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LETTER TO THE EDITOR Open Access

# Optical absorption at one-particle states of charge carriers in semiconductor quantum dots

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**Abstract**

It is shown that the oscillator strengths and dipole moments of the transitions for one-particle states of charge carriers that originated in semiconductor quantum dots take on large values, significantly exceeding the typical values of the corresponding values for semiconductor materials. In the framework of the dipole approach, it is shown that large values of the cross section of light in the studied quasi-zero-dimensional systems enable the use of such nanosystems as new strongly absorbing nanomaterials.

**Keywords:** Optical absorption; Charge carriers; Semiconductor quantum dots

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## Correspondence/Findings

### Introduction

At present, the optical [1-5] and electro-optical [6-8] properties of quasi-zero-dimensional structures are extensively studied. Such structures commonly consist of spherical semiconductor nanocrystals (the so-called quantum dots (QDs)) with a radius $a \approx 1$ to $10^2$ nm grown in semiconductor (or dielectric) matrices. The studies in this field are motivated by the fact that such heterophase systems represent as new promising materials for the development of new components of nonlinear optoelectronics to be used, specifically, for controlling optical signals in optical computers [9-12] or for manufacturing active layers of semiconductor injection lasers [5,13-17].

The optical and electro-optical properties of such quasi-zero-dimensional structures depend on the energy spectrum of a spatially confined electron–hole pair (EHP), i.e., an exciton [1-8]. By the methods of optical spectroscopy, the effects of quantum confinement on the energy spectra of electrons [1,2] and excitons [3,4] were revealed in these heterophase structures.

Previously [18], the conditions for the localization of charge carriers near the spherical interface between two dielectric media were analyzed. In this case, the polarization interaction of a charge carrier with the surface charge induced at the spherical interface, $U(r,a)$, depends on the relative permittivity $\varepsilon = \varepsilon_1/\varepsilon_2$. Here, $r$ is the spacing between the charge carrier and the center of the dielectric particle, $a$ is the radius of the particle, and $\varepsilon_1$ and $\varepsilon_2$ are the permittivities of the surrounding medium and of the dielectric particle embedded in the medium, respectively.

For the charge carriers in motion near the dielectric particle, there are two possibilities:

1. Due to the polarization interaction $U(r,a)$, the carriers can be attracted to the particle surface (to the outer or the inner surface at $\varepsilon < 1$ or $\varepsilon > 1$, respectively), with the formation of outer [19,20] or inner [21] surface states.
2. Due to the polarization interaction $U(r,a)$, the carriers can be, at $\varepsilon < 1$, repelled from the inner surface of the dielectric particle, with the formation of bulk local states inside the particle bulk [22-24]; in this case, the spectrum of the low-energy bulk states is of an oscillatory shape.

It has been shown [18-24] that the formation of the above-mentioned local states is of a threshold-type nature and is possible if the radius of the dielectric particle $a$ is larger than a certain critical radius $\alpha_c$:

$$a \geq a_c = 6|\beta|^{-1} a_{Bi}, \quad (1)$$

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where $\beta = (\varepsilon_1 - \varepsilon_2)/(\varepsilon_1 + \varepsilon_2)$ and $\alpha_{Bi}$ is the Bohr radius of a charge carrier in a medium with the permittivity $\varepsilon_1$ ($i$ = 1, 2).

In [14-17], the optical properties of an array of InAs and InSb QDs in GaAs and GaSb matrices and the corresponding operational characteristics of injection lasers, with the active region on the basis of this array, were studied experimentally. In these studies, a large short-wavelength shift of the laser emission line was observed for the array of QDs. In such an array, the energy spectrum of charge carriers is completely discrete [21-24] if the QDs are smaller than $a \approx 1$ to 7 nm in size. In the first-order approximation, the spectrum of such quantum-confined states can be described as a spectrum of a charge carrier in a spherically symmetric well with infinitely high walls.

Until now, there have been no theoretical investigations on optical absorption and scattering at such discrete states in arrays of QDs. To close the gap in this area, here we develop a theory of the interaction of an electromagnetic with one-particle quantum-confined states of charge carriers that originate in the bulk of a semiconductor QD.

In conclusion, we briefly discuss possible physical situations in which the results obtained above can be used for interpreting the experimental data.

### Hamiltonian for an electron–hole pair in a quantum dot

We consider a simple model in which a quasi-zero-dimensional system is defined as a neutral spherical semiconductor QD of radius $a$ and permittivity $\varepsilon_2$, embedded in a surrounding with permittivity $\varepsilon_1$. Let an electron ($e$) and hole ($h$), whose effective masses are, correspondingly, $m_e$ and $m_h$, be the motion in this QD. Let the spacing between the electron or hole and the QD center be $r_e$ or $r_h$, respectively. We assume that the bands for electrons and hole are parabolic. Along with the QD radius $a$, the characteristic lengths of the problem are $a_e$, $a_h$, and $a_{ex}$, where

$$a_e = \varepsilon_2\hbar^2/m_e e^2, a_h = \varepsilon_2\hbar^2/m_h e^2, a_{ex} = \varepsilon_2\hbar^2/\mu e^2 \quad (2)$$

are the Bohr radii of the electron, hole, and exciton in the semiconductor with the permittivity $\varepsilon_2$, respectively, and $\mu = m_e m_h/(m_e + m_h)$ is the exciton effective mass. All of the characteristic lengths of the problem are much larger than the interatomic spacing $\alpha_0$:

$$a, a_e, a_h, a_{ex} \gg a_0. \quad (3)$$

This allows us to treat the motion of the electron and hole in the QD in the effective-mass approximation.

In the context of the above-described model and approximations for a quasi-zero-dimensional system, the Hamiltonian of the EHP is

$$H = -\frac{\hbar^2}{2m_e}\Delta_e - \frac{\hbar^2}{2m_h}\Delta_h + V_{eh}(r_e, r_h) + U(r_e, r_h, a) + E_g, \quad (4)$$

where the first two terms in the sum define the kinetic energy of the electron and hole, $E_g$ is the energy bandgap in the bulk (unbounded) semiconductor with the permittivity $\varepsilon_2$, $V_{eh}(\boldsymbol{r_e},\boldsymbol{r_h})$ is the energy of the electron Coulomb interaction:

$$V_{eh}(r_e, r_h) = \frac{e^2}{2\varepsilon_2 a}\frac{2a}{\left(r_e^2 - 2r_e, r_h\cos\theta + r_h^2\right)^{1/2}} \quad (5)$$

with the angle $\theta$ between the vector $\boldsymbol{r_e}$ and $\boldsymbol{r_h}$, and $U(\boldsymbol{r_e},\boldsymbol{r_h},a)$ is the energy of the interaction of the electron and hole with the polarization field induced by the electron and hole at the spherical interface between the two media. For arbitrary values of $\varepsilon_1$ and $\varepsilon_2$, the interaction energy $U(\boldsymbol{r_e},\boldsymbol{r_h},a)$ can be represented analytically as [18,24]

$$U(r_e, r_h, a) = -\frac{e^2\beta}{\varepsilon_2 a}\frac{1}{\left[(r_e r_h/a^2)^2 - 2(r_e r_h/a^2)\cos\theta + 1\right]^{1/2}} - \frac{e^2\beta}{2(\varepsilon_1+\varepsilon_2)a}\int_0^\infty \frac{dy(a^2/r_h y)^\alpha \Theta(y - a^2/r_h)}{|r_e - y(r_e/r_h)|} - \frac{e^2\beta}{2(\varepsilon_1+\varepsilon_2)a}\int_0^\infty \frac{dy(a^2/r_e y)^\alpha \Theta(y - a^2/r_e)}{|r_h - y(r_h/r_e)|}, \quad (6)$$

where $\Theta(x)$ is the unit step function and

$$\alpha = \frac{\varepsilon_1}{\varepsilon_1 + \varepsilon_2}. \quad (7)$$

### The spectrum of charge carriers in a quantum dot

In the bulk of a QD, the electron (hole) energy levels can originate. Their energies are defined as [25]

$$E_{n,l}(a) = \frac{\hbar^2}{2m_{e(h)}a^2}\varphi_{n,l}^2, \quad (8)$$

where the subscripts ($n$,$l$) refer to the corresponding quantum size-confined states. Here, $n$ and $l$ are the principal and azimuthal quantum numbers for the electron (hole), respectively, and $\varphi_{n,l}$ are the roots of the Bessel function, i.e., $J_{1+\frac{1}{2}}(\varphi_{n,l}) = 0$. For the quantum-confined levels to originate, it is necessary that in the Hamiltonian (4), the electron (hole) energy $E_{n,l}(a)$ (8) be much larger

than the energy of the interaction of the electron (hole) with the polarization field $U(a)$ (6) generated at the spherical QD-dielectric (semiconductor) matrix interface:

$$E_{n,l}(a) = \frac{\hbar^2}{2m_{e(h)}a^2}\varphi_{n,l}^2 \gg U(a) \approx \frac{e^2\beta}{2\varepsilon_2 a}. \quad (9)$$

Condition (9) is satisfied for QDs of radii

$$a \ll a_s^{e(h)} = \frac{\varphi_{n,l}^2}{\beta}a_{e(h)}. \quad (10)$$

At room temperature $T_0$, the discrete levels of the electron (hole) $E_{n,l}$ (8) in the QD are slightly broadened if the energy separation between the levels is [20-22]

$$\Delta E_{n,l}(a) = E_{n',l} - E_{n,l}(a) \ll kT_0. \quad (11)$$

Taking into account (8), we can rewrite inequality (11) as

$$\frac{\hbar^2}{2m_{e(h)}a^2}\frac{\varphi_{n',l}^2 - \varphi_{n,l}^2}{kT_0} = \eta(a) \ll 1. \quad (12)$$

Formula (8), describing the spectrum of charge carriers in a QD, is applicable to the lowest states $(n,l)$ that satisfy the inequality

$$\Delta E_{n,l}(a) \ll \Delta V_0(a), \quad (13)$$

where $\Delta V_0(a)$ is the depth of the potential well for electrons in the QD. For example, for the CdS QDs whose sizes satisfy inequality (10), the value $\Delta V_0$ is 2.3 to 2.5 eV [26].

If condition (10) is satisfied, we can use, for the electron (hole) wave function in a QD, the wave function of an electron (hole) in a spherical quantum well with infinitely high walls [27]:

$$\Psi_{n,l,m}(r,\theta,\varphi) = Y_{l,m}(\theta,\varphi)\frac{J_{1+1/2}(\varphi_{n,l})\sqrt{2}}{J_{1+3/2}(\varphi_{n,l})a\sqrt{r}}, \quad (14)$$

where $r = r_e$ or $r = r_h$ is the distance of the electron or hole from the QD center, respectively, $\theta$ and $\varphi$ are the azimuthal and polar angles that define the orientation of the radius vector of the electron (hole), respectively, $Y_{l,m}$ are the normalized spherical functions ($m$ is the magnetic quantum number of the electron or hole), and $J_\nu(x)$ are the Bessel functions that can be expressed in terms of the spherical Bessel functions $j_\nu(x)$ [27] as

$$J_{1+1/2}(\varphi_{n,l}) = \sqrt{2/\pi}\varphi_{l,n}j_l(\varphi_{l,n})$$

$$J_{1+3/2}(\varphi_{n,l}) = \sqrt{2/\pi}\varphi_{l,n}j_{l+1}(\varphi_{l,n}). \quad (15)$$

### Dipole moments of charge carrier transitions in a quantum dot

In the frequency region corresponding to the above-considered states of charge carriers in the QD bulk, $(n,l)$, the wavelength of light is much larger than the dimensions of these states ($\sim a_e$, $a_h$), that is the behavior of approximation. In this case, the operator of the dipole moment of the electron (hole) located in the QD bulk is expressed as [28]

$$D(r) = \frac{3\varepsilon_1}{2\varepsilon_1 + \varepsilon_2}D^0(r), D^0(r) = er. \quad (16)$$

To estimate the value of the dipole moment, it is sufficient to consider the transition between the lowest discrete states (8), e.g., between the ground states $|1s\rangle = (n = 1, l = 0, m = 0)$ and $|1p\rangle = (n = 1, l = 1, m = 0)$. To calculate the matrix element of the dipole moment of the charge carrier transition from the $1s$ state to the $1p$ state, $D_{1,0}(a)$, we assume that the uniform field of the light wave $E(\omega,t)$ is directed only along the axis $Z$ ($\omega$ is the wave frequency). In this case, we take the dipole moment $D(r)$ (16) induced by the field $E(\omega,t)$ as the perturbation responsible for such dipole transition. The expression for the dipole moment of the transition $D_{1,0}(a)$ follows from formulas (16) and the expression for the dipole moment of the transition in free space:

$$D_{1,0}^0(a) = \langle 1s|D^0(r)|1p\rangle = e\langle 1s|r|1p\rangle. \quad (17)$$

Taking into account (14) and (15), we can write the wave functions of the $|1s\rangle$ and $|1p\rangle$ states as

$$|1s\rangle = \Psi_{1,0,0}(r,\theta) = Y_{0,0}(\theta)\frac{2}{a^{3/2}}\frac{j_0(\varphi_{0,1}(r/a))}{j_1(\varphi_{0,1})}, \quad (18)$$

$$|1p\rangle = \psi_{1,0,0}(r,\theta) = Y_{1,0}(\theta)\frac{2}{\alpha^{3/2}}\frac{j_1(\varphi_{1,1}(r/\alpha))}{j_2(\varphi_{1,1})}. \quad (19)$$

On substitution of (18) and (19) into formula (17) and integration, we obtain the expression for the dipole moment of the transition in free space as follows:

$$\begin{aligned} D_{1,0}^0(a) &= \frac{2\pi\sqrt{6}}{3\varphi_{1,1}j_2(\varphi_{1,1})\left(\varphi_{1,1}^2 - \pi^2\right)} \\ &\times \left[\cos\varphi_{1,1} - \frac{\left(3\varphi_{1,1}^2 - \pi^2\right)}{\varphi_{1,1}\left(\varphi_{1,1}^2 - \pi^2\right)}\sin\varphi_{1,1}\right]ea \\ &= 0.433ea. \end{aligned} \quad (20)$$

Then, according to (20) and (16), the dipole moment of the transition in the QD with the permittivity $\varepsilon_2$ in the surrounding matrix with the permittivity $\varepsilon_1$ is

$$D_{1,0}(a) = \Lambda D^0_{1,0}(a) = \Lambda 0.433ea, \tag{21}$$

where

$$\Lambda = \frac{3\varepsilon_1}{2\varepsilon_1 + \varepsilon_2}. \tag{22}$$

### Optical absorption on light at one-particle states of charge carriers in quantum dots

Using the above results for the matrix element of the dipole moment of the transition $D_{1,0}$ (formulas (21) and (22)), we can elucidate the behavior of the semiconductor quasi-zero-dimensional systems on absorbing the energy of the electromagnetic field in the frequency region corresponding to the energies of the quantum-confined states in the QD $E_{n,l}$ (8). The absorption cross section of a spherical QD of radius $a$ can be expressed in terms of the polarizability of QD, $A^{\cdot\cdot}(\omega, a)$ $A$, as [28]

$$\sigma_{\mathrm{abs}}(\omega, a) = 4\pi(\omega/c)A^{\cdot\cdot}(\omega, a), \tag{23}$$

where $\omega$ is the frequency of the external electromagnetic field.

The polarizability $A^{\cdot\cdot}$ can be easily determined if the QD is considered as a single giant ion. Let the radius of QD be $a$ (10). In such QD, the quantum-confined states of charge carriers $(n,l)$ (8) are formed. At room temperature, these states are slightly broadened, satisfying inequality (12). In this case, the polarizability of the charged QD, $A^{\cdot\cdot}$, can be expressed in terms of the matrix element of the dipole moment of the transition $D_{1,0}(a)$ (21) between the lowest $1s$ and $1p$ states [29,30]:

$$A^{\cdot\cdot}(\omega, a) = \frac{e^2}{m_{\mathrm{e(h)}}} \frac{f_{0,1}}{\omega_1^2(a) - \omega^2 - i\omega\Gamma_1(a)}, \tag{24}$$

where

$$f_{0,1} = \frac{2m_{\mathrm{e(h)}}}{\hbar e^2}[\omega_1(a) - \omega_0(a)]\left|D_{1,0}(a)\right|^2 \tag{25}$$

is the oscillator strength of the transition of a charge carrier with the effective mass $m_{\mathrm{e}}$ (or $m_{\mathrm{h}}$) from the ground $1s$ state to the $1p$ state, $\omega_1 = E_{1.1}$ and $\omega_0 = E_{1.0}$ are, correspondingly, the energies of the discrete $1s$ and $1p$ levels by formula (8), and $\Gamma_1(a)$ is the width of the $1p$ level [20,22]. Taking into account formulas (8) and (21), we can express the oscillator strength (25) of the transition as

$$f_{0,1} = \left(\varphi^2_{1,1} - \pi^2\right)\frac{\Lambda^2\left(D^0_{1,0}(a)\right)^2}{e^2a^2} = 0.967 \times 10^{-1}\Lambda^2 \tag{26}$$

We assume that the frequency $\omega$ of the wave of light is far from the resonance frequency $\omega_1$ of the discrete $1p$ state and, in addition, that the broadening $\Gamma_1$ of the $1p$ level is small, i.e., $\Gamma_1/\omega_1 \ll 1$ [20-22]. Then, for the qualitative estimate of the QD polarizability $A^{\cdot\cdot}(\omega, a)$ (24), we obtain, with regard to (8), the following expression:

$$A^{\cdot\cdot}(a) = \frac{4f_{0,1}}{\varphi^4_{1,1}} \frac{m_{\mathrm{e(h)}}}{m_0}\left(\frac{a}{a_{\mathrm{B}}}\right)^4 a^3_{\mathrm{B}}, \tag{27}$$

where $a_{\mathrm{B}} = \hbar^2/m_0e^2$ is the Bohr radius of an electron in free space. Now we write the expression for the cross section of elastic scattering of the electromagnetic wave with frequency $\omega$ by the QD of radius $a$ [28] as

$$\sigma_{\mathrm{sc}}(\omega, a) = 2^7 3^{-3}\pi^3(\omega/c)^4\left|A^{\cdot\cdot}(a)\right|^2. \tag{28}$$

The theory developed here for QDs is applicable only to the intraband transitions of electrons (holes) whose spectrum $E^{e(h)}_{nl}(a)$ is defined by formula (8). The optical attenuation coefficient, which involves both the absorption and scattering of light by one-particle quantum-confined states of charge carriers in the QD bulk of radius $a$, can be written, for the concentration of QDs $N$, as [29,30]

$$\gamma(\omega, a) = N[\sigma_{\mathrm{abs}}(\omega, a) + \sigma_{\mathrm{sc}}(\omega, a)] \tag{29}$$

Formula (29) is applicable to an array of QDs that do not interact with each other. The condition, such that the QDs of radius $a$ and concentration $N$ do not interact with each other, is reduced to the requirement that the spacing between the QDs ($\sim N^{-1/3}$) is much larger than the dimensions of the above-considered one-particle states ($\sim a_{\mathrm{e(h)}}$), i.e.,

$$a_{\mathrm{e(h)}}N^{-1/3} \ll 1. \tag{30}$$

With $a_{\mathrm{e(h)}} \approx 5.0$ nm criterion (30) is satisfied for the concentrations of QDs $N \le 10^{15}$ cm$^{-3}$ achievable under the experimental conditions of [13-17] for a number of III-V semiconductors.

### Comparison of the theory with experiments

Similar to [26], we can assume that under the experimental conditions of [13-17], the annealing of the arrays of InAs and InSb QDs in the GaAs and GaSb matrices at the temperature $T$ = 273 K induces the thermal emission of a light electron, so that a hole alone remains in the QD bulk. In this case, the electron may be localized at a deep trap in the matrix. If the distance $d$, from this

**Table 1 Parameters of quantum-confined states of hole in motion in quantum dots**

| Quantum dots | | | | Matrices | | $a_{Bh}$ | Λ | $a$ | $f_{0,1}$ | $D^0_{1,0}$ | $D_{1,0} = \Lambda D^0_{1,0}$ | $\eta$ | $A^{``}$ | $\sigma_{abs}$ |
|---|---|---|---|---|---|---|---|---|---|---|---|---|---|---|
| Semiconductor | $\varepsilon_2$ | $\frac{m_e}{m_0}$ | $\frac{m_h}{m_0}$ | Semiconductor | $\varepsilon_1$ | (nm) | | (nm) | ($10^{-1}$) | (D) | (D) | (%) | ($10^{-22}$ cm$^3$) | ($10^{-18}$ cm$^2$) |
| InAs | 12.5 | 0.028 | 0.33 | GaAs | 12 | 2.0 | 0.99 | 2.0 | 0.95 | 6.12 | 6.04 | 8.7 | 9.2 | 7.5 |
| | | | | | | | | 3.0 | | 9.2 | 9.06 | 19.6 | 47 | 38 |
| | | | | | | | | 4.0 | | 12.2 | 12.1 | 35 | 140 | 120 |
| InSb | 18 | 0.013 | 0.18 | GaSb | 12 | 5.3 | 0.94 | 2.0 | 0.85 | 6.12 | 5.74 | 4.7 | 4.6 | 3.7 |
| | | | | | | | | 3.0 | | 9.2 | 8.61 | 10.6 | 23 | 19 |
| | | | | | | | | 4.0 | | 12.2 | 11.5 | 19 | 74 | 61 |
| | | | | | | | | | | 15.3 | 14.4 | 29.4 | 180 | 160 |

The radius $a$ and permittivity $\varepsilon_2$ are embedded in semiconductor matrices with permittivity $\varepsilon_1$. Here, $m_e$ and $m_h$ are the effective masses of an electron and a hole in the QD, respectively; $a_{Bh}$ is the Bohr radius of a hole in the QD; $\Lambda = 3\varepsilon_1/(2\varepsilon_1 + \varepsilon_2)$; $f_{0,1}$ is the oscillator strength of the transition; $D^0_{1,0}$ and $D_{1,0}$ are the dipole moments of the transitions in free space and in the QD, respectively; $\eta$ is the parameter of the broadening of states at room temperature; $A^{``}$ is the polarizability of the QD; and $\sigma_{abs}$ is the absorption cross section. The unit of measure of the dipole moment is debye (D), (i.e., the dipole moment $D = er$ with $r$ = 1 Å (debye)).

trap to the QD center, is large compared to the QD radius $a$ ($d \gg a$), the Coulomb electron–hole interaction $V_{eh}(r_e,r_h)$ (5) in the Hamiltonian (4) can be disregarded. As a result, the one-particle quantum-confined hole states ($n$,$l$), with the energy spectrum $E_{n,l}(a)$ described by formula (8), appear in the QD bulk. Now we roughly estimate the cross sections of optical absorption $\sigma_{abs}$ ((25) and (29)) and $\sigma_{sc}$ (28) at the quantum-confined hole state in the QDs for the selected $1s \rightarrow 1p$ transition under the experimental conditions of [13,17]. For the rough estimation of the cross sections of optical absorption and scattering, we use expressions (23), (27), and (28) on the assumptions that the frequency of the light wave $\omega$ is far from the resonance frequency $\omega_1$ of the discrete hole state in the QD and that the broadening $\Gamma_1$ of the energy level $E_{1.1} = \omega_1$ (8) is small [20-22] ($\Gamma_1/\omega_1 \ll 1$). In this case, the absorption cross section $\sigma_{abs}$ and the scattering cross section $\sigma_{sc}$ take the following forms:

$$\sigma_{abs}(\omega,\alpha) = \frac{16\pi f_{0,1}}{\varphi^4_{1,1}} \frac{\omega}{c} \frac{m_h}{m_0} \left(\frac{\alpha}{\alpha_B}\right)^4 \alpha_B^3, \qquad (31)$$

$$\sigma_{sc}(\omega,a) = \frac{1^{11}\pi^3 {f_{0,1}}^2}{3^3\varphi^8_{1,1}} \left(\frac{\omega}{c}\right)^4 \left(\frac{m_h}{m_0}\right)^2 \left(\frac{a}{a_B}\right)^8 a_B^6. \qquad (32)$$

The estimated parameters of the hole states in QDs dispersed in the III-V semiconductors are listed in Table 1. It is worth noting that at room temperature, the quantum-confined hole states are slightly broadened; the parameter $\eta(a)$ (14) does not exceed 35%.

The quasi-zero-dimensional systems considered here are essential nonlinear media with respect to infrared radiation. In fact, the dipole moments of the transitions in QDs of radii $a \approx 20$ to 40 Å are rather large in magnitude: $D_{1,0} \approx 10$ D (see table), being many times larger than the values $D \approx 0.1$ D typical for bulk III-V semiconductors [31,32]. In addition, according to the selection rules for QDs in the electromagnetic field, the dipole transitions between the nearest levels $E_{n,l}$ (8) are allowed [27]. Under these transitions, the azimuthal quantum number changes by unity.

From the estimates presented in the table, it follows that for QDs of radii $a \approx 2.0$ to 4.0 nm, the absorption cross section can be as large as $\sigma_{abs} \approx 10^{16}$ cm$^2$. This value is 8 orders of magnitude larger than the typical absorption cross sections for atoms [33]. Since the scattering cross sections $\sigma_{sc}$ (32) under the experimental conditions of [13-17] are negligible compared to the corresponding absorption cross sections $\sigma_{abs}$ (31) ($\sigma_{sc}/\sigma_{abs} \approx 10^{-12}$), the estimates for $\sigma_{sc}$ are not included in the table.

Thus, the optical attenuation coefficient $\gamma$ (31) for QDs is controlled mainly by the absorption process at the quantum-confined states of charge carriers ($n$,$l$) (8). If the concentration of QDs of radii $a$ satisfying condition (38) is $N \approx 10^{15}$ cm$^{-3}$, the optical attenuation coefficient, on absorption at the QDs with the parameters listed in the table, takes the value $\gamma \approx 0.1$ cm$^{-1}$.

The large optical absorption cross section and attenuation coefficients in the quasi-zero-dimensional systems treated above allow the use of such heterophase structures as new efficient absorbers of electromagnetic waves in a wide wavelength range variable over wide limits in accordance with the nature of materials in contact [34].

#### Competing interests
The author declares that he has no competing interests.

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