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# Optical properties of near-surface exciton quantum wells

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An overview of theoretical investigations on near-surface semiconductor quantum wells, whose optical properties are considerably affected by the interaction of the exciton with the sample surface is given. Near-surface quantum wells with both weak and strong quantum confinement of excitons are considered. When the exciton quantum well is very close to the sample surface, exciton dynamics is determined not only by characteristics of the quantum well, but also by the interaction of the exciton with the sample surface. Optical spectra for various systems of near-surface quantum wells are discussed. In addition, the optical manifestation of magnetoexcitons in near-surface quantum wells is also commented.

*Keywords*: Exciton quantum wells; Quantum well surface optical properties; Magnetoexcitons

## 1. Introduction

Excitons in confined systems, such as quantum wells (QWs), have attracted the attention of researches because of their potential applications to optoelectronic devices. There exist two qualitatively different regimes of exciton confinement in quantum wells: weak and strong confinement regimes. In the regime of weak confinement of exciton (also known as thin film regime), the effective width $d$ of the quantum well is much larger than the exciton radius $a_0$ ( $d >> a_0$ ) and, therefore, the center-of-mass motion of the exciton is quantized. In this regime the relative motion of the electron-hole pair is practically the same as in the bulk except for a small distortion near film boundaries [1,2]. Because of the interaction of the exciton with the interfaces, transition (exciton-free) layers appear near film boundaries. Therefore, the confinement of the exciton center-of-mass motion occurs in an effective length $d_{\text{eff}}$ smaller than the actual film thickness ( $d_{\text{eff}} \approx d - 2\,l$, $l$ is the effective size of one exciton-free layer inside the film ). The quantization of the exciton center-of-mass in thin films manifests itself as a resonance structure in optical spectra (reflectivity, transmission, and absorption). It is interesting that this resonance structure can also be explained as an effect of multimode interference of exciton polaritons [2-4].

In the regime of strong exciton confinement, the thickness of the quantum well $d$ is smaller than the exciton radius $a_0$ ( $d \le a_0$ ) and the motion of the electron and hole is quantized separately in the direction perpendicular to the quantum-well plane. Due to the in-plane Coulomb interaction, the electron-hole pair in a quantum well forms a quasi-2D exciton. The optical response of excitons in quantum wells has been investigated in numerous works (see, for example, [5-7] ). Quasi-2D excitons in quantum wells couple to light and produce resonances in reflectivity spectra, which correspond to exciton-polariton modes with polarizations perpendicular to the plane of incidence (T mode), parallel to the in-plane component $k\|$ of the wave vector for the incident light (L mode), and perpendicular to the quantum-well plane (Z mode).

Semiconductor thin films and quantum wells are commonly grown on substrates and may have a cap layer, overlying them. Therefore, their optical properties depend on the characteristics of both the substrate and the cap layer. If the thickness of the cap layer is sufficiently small (of the order of the exciton radius), the optical properties of the semiconductor structure may be affected considerably by the interaction of the exciton with the sample surface. In this paper we present an overview of the theoretical approaches (Secs. 2-4) applied for interpreting optical properties of near-surface exciton quantum wells in both weak (Sec. 3) and strong (Sec. 4) confinement regimes. Also, we discuss about optical techniques (reflectivity, 45 degrees reflectometry, diffuse reflection, for example) which are quite appropriate to characterize that kind of confined systems.

## 2. General formulation of the problem

The optical properties of excitons in near-surface quantum wells (quantum wells on substrate) can be described by employing the density matrix approach of Stahl and Balslev [8], which leads to the set of differential equations for the coherent-amplitude vector $\vec{Y}(\vec{r}_e, \vec{r}_e)$ of electron and hole with coordinates $r_e$ and $\vec{r}_h$, respectively:

$$\left(H_{vc} - \hbar(w + in)\right)\vec{Y} = \hat{M}_{vc}(\vec{r})\vec{E}(\vec{R}), \qquad (1)$$

where $\vec{E}$ is the electric field of the electromagnetic wave of frequency ω, $\hat{M}_{vc}$ is the interband-transition dipole density, ν is a phenomelogical damping coefficient, $H_{vc}$ is a two-band Hamiltonian with gap $E_g$, which is given by

$$H_{vc} = E_g - \frac{\hbar^2}{2m_e}\nabla_e^2 - \frac{\hbar^2}{2m_h}\nabla_h^2 + V_e(\vec{r}_e) + V_h(\vec{r}_h) +$$

$$+V_c(r) + V_s(\vec{r}_e, \vec{r}_h). \qquad (2)$$

Here, $V_c(r) = -e^2/e_s r$ is the Coulomb potential, $r$ is the magnitude of the vector $\vec{r} = \vec{r}_e - \vec{r}_h = (\vec{r}, z)$, $\vec{R} = (m_e\vec{r}_e + m_h\vec{r}_h)/(m_e + m_h)$ is the radius vector of the exciton center-of-mass, $\varepsilon_s$ is the low-frequency dielectric constant of the medium, $m_e$ and $m_h$ are respectively the electron and hole effective masses. The quantities $V_e(\vec{r}_e)$ and $V_h(\vec{r}_h)$ are the confining potentials for the electron and the hole. In the case of QW geometry these confining potentials have steplike form in the *z* direction. Finally, the quantity $V_s(r_e, r_h)$ denotes the potential due to the surface. An intrinsic contribution to $V_s(\vec{r}_e, \vec{r}_h)$ is the image potential, $V_{im}(r_e, r_h)$, which appears because of the dielectric mismatch at the sample surface (*Z*=0):

$$V_{im}(r_e, r_h) = \frac{e^2}{e_s}\frac{e_s - 1}{e_s + 1}\left[\frac{1}{4}\left(\frac{1}{z_e} + \frac{1}{z_h}\right) - \frac{1}{\sqrt{r^2 + (z_e + z_h)^2}}\right] \quad (3)$$

The polarization of the medium is related to the coherent amplitudes as

$$\vec{P}(\vec{R}) = 2\int \hat{M}_{vc}(\vec{r})\vec{Y}(\vec{R}, \vec{r})d\vec{r}\,. \quad (4)$$

In order to calculate optical functions of near-surface quantum wells, the system of equations (1)-(4) together with Maxwell equations,

$$\nabla\times\nabla\times\vec{E} - e_\infty\frac{w^2}{c^2}\vec{E} = 4p\frac{w^2}{c^2}\vec{P}\ , \quad (5)$$

should be solved. The quantity $\varepsilon_\infty$ in Eq. (5) stands for the high-frequency dielectric constant. In solving Eqs. (1)-(5), the boundary condition $\vec{Y} = 0$, when the electron or the hole is at the vacuum-QW interface ($z_e$=0 or $z_h$=0 ), as well as the condition of continuity for tangential components of the electric and magnetic fields of the electromagnetic wave are applied.

## 3. Weak confinement of exciton

### *3.1. Thin films on substrate*

Several theories (see, for example, Refs. [1-4, 9-13]) have been developed to describe the optical manifestation of the quantization of the exciton center-of-mass in thin semiconductor films. The main problem encountered here is associated with the complicated behavior of the exciton near film boundaries because of its interaction with them. The exciton-interface interaction leads to the coupling of the relative and translational motions of the electron-hole pair. One of the mostly applied models to interpret optical spectra in terms of the concept of the exciton center-of-mass quantization is the so-called adiabatic approach [9]. Within this model, the coherent amplitude $\vec{Y}(\vec{r}_e, \vec{r}_h)$ is expanded as

$$\vec{Y}(\vec{r}_e, \vec{r}_h) = \sum_l \vec{Y}_l(\vec{R})j_l(Z, \vec{r}), \quad (6)$$

where $l = \{n, l, m\}$ is the complete set of quantum numbers, $j_l(Z, \vec{r})$ are solutions of the Schrödinger equation given by

$$\left[-\frac{\hbar^2}{2m}\nabla_r^2 - \frac{e^2}{e_s r} + V_s(Z, \vec{r})\right]j_l(Z, \vec{r}) =$$

$$= E_l^r(Z)j_l(Z, \vec{r}) \quad (7)$$

Here, the eigenfunctions $j_l(Z, \vec{r})$ and the eigenenergies $E_l^r(Z)$ associated with the relative motion depend parametrically on the center-of-mass coordinate *Z*. According to the adiabatic approximation, the expansion coefficients $\vec{Y}_l(\vec{R})$ in Eq. (6) satisfy the equation:

$$\left\{E_g + E_l^r(Z) - \hbar(w + in) - \frac{\hbar^2}{2M}\nabla^2\right\}\vec{Y}_l(\vec{R}) =$$

$$= \langle\hat{M}_{vc}\rangle_l^*\vec{E}(\vec{R}), \quad (8)$$

where

$$\langle\hat{M}_{vc}\rangle_l^* = \int d\vec{r}\, j_l^*(Z, \vec{r})\hat{M}_{vc}(\vec{r})\,. \quad (9)$$

A simple model for the transition dipole density is the shell model [8] which is defined as

$$\hat{M}_{vc}(\vec{r}) = \frac{M_0}{4pr^2}d(r - r_0)\,, \quad (10)$$

with $r_0 \to 0$. Hence, only *s*-excitons couple to light. Sufficiently far from the surface $(V_s \to 0)$ the energy $E_l^r \approx -E_{b,l}^{3D}$, where $E_{b,l}^{3D}$ is the binding energy for the λ-sate 3D exciton. It is convenient to introduce the quantity $U_l(Z) = E_l^r(Z) + E_{b,l}^{3D}$ which plays the role of an effective surface potential in Eq. (8) for the exciton translational motion. Finally, the excitonic polarization can be written as

$$\vec{P}(\vec{R}) = 2\sum_l\langle\hat{M}\rangle_l\vec{Y}_l(\vec{R})\,. \quad (11)$$

Below, we shall consider only the contribution of the ground-state exciton (λ={1,0,0}) to the polarization $\vec{P}$ [Eq.

(11)]. For simplicity, the sub-index λ={1,0,0}, indicating the exciton ground state, will be omitted. So, in this case the polarization $\vec{P}(\vec{R}) = 2\langle\hat{M}\rangle\vec{Y}(\vec{R})$.

In high-quality semiconductors, the no-escape boundary condition for the electron and hole at the surface and the image potential make the exciton surface potential $U(Z)$ be repulsive [9,14]. This intrinsic potential decays towards the interior of the film over distances of the order of the exciton radius. The simplest model for an intrinsic potential is an infinite barrier at certain distance *l* from the surface of the order of $a_0$. Therefore, in films of III-V and II-VI compounds, which are characterized by a relatively large exciton radius $a_0 \sim$ 3-13 nm, the size *l* of exciton-free layers reduces notably the effective thickness of the nonlocal film [1-5]. Besides, the resonance structure of the optical spectra of this kind of materials turns out to be very complicated due to the quantization of the center-of-mass of both heavy-hole (hh) and light-hole excitons [5]. The study of exciton center-of-mass quantization is simpler in CuCl thin films [11-13] which have a large binding energy $E_b^{3D} \sim$ 190 meV and a very small exciton raidus $a_0 \sim$ 7 angstroms. Hence, CuCl thin films allow to maintain the bulk-like behavior up to a very small film thickness ( $d \sim$ 7 nm $>> a_0$ ). In the case of CuCl, the presence of exciton-free layers near film boundaries can be neglected in interpreting its optical spectra. So, inside the whole CuCl film ( $0 < Z < d$ ) one can assume that $E^r(Z) = -E_b^{3D}$ (i.e. $U(Z) = 0$ ) in Eq. (8). Then, the exciton-surface interaction is described by using only the boundary condition $\vec{P} = 0$ (or $\vec{Y} = 0$ ) at *Z*=0 and *Z*=*d*.

In the recent work [15], spectra of *s*- and *p*-polarized reflectivity ($R_p$ , $R_s$) at 45-degrees angle of incidence, and spectra of 45-degrees reflectometry [16] ( $\Delta_{45} = R_p - R_s^2$ ) for a CuCl film (*d*= 16.5 nm) on a MgO substrate were calculated by using a nonlocal dielectric function,

$$e(\vec{k},w) = e_\infty + 4pc(\vec{k},w), \qquad (12)$$

$$c(\vec{k},w) = \frac{4w_T}{\hbar}\frac{M_0^2 \mid j_{100}(r=0)\mid^2}{w_T^2 + \frac{\hbar w_T}{M}k^2 - w^2 - 2inw}, \qquad (13)$$

where

$$w_T = (E_g - E_b^{3D})/\hbar. \qquad (14)$$

The spatially dispersive dielectric function [Eq. (12)] is straightforwardly obtained from the system of equations (5), (8)-(11) with $U(Z) = 0$ and polaritonic fields, $\vec{P}(\vec{R})$ and $\vec{E}(\vec{R})$, proportional to $\exp[i\vec{k}\cdot\vec{R}]$, and assuming $\mid w - w_T \mid << w_T$ (a similar derivation of the dielectric function is presented in Ref. [17]).

The reflectivity $R_s$, calculated in Ref. [15], exhibits a rich resonance structure associated with size-quantized transverse exciton polaritons. The *z* component of the wave vector for these polaritonic modes satisfies the Fabry-Perot relation Re $k_z d = np$ ( *n* is an integer ) at frequencies very close to the eigenvalues

$$\hbar w_{Tn} = E_g - E_b^{3D} + \left(\frac{\hbar^2 p^2 n^2}{2Md^2}\right) \qquad (15)$$

of the quantized transverse exciton in the thin film. In the *p*-polarized geometry longitudinal polariton modes, besides transverse modes, are also excited. The frequencies corresponding to the size quantization of the longitudinal polaritonic modes are

$$w_{Ln} = w_{Tn} + w_{LT}, \qquad (16)$$

where $w_{LT}$ is the longitudinal-transverse splitting. The longitudinal resonances are not easily detected in the spectrum of *p*-polarization reflectivity, $R_p$ , because of the relatively large damping at frequencies above the longitudinal one ($\omega > \omega_L$ )[11-13] and their interference with transverse resonances. In spectra $\Delta_{45} = R_p - R_s^2$ of 45-degrees reflectometry, transverse resonances are not so pronounced as in the simple reflectivities $R_p$ and $R_s$ [15]. This fact permits to observe the longitudinal resonances (dips) clearly in $\Delta_{45}$ spectra.

### *3.2. Extrinsic transition layers*

The center-of-mass quantization of excitons can also occur in extrinsic transition layers of bulk semiconductors [9,16,18-23] and thin films [4,24,25]. Indeed, the form of the surface potential $U(Z)$ can be altered either by surface treatments such as electron and ion bombardment, intense illumination, heating, doping, and a bias electric field, or unintentionally during the process of crystal growth. The extrinsic transition layers are characterized by space charge that produces an inhomogeneous macroscopic electric field $\vec{E}_{in}(Z)$ with which excitons interact. Hence, the potential $V_s(\vec{r}_e, \vec{r}_h)$ in Eq. (2), describing the interaction of the electron-hole pair with the surface, is given by

$$V_s(\vec{r}_e, \vec{r}_h) = V_{im}(Z, \vec{r}) + e\vec{E}_{in}(Z)\cdot\vec{r}. \qquad (17)$$

Far enough from the surface the mechanism of this interaction is related with the quadratic Stark effect. It introduces a coordinate-dependent shift of the exciton binding energy, and consequently, an extrinsic contribution to the excitonic surface potential $U(Z)$. This contribution can be attractive, unlike intrinsic potential, and a potential well, where the exciton center-of-mass motion is quantized,

can appear. The confinement of excitons in extrinsic transitions layers and its manifestation in optical spectra have been widely investigated in the last two decades [4,9,16,18-25]. Now, it is well established that excitons, confined in an extrinsic potential well, couple to light and produce transverse resonances (broad peaks) in reflectivity spectra at frequencies very close to the eigenenergies of the exciton bound states. Besides transverse resonances, longitudinal resonances (dips), associated with quantized longitudinal polarization waves within the extrinsic surface potential well, can be observed in spectra of *p*-polarization reflectivity [21,23,26] and, particularly, of 45° reflectometry [16,22]. Near-surface localized excitons and their corresponding longitudinal polarization waves are also manifest in frequency and angle dependencies of the cross section for light scattering by a randomly rough surface [27,28]. These dependencies are very sensitive to the degree of correlation between surface roughness and potential-well fluctuations [28].

## 4. Strong confinement of exciton

### *4.1 Quasi-2D quantum wells*

If the quantum-well width is smaller than the exciton radius, then the Coulomb attraction between the electron and hole is suppressed by the confining potentials in the direction perpendicular to the well plane. In this case the Coulomb potential $V_c(r)$ in the Hamiltonian (2) can be written, to good approximation, as

$$V_c \approx -\frac{e^2}{e_s r}. \qquad (18)$$

Hence, the coherent amplitude can be expanded as

$$\vec{Y}(\vec{r}_e,\vec{r}_h) = \sum_l \vec{Y}_l(\vec{R}_\parallel) j_{nm}(\vec{r}) y_p^{(e)}(z_e) y_q^{(h)}(z_h), \qquad (19)$$

where $l=\{n,m;p,q\}$ is the complete set of quantum numbers, $j_{nm}(\vec{r})$ are the eigenfunctions for the relative electron-hole motion in the 2D Coulomb potential:

$$\left[-\frac{\hbar^2}{2m}\nabla_r^2 - \frac{e^2}{e_s r}\right] j_{nm}(\vec{r}) = E_n^{2D} j_{nm}(\vec{r}). \qquad (20)$$

Functions $y_p^{(e)}(z_e)$ and $y_q^{(h)}(z_h)$ in Eq. (19) are the confinement wave functions and satisfy respectively the equations:

$$\left[-\frac{\hbar^2}{2m_e}\frac{\partial^2}{\partial z_e^2} + V_e(z_e)\right] y_p^{(e)}(z_e) = E_p^e y_p^{(e)}(z_e), \qquad (21)$$

$$\left[-\frac{\hbar^2}{2m_h}\frac{\partial^2}{\partial z_h^2} + V_h(z_h)\right] y_q^{(h)}(z_h) = E_q^h y_q^{(h)}(z_h). \qquad (22)$$

After substituting the expansion for the coherent amplitude (19) into Eq. (1), we obtain the equation for the expansion coefficients $\vec{Y}_l(\vec{R}_\parallel)$ in the form

$$\left[E_g + E_n^{2D} + E_p^e + E_q^h - \hbar(w+in) - \frac{\hbar^2}{2M}\nabla_{R_\parallel}^2\right]\vec{Y}_l(\vec{R}_\parallel) =$$

$$= \int dz_e \int dz_h \int d\vec{r}\left[j_{nm}(\vec{r}) y_p^{(e)}(z_e) y_q^{(h)}(z_h)\right]^* \times$$

$$\times \hat{M}(\vec{r})\vec{E}(\vec{R}). \qquad (23)$$

The interband-transition dipole density is well described by the shell model [29,30]:

$$\hat{M}_{vc}(\vec{r}) = \frac{\hat{M}_0}{2pr} d(r - r_0) d(z), \qquad (24)$$

with $r_0 \to 0$. Hence, only excitons with $m=0$ couple to light and the excitonic polarization can be expressed as

$$\vec{P}(\vec{R}) = \hat{M}_0 \sum_l d_{m0} j_{n0}(0) y_p^{(e)}(Z) y_q^{(h)}(Z) \vec{Y}_l(\vec{R}_\parallel), \qquad (25)$$

where $d_{mm'}$ is the Kronecker's symbol. The solution of the system of equations (5), (23) and (25) allows to calculate optical functions for near-surface quasi-2D quantum wells. Since the component parallel to the quantum-well plane, $\vec{k}_\parallel$, of the wave vector of the incident light is conserved, the polaritonic fields $\vec{P}$, $\vec{E}$, as well as the coherent amplitude $\vec{Y}$ are proportional to $\exp[i\vec{k}_\parallel \cdot \vec{R}_\parallel]$. Considering only the ground-state exciton (λ={0,0;1,1}), the polarization field can be expressed as

$$\vec{P}(Z) = \int dZ' \hat{c}(Z,Z')\vec{E}(Z'), \qquad (26)$$

where the nonlocal susceptibility is given by

$$\hat{c}(Z,Z') = \hat{c}(w)\Phi(Z)\Phi^*(Z'), \qquad (27)$$

where

$$\Phi(Z) = y_1^{(e)}(Z) y_1^{(h)}(Z), \tag{28}$$

$$\hat{c}(w) = \hat{M}_0^2 \frac{2}{\hbar} \frac{|j_{00}(0)|^2}{w_0(k_\parallel) - w - in}. \tag{29}$$

Here, the resonance frequency $w_0$ is defined as

$$\hbar w_0(k_\parallel) = E_g + E_0^{2D} + E_1^e + E_1^h + \frac{\hbar^2 k_\parallel^2}{2M}. \tag{30}$$

In the work [6], using a nonlocal susceptibility with microscopic parameters as that of Eq. (27), it was proved that reflectivity measurements as a function of the angle of incidence allow the measurement of the exciton-polariton modes in quantum wells. In particular, the comparison of the reflectivities for *s* and *p* polarization is useful for detecting the different polariton modes (T, L and Z) and to measure their dispersion relations. The optical properties of quantum wells of any shape, applying a mean-field polariton theory were investigated in Ref. [7]. There, it was shown that the splitting between the Z mode and the T mode depends mostly on the well thickness in fair agreement with the experiments of Ref. [31]. As is shown in [32], the spectra $\Delta_{45} = R_p - R_s^2$ of 45-degrees reflectometry may provide qualitatively new information on the exciton-polariton modes in a quantum well. The spectra $\Delta_{45}$ turn out to be very sensitive to the zeros of the dielectric function along the quantum-well growth direction and, therefore, allow to identify clearly the resonances associated with the Z mode, which are difficult to observe in simple spectra of reflectivity $R_p$. In particular, the 45-degrees reflectometry could be a powerful tool for studying Z exciton-polariton modes in near-surface quantum wells.

The optical properties of quasi-2D exciton quantum wells are generally studied assuming the QW interfaces to be ideally flat. However, realistic structures of quantum wells have inherent interface roughness, which gives rise to fluctuations in the confinement potentials $V_e(\vec{r}_e)$ and $V_h(\vec{r}_h)$ and, consequently, to a considerable increase of the relaxation frequency ($n$) as well as a shift of the resonance frequency ($w_0$). The variation of $n$ and $w_0$, due to interface roughness, depends on the frequency $w$ and, therefore, can alter the line-shape of resonances in optical spectra such as absorption [33], reflectivity [34], and diffuse reflection [35].

### *4.2. Magnetoexcitons in near-surface quantum wells*

The study of excitons in quantum wells under the action of a strong external transverse magnetic field is of great interest at present [36-40]. The discrete energy spectrum of such excitons is due to the size quantization in the direction perpendicular to the well plane and the action of the transverse magnetic field (Landau quantization). Both the confinement potential and the transverse magnetic field inhibit the Coulomb attraction between the electron and hole.

Optical spectra of exciton quantum wells subjected to an external magnetic field, $\vec{H}_0 \parallel \hat{z}$, can be calculated by using the system of equations (1), (4), and (5) with Hamiltoninan $H_{vc}$ given by

$$H_{vc} = E_g - \frac{\hbar^2}{2m_e}\frac{\partial^2}{\partial z_e^2} - \frac{\hbar^2}{2m_h}\frac{\partial^2}{\partial z_h^2} + V_e(\vec{r}_e) + V_h(\vec{r}_h) +$$

$$+ \frac{1}{2m_e}\left(-i\hbar\nabla_{re} + \frac{e}{c}\vec{A}_e\right)^2 + \frac{1}{2m_h}\left(-i\hbar\nabla_{rh} - \frac{e}{c}\vec{A}_h\right)^2 +$$

$$+ V_c(r) + V_s(\vec{r}_e, \vec{r}_h). \tag{31}$$

Here, $\vec{A}_{e,h} = (1/2)\vec{H}_0 \times \vec{r}_{e,h}$ is the vector potential, other quantities are defined above.

As is shown in Refs. [36-38,40] the magnetoexciton resonances in optical spectra (photoluminescence [36-38], and reflectivity [40]) for near-surface quantum wells undergo a blue shift as the cap layer thickness decreases. The blue shift of magnetoexciton levels is attributed to two factors. The first one is connected to the influence of the high vacuum potential barrier close to the quantum well, i.e. to the condition of vanishing of the coherent amplitude $\vec{Y}(\vec{r}_e, \vec{r}_h)$ at the sample surface. The second factor is owing to the interaction of the electron and hole with their image charges, i.e. to the image potential $V_{im}(\vec{r}_e, \vec{r}_h)$ (3).

## 5. Conclusion

We have presented the microscopic formalism of the coherent amplitude, which is useful for calculating optical functions of quantum wells in the regimes of both strong and weak exciton confinement. In numerous experimental and theoretical works it has been established that the optical properties of near-surface quantum wells (thin films on substrate, extrinsic transition layers, quasi-2D quantum wells with narrow cap layers) are strongly affected by the interaction between the exciton and the sample surface. The different mechanisms of such an interaction were discussed.

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