\bar{G}(z-z')\,\mathrm{d}z' - \frac{1}{k^2}k_0^2 T\Lambda_z\Phi(z). \tag{27}$$

Then the quantities Λ_α are explicitly found by multiplying the electric field projections by the function $\Phi(z)$ and integrating over z. Hence, defining $L = \int\int \Phi(z)\Phi(z')\bar{G}(z-z')\,\mathrm{d}z\,\mathrm{d}z'$,

$$\Lambda_y = \frac{E_{0,y}\int \Phi(z)\mathrm{e}^{-\chi z}\,\mathrm{d}z}{1 - k_0^2 T L}, \tag{28}$$

$$\Lambda_x = \frac{E_{0,x}\int \Phi(z)\mathrm{e}^{-\chi z}\,\mathrm{d}z}{1 + k_0^2 T \frac{\chi^2}{k^2} L}, \tag{29}$$

$$\Lambda_z = \frac{E_{0,z}\int \Phi(z)\mathrm{e}^{-\chi z}\,\mathrm{d}z}{1 - \frac{k_x^2}{k^2}k_0^2 T L + \frac{1}{k^2}k_0^2 T\int[\Phi(z)]^2\,\mathrm{d}z}. \tag{30}$$

In order to obtain the reflection coefficient, we must have the ratio of the amplitude of the left-decaying wave (coefficient of $\mathrm{e}^{+\chi z}$) to the amplitude of the right-decaying wave (coefficient of $\mathrm{e}^{-\chi z}$) at $z \to -\infty$. The transmission coefficient is similarly found from the ratios of the right-decaying waves (coefficients of $\mathrm{e}^{-\chi z}$) at each side of the QW.

Defining (compare equation (13))

$$\Gamma_{\mathrm{e}0} = \frac{k}{2}\omega_{\mathrm{LT}}\pi a_{\mathrm{B}}^3\left(\int \Phi(z)\cosh\chi z\,\mathrm{d}z\right)^2, \tag{31}$$

we express the ‘reflection’ and ‘transmission’ coefficients for the TE polarization as

$$r_{\mathrm{QW}} = \frac{\tilde{\Gamma}_{\mathrm{e}}}{\tilde{\omega}_0 - \tilde{\Gamma}_{\mathrm{e}} - \omega - \mathrm{i}\Gamma}, \tag{32}$$

$$t_{\mathrm{QW}} = 1 + \frac{\tilde{\Gamma}_{\mathrm{e}}}{\tilde{\omega}_0 - \tilde{\Gamma}_{\mathrm{e}} - \omega - \mathrm{i}\Gamma} = 1 + r_{\mathrm{QW}}, \tag{33}$$

respectively. Here

$$\tilde{\Gamma}_{\mathrm{e}} = \frac{k}{\chi}\Gamma_{e0}, \tag{34}$$

$$\tilde{\omega}_0 = \omega_0 - \frac{k^2}{2\chi}\omega_{\mathrm{LT}}\pi a_{\mathrm{B}}^3 \int\int \Phi(z)\Phi(z')\sinh\chi|z-z'|\,\mathrm{d}z\,\mathrm{d}z'. \tag{35}$$

The reflection and transmission coefficients for the TM polarization are, respectively

$$r_{\mathrm{QW}} = -\frac{\tilde{\Gamma}'_{\mathrm{e}}}{\tilde{\omega}'_0 - \tilde{\Gamma}'_{\mathrm{e}} - \omega - \mathrm{i}\Gamma} - \frac{\tilde{\Gamma}''_{\mathrm{e}}}{\tilde{\omega}''_0 - \tilde{\Gamma}''_{\mathrm{e}} - \omega - \mathrm{i}\Gamma} = r_1 - r_2, \tag{36}$$

$$t_{\rm QW} = 1 - \frac{\tilde{\Gamma}'_{\rm e}}{\tilde{\omega}'_0 - \tilde{\Gamma}'_{\rm e} - \omega - {\rm i}\Gamma} + \frac{\tilde{\Gamma}''_{\rm e}}{\tilde{\omega}''_0 - \tilde{\Gamma}''_{\rm e} - \omega - {\rm i}\Gamma} = 1 + r_1 + r_2, \tag{37}$$

where

$$\tilde{\Gamma}'_{\rm e} = \frac{\chi}{k}\Gamma_{e0}, \tag{38}$$

$$\tilde{\omega}'_0 = \omega_0 + \frac{\chi}{2}\omega_{\rm LT}\pi a_{\rm B}^3 \int\int \Phi(z)\Phi(z')\sinh\chi|z - z'|\,{\rm d}z'\,{\rm d}z, \tag{39}$$

$$\tilde{\Gamma}''_{\rm e} = \frac{k_x^2}{\chi k}\Gamma_{e0}, \tag{40}$$

$$\tilde{\omega}''_0 = \omega_0 - \frac{k_x^2}{2\chi}\omega_{\rm LT}\pi a_{\rm B}^3 \int\int \Phi(z)\Phi(z')\sinh\chi|z - z'|\,{\rm d}z'\,{\rm d}z + \omega_{\rm LT}\pi a_{\rm B}^3 \int [\Phi(z)]^2\,{\rm d}z. \tag{41}$$

The exciton polariton eigenenergies are given by the poles of $r_{\rm QW}$ for the TE case (equation (32)) and TM case (equation (36)). Thus the transverse (T) polariton has a resonant frequency $\tilde{\omega}_0$, and a width given by the non-radiative damping Γ. Likewise, the L- and Z-polaritons have resonant frequencies of $\tilde{\omega}'_0$ and $\tilde{\omega}''_0$, respectively.

It is apparent from equations (32) and (36) that, in the evanescent field case, there is no radiative damping of the exciton, in contrast to the case when a freely propagating wave is incident upon the QW. Mathematically, this is because in the denominator the quantity corresponding to $\tilde{\Gamma}_{\rm e}$ becomes real in the case of evanescent waves and contributes only to the resonance frequency renormalization.

This result can also be easily understood from a physical point of view. The radiative damping in the propagating wave case exists because the photons can propagate infinitely far away from the QW plane. However, in the evanescent wave case, light is no longer free to propagate in the direction out of the plane of the QW, and therefore the radiative exciton damping is zero. Consequently, surface exciton–polaritons are 'dark', as they have an infinite radiative lifetime [7].

Note that, if the angle of incidence corresponds to the critical angle of total internal reflection (e.g. $k_z \to 0$ or $\chi \to 0$), so that the polariton state is at the edge of the light-cone, the reflection coefficients for the cases of propagating (14) and evanescent (32) waves coincide and are both equal to -1.

Normally the electric field does not vary significantly over the distance that $\Phi(z)$ decays exponentially to an insignificant value, and hence $\Gamma_{e0} \approx \Gamma_0$. Using equation (5) and replacing Γ_{e0} with Γ_0 in equations (31), (38) and (40), we find that to a good approximation one can use equations (14) and (18) with a complex value for the angle of propagation in the medium, instead of equations (32) and (36).

4. Generalized transfer matrix

In a basis of left- and right-decaying evanescent waves, the transfer matrix relating the electric field on either side of the quantum well is identical to the well-known one for propagating waves [1]:

$$T = \frac{1}{t_{\rm QW}}\begin{pmatrix} t_{\rm QW}^2 - r_{\rm QW}^2 & r_{\rm QW} \\ -r_{\rm QW} & 1 \end{pmatrix}. \tag{42}$$

When a complex angle is used, the matrix in the basis of tangential components of electric and magnetic fields is identical to the one used for propagating waves. For example, the matrix in the TE polarization is

$$\hat{T}_i = \begin{pmatrix} 1 & 0 \\ -2n\cos\theta\frac{r}{1+r} & 1 \end{pmatrix}. \tag{43}$$

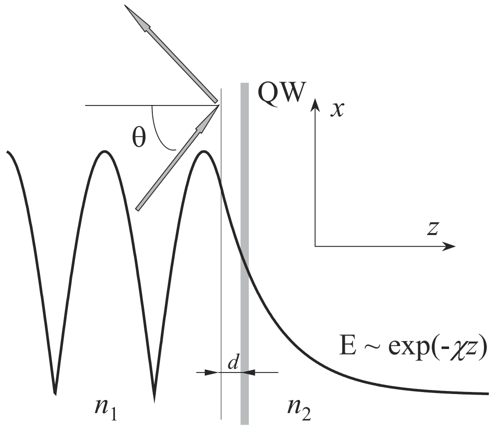

Figure 1. A profile of an envelope of the electric field in the case of total internal reflection of the incident light at the interfaces between the media having refractive indices n_1 and n_2. A quantum well is placed behind the interface.

Such results are significant because they mean that transfer matrix methods which facilitate the calculation of the transmission and reflection spectra of the layered structures can be used to consider the interaction of evanescent electromagnetic waves with the QW.

For the structure shown in figure 1, the amplitude reflection coefficient R can be easily expressed in terms of the amplitude reflection coefficients r_i and r_{QW} of the interface between the media 1 and 2 and of the quantum well, respectively:

$$R = \frac{r_i + r_{QW}\exp(-\chi d)}{1 + r_i r_{QW}\exp(-\chi d)}, \tag{44}$$

where d is the distance between the interface and the nearest edge of the quantum well.

Figure 2 shows some reflection spectra of TE-polarized light for the structure shown in figure 1. The parameters of the structure are chosen to correspond to an experimentally feasible III–V heterostructure. The refractive indices $n_1 = 3.7$ and $n_2 = 3.0$ correspond to GaAs and AlAs, respectively, and the critical angle of total internal reflection for the interface, given by equation (1), is $\theta_c = 54.176°$. The parameters of the quantum well are chosen to be those of an InGaAs quantum well of thickness 50 nm, with $\hbar\omega_0 = 0.9$ eV, $\hbar\Gamma_0 = 0.05$ meV, and $\hbar\Gamma = 1$ meV. The distance between the interface and the quantum well is taken as $d = 100$ nm. Substantial angles of incidence on the GaAs/AlAs interface can be achieved by using a prism or diffraction grating on the external boundary of the sample.

It follows from equations (16) and (34), that the interaction of light and a quantum well exciton is enhanced dramatically if the angle of incidence is close to the critical value θ_c. Figure 2 shows that if the incidence angle θ is much less than θ_c, the reflection spectrum is represented by the usual type of resonant curve with a relatively small magnitude modulation [8]. However, an increase of θ leads to a growth of the modulation, and the shape of the spectrum changes substantially. When θ becomes close to θ_c, the background reflectivity tends to unity and the Lorentzian dip appears at the frequency of the exciton resonance. Note that the position of the dip corresponds to the exciton resonance frequency ω_0, despite the fact that the pole in r_{QW} (equation (32)) is shifted from ω_0 by $\tilde{\Gamma}_e$, which itself is quite large if the incidence angle is close to θ_c. A subsequent increase of the incidence angle leads to a decrease of the depth of the excitonic dip.

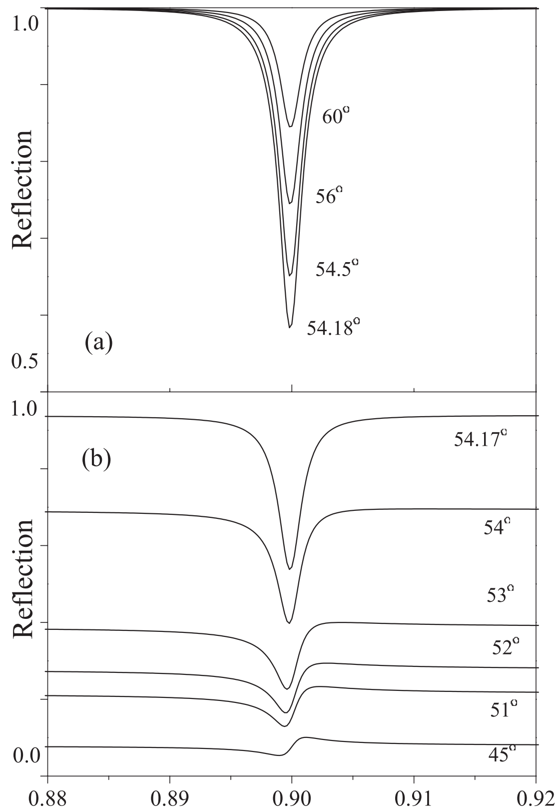

Figure 2. Reflection spectra of the structure shown in figure 1. (a) For incidence angles equal to the critical angle and exceeding it: 54.18°; 54.5°; 56°; 60°; (b) for incidence angles below the critical angle: 54.17°; 54°; 53°; 52°; 51°, 45°.

Note that if θ is exactly equal to the critical value θ_c, the relative depth of the excitonic dip exceeds 50% for the chosen parameters, despite the non-radiative damping of the exciton being 20 times stronger than its radiative counterpart. Such amplification of the excitonic modulation of reflection spectra makes the spectroscopy of quantum well excitons in the geometry of total internal reflection a promising tool for the precise determination of excitonic parameters.

5. Conclusions

We have developed a semi-classical formalism providing a description of the interaction of a quantum well exciton with an evanescent optical wave in the waveguiding regime. We have considered structures having a quantum well placed behind an interface at which light experiences total internal reflection or in the cladding layer of a planar waveguide. We have obtained an exact solution of Maxwell equations in the TE- and TM-polarizations and have given simplified expressions for the effective reflection and transmission coefficients of the QW using a generalized form of Snell's law. We have shown that if the angle of incidence is equal to the critical angle of total internal reflection, the excitonic resonance in the reflection spectrum becomes more pronounced and is an order of magnitude greater than in the normal incidence case.

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