Chapter 4
Dielectric Confinement Affected Exciton-Polariton Properties of the Semiconductor Nanowires

K. H. Aharonyan, E. M. Kazaryan, and E. P. Kokanyan

Abstract The nonlocal susceptibility with the cylindrical symmetry in semiconductor quantum wires embedded in the dielectric barrier environment (DQWW) is calculated. The strong dependence from the wire radius R due to the dielectric mismatch effect is established ($\sim R^{-8/3}$). It has been received that the oscillator strength of the one-dimensional (1D) excitons in DQWW increase strongly with decreasing R. The Maxwell's equations are solved in presence with the nonlocal excitonic response to the dielectric polarization for the DQWW and in a result the effective boundary conditions just considering a quantum wire presence in the structure have been established. The scattering coefficient of the light incident on the DQWW is obtained which strongly depends on the DQWW material parameters such as: (a) quantum wire radius R and dielectric constant ε_w, (b) barrier environment dielectric constant ε_b. This made it possible to obtain the dispersion spectrum and lifetime for the 1D exciton-polaritons of the DQWW near the exciton resonance with the opportunities of the valuable manipulations along with the magnitudes of the DQWW material parameters.

4.1 Introduction

The strong coupling of the low-dimensional excitons and localized photon states (exciton-polaritons) has been a subject of considerable interest for a long time in the optics of semiconductor nanostructures [1–3]. A number of unique physical phenomena were discovered (a Rabbi splitting, a large decrease of the exciton resonance lifetime [4, 5]) which can be directly aimed at promising applications in semiconductor optoelectronics. In this regard, intense theoretical efforts [6–10] have been carried

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out over the past time to realize the photonic structures such as planar and cylindrical multilayered microcavities, wide band-gap semiconductor photonic structures which offer unique optical properties.

To provide a substantial enhancement of light-matter interaction with the quantum-confined excitons and thus to support high polariton stability, the photonic structures should provide both the large exciton binding energy and the oscillator strength values [1, 2, 9]. Due to this a search for ways to enhance these physical values in photonic structures is a decisive factor in this area [11–15].

In the bulk semiconductors a role of Wannier-Mott excitons to optical response is insignificant compared to interband transitions since the exciton line dimensionless oscillator strength (per unit crystal cell) is the order of $(d_0/a_{0ex})^3 \approx 10^{-4} \div 10^{-5}$, where d_0 is the lattice constant and a_{0ex} is the excitons effective radii taking the values as $a_{0ex} \gg d_0$.

Such an estimate follows from the physical condition that the dimensionless oscillator strength is of the order of exciton radiative recombination probability (i.e. localization probability of electron and hole in the same crystal lattice unit cell) proportional to the $|\phi_{ex}(0)|^2$ value as compared to the intensity of the absorption spectrum of the semiconductor, where ϕ_{ex} is the exciton wave function. Here the localization probability is inversely proportional to the exciton effective volume $V_{ex} \sim a_{ex}^3$ and the absorption intensity is characterized by the localized Wannier functions linear combinations of band states just within the crystal unit cell volume V_0 and estimating as proportional to $V_0^{-1} \sim d_0^{-3}$.

In turn, the exciton effects in semiconductor low-dimensional structures-quantum wells (QW), quantum wires (QWW) are much more prominent and accessible for experimental detection than in bulk samples both in the absorption and in the emission. This, first of all, found the confident confirmation in experiments on optical absorption and photoluminescence [16, 17]. Theoretically, this is explaining by an increase in the exciton binding energy and in the oscillator strength of the corresponding exciton transition with a decrease in the spatial dimensions of the semiconductor quantum sample when takes place the compression of the exciton wave function in the spatial confinement direction (quantum confinement (QC)). In particular, in the thin QW there is a correspondingly fourfold and eightfold increase of the exciton binding energy and the exciton transition oscillator strength [18, 19].

These results refers the physical situation, when the dielectric constants mismatch between semiconductor quantum sample (semiconductor/dielectric quantum wells (DQW) and quantum wires (DQWW)) and surrounding dielectric barrier environment (ε_w, ε_b) is small or neglected. The difference between ε_w and ε_b leads to an unhomogeneous polarization of the quantum structure. In result for the case of $\varepsilon_w \gg \varepsilon_b$ and the distances between the charges as large as the QC spatial size (DQW thickness d or DQWW radius R) the charges produced field in the barrier begins to play a perceptible role and force electron and hole into the middle of the quantum sample and, so, to modify and enhance their interaction (dielectric confinement effect (DC)) [20–22]. As correspondingly are shown in Refs. [23, 24] and [22] the effective exciton volume in DQW and DQWW might radically reduced since in the direction normal to the DQW the exciton dimension is $\sim d$ and the in-plane effective radius

Fig. 4.1 The cylindrical DQWW of a radius R
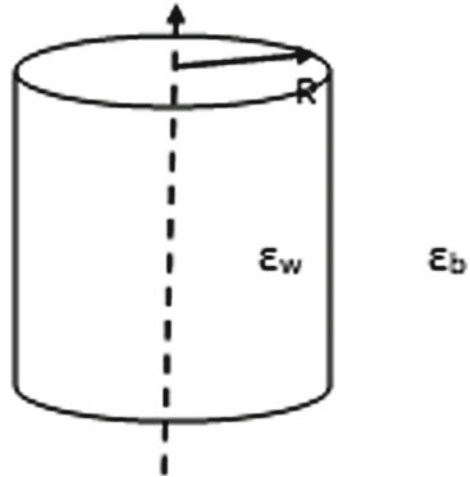
due to DC effect takes the value $\sim \sqrt{a0_{ex}d}$, while in the direction normal to the DQWW axes the exciton dimension is just $\sim R$ and in the direction parallel to axes the exciton effective radius becomes the value $(a_{0ex}R^2)^{1/3}$.

As follows, in addition to the existing enhancements from the QC effect, the DC effect in turn might lead to the further strong increase of the exciton binding energy and the exciton transition oscillator strength with decreasing of the DQW thickness d or DQWW radius R. The latter creates the more favorable conditions in these structures to provide a substantial enhancement of the light-matter interaction to provide a substantial enhancement of the light-matter interaction with the quantum-confined excitons and thus to support high polariton stability. In this connection in Ref. [11] the contribution of the excitonic transition to the electromagnetic response of the DQW and the possibilities of the propagation of the polariton waves are considered (Fig. 4.1).

The aim of the present paper is to develop an analogous formalism for the analysis of the particular features of the light-matter interaction with the one-dimensional (1D) DC enhanced dipole-allowed excitons in the DQWW with the cylindrical symmetry confinement. Specific azimuthal symmetry here substantially simplifies the convolution of the problem, as together with the 1D wave vector an additional quantum number (azimuthal) is available now.

The outline of presented paper is as follows. In Sect. 4.2, we describe the macroscopic theory of the electromagnetic response of the DQWW to the electromagnetic field with frequency ω and projection of wave vector $\boldsymbol{k}$ on the wire axes ($k_\parallel$). The next section is devoted to obtaining explicitly effective boundary conditions which are considering the dependence of the dielectric properties on the wire radius R. The Sect. 4.4 consists of an application of the obtained effective boundary conditions for the calculation of the light reflection coefficient from the structure under consideration.

4.2 Electromagnetic Linear Response to the Electromagnetic Field in the Macroscopically Homogeneous and Isotropic DQWW

We are considering an infinitely long cylindrical DQWW of a radius R filled by an active material with the dielectric constant ε_w and immersed in a dielectric barrier environment with the dielectric constant ε_b. Let take in the cylindrical polar coordinates (ρ, ϕ, z), where z axes coincides with the DQWW axes, and the plane of incidence (ρ, ϕ) be normal to the wire axes. Here a strong spatial confinement regime would be assumed presupposing that a quantum wire radius is small compared with the exciton Bohr radius $a_{0ex} = \varepsilon_w \hbar^2 / \mu^*_{ex} e^2$ for bulk samples ($a_{0ex} \gg R$), μ^*_{ex} is the exciton reduced mass. Then, as follows, the distances along the wire axes $|z| \gg R$ would be essential in discussed case and therefore the one-dimensional (1D) long wave region $k_\parallel \ll R^{-1}$ could be appropriate here. For this case, it is quite reasonable to introduce a spatially separated exciton wave function in the form

$$\Psi_{k_\parallel, M_e, M_h} (\boldsymbol{\rho_e}, \boldsymbol{\rho_h}, z_e, z_h) = \phi_l (z_e - z_h) \, \Upsilon_{j, M_e} (\boldsymbol{\rho_e}) \, \Upsilon_{j, M_h} (\boldsymbol{\rho_h}). \qquad (4.1)$$

Here $\phi(z)$ is the 1D wave envelope function of the exciton pair relative motion, $l = 1, 2, \ldots$ numbers the 1D exciton sublevels, $\Upsilon(\rho, \phi)$ is the one-particle wave function describing the electron or hole transverse relative to the DQWW axes motion and in the QC model of a square well with infinite walls is characterized by [22]

$$\Upsilon_{j, M} (\boldsymbol{\rho}, \phi) = \frac{1}{s^{1/2}} e^{i|M|\phi} J_{|M|} \left(\lambda_j^{|M|} \frac{\rho}{R} \right) / J_{|M|+1} \left(\lambda_j^{|M|} \right), \qquad (4.2)$$

where $M_{e(h)} = 0, 1, 2, \ldots$ is the one-particle angular momenta and $j = 1, 2, \ldots$ numbers the QWW subbands, $\lambda_j^{|M|}$ are the roots of $J_{|M|}(x) = 0$.

As shown in Ref. [22], for the DQWW satisfying the system of conditions $a_{0ex} \gg R \gg a_{0ex} / (\varepsilon_r \ln \varepsilon_r)^2$, where $\varepsilon_r = \varepsilon_w / \varepsilon_b$, the effective radii of the exciton ground and first excited bound states fall into the range of distances $\varepsilon_r (R/z)^2 \ln (|z|/R) \gg 1$ with $|z| \gg R$, where exciton interaction potential has the form

$$V(z) = -\frac{e^2}{\varepsilon_s R} \sqrt{\frac{\varepsilon_r \ln \varepsilon_r}{2}} \left[1 - \frac{|z|}{R} \sqrt{\frac{2}{\varepsilon_r \ln \varepsilon_r}} \right] \qquad (4.3)$$

and for that the 1D wave function of exciton pair relative motion and 1D exciton bound state energy spectrum have the form [25, 26]

$$\phi_l(z) = \frac{N(\mu_l)}{a_{ex}^{1/2}} \Phi \left(\frac{|z|}{a_{ex}} + \mu_l \right), l = 1, 2, 3, \ldots, \text{- even states}, \qquad (4.4)$$

$$\phi_{l+1}(z) = \frac{N(\xi_l)}{a_{ex}^{1/2}} \frac{|z|}{z} \Phi\left(\frac{|z|}{a_{ex}} + \xi_{l+1}\right), l = 1, 2, 3, \ldots, \text{- odd states,} \qquad (4.5)$$

$$E_l^{even(odd)} = -\frac{e^2}{\varepsilon_s R}\sqrt{\frac{\varepsilon_r \ln \varepsilon_r}{2}}\left[1 - \frac{\mu_l(\xi_{l+1})}{2^{1/3}}\sqrt{\frac{2}{\varepsilon_r \ln \varepsilon_r}}\frac{a_{0ex}}{R}\right], \qquad (4.6)$$

where $\Phi(z)$ is the Airy function, μ_l are the solutions of $\Phi'(z) = 0$, ξ_l are the solutions of $\Phi(z) = 0$, $a_{ex} = (a_{0ex}R^2/2)^{1/3}$, N is the normalization constant (for more details see [25, 26]). As follows from Eq. (4.6), the exciton binding energy of the DQWW is determined by the dielectric properties of the surrounding wire medium and increases as $\sim R^{-1}$.

The oscillator strength (per unit length of the DQWW) of the associated optical transition is

$$\frac{f_x}{L} = i\omega \frac{2m_0}{\hbar e^2} \mu_{cv} \left|\int \Psi_M(\rho_e = \rho_h, z_e = z_h)\, d\rho_e dz_e\right|^2, \qquad (4.7)$$

where μ_{cv} is the dipole matrix element in the bulk material, L is the wire length, m_0 is the electron free mass. In Eq. (4.7) only the $1D$ excitons with $M = M_e + M_h = 0$ are allowed in the dipole approximation because the others have zero exciton wave function at the origin of the relative coordinates. As follows from Eqs. (4.1), (4.2) and (4.3) the $1D$ exciton effective volume V_{ex} in DQWW decreases as $(R/a_{0ex})^{8/3}$ and since $|\Psi_M(0)|^2 \sim V_{ex}^{-1}$, so the $1D$ exciton line oscillator strength f_{ex} increases as $R^{-8/3}$.

The $1D$ excitonic transition contribution to the electromagnetic linear response of the cylindrical DQWW induces the dielectric polarization connected with the electric field E as

$$P_\alpha(\rho) = \int \chi_{\alpha\beta}(k_\parallel, \omega, \rho, \rho')\, E_{\beta k_\parallel \omega}(\rho')\, d\rho', \qquad (4.8)$$

where the nonlocal dielectric Fourier-transformed polarizability $\chi_{\alpha\beta}$ is given by

$$\chi_{\alpha\beta}(k_\parallel, \omega, \rho, \rho') = \chi_{\alpha\beta}(k_\parallel, \omega)\left[\Upsilon_{j,M}(\rho)\right]^2\left[\Upsilon_{j,M}(\rho')\right]^2, \qquad (4.9)$$

which will be calculated in accordance with the second order perturbation theory.

Here we assume that the light wave propagate in the plane orthogonal to the DQWW axes and consider only the exciton ground state in the wire which is characterized by $M = 0$ angular momentum. Given the cylindrical symmetry of the exciton ground state, the integral in Eq. (4.8) is nonzero only in the case of cylindrical light waves having zero angular momentum. Thus in the following we limit the discussion to that situation only.

After necessary actions with the excitonic wave functions (4.1), (4.2) and (4.4) for the Fourier-transformed dielectric polarization $P_\alpha(k_\parallel, \omega, \rho)$ obtain

$$P_\alpha \left(k_\parallel, \omega, \boldsymbol{\rho}\right) = \Lambda_{\alpha\beta\chi} \left(k_\parallel, \omega; \rho\right) \bar{E}_\beta, \tag{4.10}$$

where

$$\chi \left(k_\parallel, \omega; \boldsymbol{\rho}\right) = \chi \left(k_\parallel, \omega\right) \left(\frac{J_{|M|} \left(\lambda_j^{|M|} \frac{\rho}{R}\right)}{J_{|M|+1} \left(\lambda_j^{|M|}\right)}\right)^2, \tag{4.11}$$

and $\bar{E} = \int \left[\Upsilon_{j,M} \left(\boldsymbol{\rho}'\right)\right]^2 E_{k_\parallel \omega} \left(\boldsymbol{\rho}'\right) d\boldsymbol{\rho}'$.

In the Eq. (4.11)

$$\chi \left(k_\parallel, \omega\right) = \frac{\left(2\mu_{ex}^* e^8 \hbar^4\right)^{1/3} \left[N \left(\mu_l\right) \Phi \left(\mu_l\right)\right]^2}{\pi \varepsilon_s^{1/3} R^{8/3}} \frac{|V_{cv}|^2}{\varepsilon_{k_\parallel} \left(\varepsilon_{k_\parallel}^2 - (\hbar\omega)^2\right)}, \tag{4.12}$$

where V_{cv} is the matrix element of the velocity operator corresponding to the transition from the valence to conduction band, $\varepsilon_{k_\parallel}$ is the resonant energy of exciton creation, including its kinetic energy in the state with momentum $\hbar k_\parallel$.

In Eq. (4.10) the coefficients $\Lambda_{\alpha\beta}$, where $(\alpha, \beta) = (\rho, \phi, z)$, accounts the polarization structure of the exciton transition, i.e., the symmetry of wave functions of the c and v bands. As a rule, in assumption that the wire and barrier environments are optically isotropic, these coefficients contain just two characteristic terms: $\Lambda_z = \Lambda_\parallel$ and $\Lambda_\rho = \Lambda_\perp$.

Equation (4.11) exhibits clearly the strong increase of the excitonic oscillator strength in DQWW as R decreases. Except that Eqs. (4.10)–(4.11) expresses the nonlocal structure in **rho**, reflecting, as in the DQW [11–13], the peculiar "firmness" of the exciton transverse to wire axes wave function (4.1). Thereby the exciton creation probability defines by the field strength averaged over the electron and hole transverse to wire axes wave func-tions and the spatial distribution of the induced current itself is proportional to the same wave functions at $\rho_e = \rho_h = \rho'$.

4.3 Effective Boundary Conditions for the Electromagnetic Fields in the Macroscopically Homogeneous and Isotropic DQWW

The distribution of the electromagnetic field in the DQWW in presence of such polarization will be found in analogy with the well established method related the effective boundary conditions in QW systems [1, 11, 27] connected with the field magnitudes in the neighboring media and are followed from the Maxwell's equations. Here we expand this method on the DQWW system case in the first time.

Thereby in accordance with the Eqs. (4.8), (4.9) and (4.10) the corresponding electromagnetic field induction vector is given by

$$D_\alpha \left(k_\parallel, \omega; \rho\right) = \varepsilon_w E_{\alpha k_\parallel \omega} \left(\rho\right) + 4\pi \Lambda_{\alpha\beta} \chi \left(k_\parallel, \omega, \rho\right) \bar{E}_\beta, \qquad (4.13)$$

With the aim of to obtain the boundary conditions let now integrate the Maxwell equations and take into consideration only the leading terms in the small quantities $k_\parallel R$ and $(\omega/c)R$ together with the Eq. (4.13). In these equations we are dealing with the monochromatic fields $\boldsymbol{E}$, $\boldsymbol{D}$ and $\boldsymbol{H}$ vary in space according to the law

$$\boldsymbol{A}\left(\boldsymbol{r}, t\right) = \boldsymbol{A}\left(\rho\right) e^{i(k_\parallel z - \omega t)} = \sum_{|M|} \boldsymbol{A}\left(\rho\right) e^{i|M|\phi} e^{i(k_\parallel z - \omega t)}. \qquad (4.14)$$

The resulting solutions will be matched in the barrier and QWW regions and the latter will be occurred only by these boundary conditions. Below, using these boundary conditions, we will analyze the reflection coefficient from the DQWW structure.

Let at first start from the integrating of the Maxwell equation $\text{div}\boldsymbol{D} = 0$. The latter in the polar planar coordinates (ρ, ϕ) for the case of exciton ground state with the condition $M = 0$ takes the form

$$\boldsymbol{e}_\rho \int_0^R \int_0^{2\pi} \frac{1}{\rho} \frac{\partial}{\partial \rho} \left(\rho D_\rho^{(0)}\left(\rho\right) e^{ik_\parallel z}\right) \rho d\rho d\phi + \boldsymbol{e}_z \int_0^R \int_0^{2\pi} \frac{\partial}{\partial z} \left(D_z^{(0)}\left(\rho\right) e^{ik_\parallel z}\right) \rho d\rho d\phi = 0, \qquad (4.15)$$

which after integrating together with Eq. (4.13) becomes

$$\boldsymbol{E}_\rho^{(0)}\left(R\right) = -\frac{iR}{\varepsilon_w} \varepsilon_\parallel \left(k_\parallel \bar{E}_z^{(0)}\left(R\right)\right) \boldsymbol{e}_\rho, \qquad (4.16)$$

where $\bar{E}_{1z}^{(0)}\left(R\right) = \int_0^R E_z^{(0)}\left(\rho\right) \rho d\rho$ and

$$\varepsilon_\parallel = \varepsilon_w + \frac{4\pi \Lambda_\parallel \chi \left(k_\parallel, \omega\right)}{J_1^2 \left(\lambda_1^0\right)}, \qquad (4.17)$$

In Eq. (4.15) the term associated with an angular variation of the electric field components is obviously omitted. Here Eq. (4.16) links the corresponding components of the electric field.

To receive the complete set of these components let now deal with the equation $\text{rot}\boldsymbol{E} = \frac{i\omega}{c}\boldsymbol{H}$ as well. Analogous to the Eq. (4.15) for this case we have

$$-\boldsymbol{e}_\rho \int_0^R \int_0^{2\pi} \frac{\partial E_\phi^{(0)}}{\partial z} \rho d\rho d\phi + \boldsymbol{e}_\phi \int_0^R \int_0^{2\pi} \left(\frac{\partial E_\rho^{(0)}}{\partial z} - \frac{\partial E_z^{(0)}}{\partial \rho}\right) \rho d\rho d\phi + \boldsymbol{e}_z \int_0^R \int_0^{2\pi} \frac{1}{\rho} \frac{\partial(\rho E_\phi^{(0)})}{\partial \rho} \rho d\rho d\phi, \qquad (4.18)$$

where $\boldsymbol{e}_\rho$, $\boldsymbol{e}_\phi$ and $\boldsymbol{e}_z$ are the polar directional unit vectors.

Let now multiply Eq. (4.18) vectorially by $\boldsymbol{e}_\rho$ and make use of the Eq. (4.14) then we receive

$$E_\phi^{(0)} e_\phi + E_z^{(0)} e_z = \frac{i\omega R}{c} \left[\bar{H}^{(0)}(R) e_\rho \right] + \bar{E}_{2z}^{(0)} e_z + \frac{ik_\parallel R}{2\varepsilon_w} \bar{D}_\rho^{(0)} \frac{\varepsilon_w + 8\pi\Lambda_\perp\chi \left[J_1^{-2} - J_1^{-4} \right]}{\varepsilon_w + 8\pi\Lambda_\perp\chi} e_z,$$
$$(4.19)$$

where $\quad \bar{H}^{(0)}(R) = \int_0^R H^{(0)}(\rho)\,\rho d\rho, \quad \bar{E}_{2z}^{(0)}(R) = \int_0^R E_z^{(0)}(\rho)\,d\rho, \quad \bar{D}_\rho^{(0)}(R) = \int_0^R D_\rho^{(0)}(\rho)\,\rho d\rho.$

If combining the Eqs. (4.16) and (4.19) we find the electric field boundary conditions as

$$\boldsymbol{E}(R) = -\frac{iR}{\varepsilon_w}\varepsilon_\parallel \left(k_\parallel \bar{E}_{1z}^{(0)}(R) \right) e_\rho - \frac{i\omega R}{c} \left[\bar{H}^{(0)} e_\rho \right] + \bar{E}_{2z}^{(0)} e_z + \frac{ik_\parallel R}{\varepsilon_s} \bar{D}_\rho \frac{\varepsilon_w + 8\pi\Lambda_\perp\chi \left[J_1^{-2} - J_1^{-4} \right]}{\varepsilon_w + 8\pi\Lambda_\perp\chi} e_z,$$
$$(4.20)$$

From the magnetic field equations $\text{div}\,\boldsymbol{D} = 0$, $\text{rot}\,\boldsymbol{D} = -\frac{i\omega}{c}\boldsymbol{D}$ the similar manipulations produce finally the corresponding boundary conditions as

$$\boldsymbol{H}(R) = -iR \left(k_\parallel \bar{H}_z^{(0)}(R) \right) e_\rho - \frac{i\omega R}{c}\varepsilon_\parallel \left[\bar{\boldsymbol{E}}^{(0)}(R) e_\rho \right] + \bar{H}_z^{(0)} e_z + iR \left(k_\parallel H_\rho^{(0)} \right) e_z,$$
$$(4.21)$$

where $\quad \bar{H}_z^{(0)}(R) = \int_0^R H_z^{(0)}(\rho)\,\rho d\rho, \quad \bar{\boldsymbol{E}}^{(0)}(R) = \int_0^R \boldsymbol{E}^{(0)}(\rho)\,\rho d\rho, \quad \bar{H}_\rho^{(0)}(R) = \int_0^R H_\rho^{(0)}(\rho)\,\rho d\rho.$

In Eqs. (4.20) and (4.21) $\boldsymbol{E}(R)$ and $\boldsymbol{H}(R)$ are the boundary values of the electric and magnetic fields in the barrier region, c is the light velocity. Under boundary conditions after Exps. (4.20) and (4.21), a presence of a DQWW, as already emphasized, is taken into account up to the first-order terms in small parameters $\sim k_\parallel R$ and $\omega R/c$.

At the same time due to the Eq. (4.17) the first terms of the right hand sides in the Eqs. (4.20) and (4.21), i.e. the terms $\sim \varepsilon_\parallel \frac{\omega R}{c}$, would be hold only in the discussions. In turn, the last term in Eq. (4.20) will be taken into account in the narrow frequency range with $\varepsilon_w + 8\pi\Lambda_\perp\chi \ll 1$.

4.4 The Scattering Coefficient of the Light from the Macroscopically Homogeneous and Isotropic DQWW

In this section let calculate the scattering coefficient r_{DQWW} of electromagnetic light wave from the discussed DQWW structure based on the effective boundary conditions after Eqs. (4.20) and (4.21).

As we admitted in Sect. 4.2 the electromagnetic waves of a cylindrical symmetry propagate in the plane orthogonal to the DQWW axes (z axes) and along the latter is directed the electric-field vector (p-polarized light). With that we consider only the exciton ground state in the quantum wire, just characterized by zero angular momentum and are interested in results close to the exciton resonance. We note that in the case with $M = 0$ the excitonic polarization lies along the DQWW axes and the mode under consideration has longitudinal nature. Accordingly, as in the [15],

we consider the exciton resonant modes appearing as resonances of the Breit-Wigner type in the r_{DQWW} of the barrier waves.

Let establish the scattering coefficient r_{DQWW} as the ratio of the amplitudes of outgoing and incoming cylindrical waves in the DQWW barrier region at $\rho = R$ for which the p-polarized electromagnetic field correspondingly has the forms [15] outgoing and incoming cylindrical waves in the DQWW barrier region at $\rho = R$ for outgoing and incoming cylindrical waves in the DQWW barrier region at $\rho = R$ for which the p-polarized electromagnetic field correspondingly has the forms [15]

$$\begin{cases} E_z^{inc} = H_0^{(1)}(k\rho) = J_0(k\rho) + i N_0(k\rho) \\ H_\phi^{inc} = i\sqrt{\varepsilon_b}\frac{k}{k_0}\left(H_0^{(1)'}(k\rho)\right) = i\sqrt{\varepsilon_b}\frac{k}{k_0}\left(J_0'(k\rho) + i N_0'(k\rho)\right) \end{cases}, \qquad (4.22)$$

$$\begin{cases} E_z^{out} = H_0^{(2)}(k\rho) = J_0(k\rho) - i N_0(k\rho) \\ H_\phi^{out} = i\sqrt{\varepsilon_b}\frac{k}{k_0}\left(H_0^{(2)'}(k\rho)\right) = i\sqrt{\varepsilon_b}\frac{k}{k_0}\left(J_0'(k\rho) - i N_0'(k\rho)\right) \end{cases}, \qquad (4.23)$$

where $= \sqrt{k_0^2 - k_\parallel^2}$, $k_0 = \sqrt{\varepsilon_b}\frac{\omega}{c}$, $J_0(x)$ and $N_0(x)$ are the zeroth-order regular and singular Bessel functions, respectively, $J_0'(x) = -J_1(x)$ and $N_0'(x) = -N_1(x)$ are the first derivatives of these functions.

For this type of barrier resonant modes when taking into account the effective boundary conditions after Eqs. (4.20) and (4.21) we have

$$\begin{cases} E_z(k\rho)\,|_{\rho=R^+} = E_{2z}^{(0)}(k\rho)\,|_{\rho=R^-} \\ H_\phi(k\rho)\,|_{\rho=R^+} = -\frac{i\omega R}{c}\varepsilon_\parallel \gamma E_{2z}^{(0)}(k\rho)\,|_{\rho=R^-} \end{cases}, \qquad (4.24)$$

where

$$E_z(k\rho)\,|_{\rho=R^+} = E_z^{inc}(k\rho)\,|_{\rho=R^+} + r_{DQWW} E_z^{out}(k\rho)\,|_{\rho=R^+}, \qquad (4.25)$$

$$H_\phi(k\rho)\,|_{\rho=R^+} = H_\phi^{inc}(k\rho)\,|_{\rho=R^+} + r_{DQWW} H_\phi^{out}(k\rho)\,|_{\rho=R^+} \qquad (4.26)$$

and $\gamma = \bar{E}_{1z}^{(0)}(R) / \bar{E}_{2z}^{(0)}(R) \sim 1$.

After substituting the Eqs. (4.22), (4.23), (4.25) and (4.26) in Eq. (4.24) we get for the scattering coefficient r_{DQWW} as

$$r_{DQWW} = \frac{[\Lambda N_0(kR) - N_1(kR)] - i[\Lambda J_0(kR) - J_1(kR)]}{[\Lambda N_0(kR) - N_1(kR)] + i[\Lambda J_0(kR) - J_1(kR)]}, \qquad (4.27)$$

where $\Lambda = \frac{\omega R}{c\sqrt{\varepsilon_b}}\varepsilon_\parallel \gamma$ and $\varepsilon_\parallel$ is defined after Eq. (4.17).

Equation (4.27) presents the scattering coefficient of p-polarized light from the DQWW structure at normal to the wire axes incidence near $\omega = \omega_{ex} = \varepsilon_{k_\parallel} \hbar^{-1}$ and strongly depends from the DQWW material parameters such as: (a) quantum wire radius R and dielectric constant ε_w, (b) barrier dielectric constant ε_b.

Since here we are dealing with an accuracy of small parameters $\sim k_\parallel R$ and $\omega R/c$ then Eq. (4.27) will take the following limiting form

$$r_{DQWW} = \frac{(2/\pi)\left[\Lambda \ln(kR) + (kR)^{-1}\right] - i\Lambda}{(2/\pi)\left[\Lambda \ln(kR) + (kR)^{-1}\right] + i\Lambda}, \qquad (4.28)$$

where the asymptotic small parameter limits of the Bessel functions are taken into account. If now substitute corresponding expressions for Λ, $\varepsilon_\parallel$ and χ parameters in Eq. (4.28) then after simple transformations we receive as in Ref. [15] a Breit-Wigner type final form for the scattering coefficient of the p-polarized light

$$r_{DQWW} = \frac{\omega - \widetilde{\omega}_p\left(k_\parallel, \omega_{ex}\right) - i\Gamma\left(k_\parallel, \omega_{ex}\right)}{\omega - \widetilde{\omega}_p\left(k_\parallel, \omega_{ex}\right) + i\Gamma\left(k_\parallel, \omega_{ex}\right)}, \qquad (4.29)$$

with

$$\widetilde{\omega}_p\left(k_\parallel, \omega_{ex}\right) = \omega_{ex} + \frac{\alpha(k_0 R)^2 \ln(kR)}{\varepsilon_b + \alpha\varepsilon_w(k_0 R)^2 \ln(kR)}\Omega_{ex}, \qquad (4.30)$$

where

$$\Omega_{ex} = \frac{(2\mu_{ex}^* e^8)^{1/3}\left[N\left(\mu_l\right)\Phi\left(\mu_l\right)\right]^2 |V_{cv}|^2}{\pi\varepsilon_w^{1/3}\hbar^5\omega_{ex}^6 R^{8/3}}. \qquad (4.31)$$

Here the dispersion law for the DQWW structure exciton-polaritons and their ra-diative broadening would be found from the following expressions

$$\omega = \widetilde{\omega}_p\left(k_\parallel, \omega_{ex}\right) \qquad (4.32)$$

and

$$\Gamma\left(k_\parallel, \omega_{ex}\right) = \frac{\alpha\varepsilon_b(k_0 R)^2 \ln(kR)}{(\varepsilon_b + \alpha\varepsilon_w(k_0 R)^2 \ln(kR))(\varepsilon_b + \alpha(k_0 R)^2 \ln(kR))}\Omega_{ex}. \qquad (4.33)$$

The Eqs. (4.30)–(4.33) made it possible to obtain the dispersion spectrum and life-time for the $1D$ exciton-polaritons of the DQWW near the exciton resonance with the opportunities of the valuable manipulations along with the magnitudes of the DQWW material parameters.

In conclusion, following the typical procedure for the optics of $2D$ system [1, 2, 27], we have developed a theory to solve Maxwell's equations together with the non-local $1D$ excitonic response to the dielectric polarization of the DQWW in the presence of the dielectric mismatch effect. This permits to calculate the dispersion spectrum of 1D exciton-polaritons of the DQWW near the excitonic resonance.

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