Optical anisotropy and the direction of polarization of exciton emissions in a semiconductor quantum dot: Effect of heavy- and light-hole mixing

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The dependence of the directions of polarization of exciton emissions, fine structure splittings (FSS), and polarization anisotropy on the light- and heavy-hole (LH–HH) mixing in semiconductor quantum dots (QDs) is investigated using a mesoscopic model. In general, all QDs have a four-fold exciton ground state. Two exciton states have directions of polarization in the growth-plane, while the other two are along the growth direction of the QD. The LH–HH mixing does affect the FSS and polarization anisotropy of bright exciton states in the growth-plane in the low symmetry QDs (e.g., C_{2V}, C_S, C_1), while it has no effect on the FSS and polarization anisotropy in high symmetry QDs (e.g., C_{3V}, D_{2d}). When the hole ground state is pure HH or LH, the bright exciton states in the growth-plane are normal to each other. The LH–HH mixing affects the relative intensities and directions of bright exciton states in the growth-plane of the QD. The polarization anisotropy of exciton emissions in the growth-plane of the QD is independent of the phase angle of LH–HH mixing but strongly depends on the magnitude of LH–HH mixing in low symmetry QDs.

Keywords: light–heavy hole mixing, excitons, quantum dots

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1. Introduction

In the standard sp-bonded semiconductors the conduction band minimum (CBM) is formed of s atomic orbitals, while the valence band maximum (VBM) is formed of p atomic orbitals. Due to the spin–orbit interactions the valence band splits into the heavy-hole (HH), light-hole (LH), and split-off (SO) bands. The HH and LH states are degenerate at the Γ point of the Brillouin zone of the bulk semiconductor. However, in a semiconductor quantum dot (QD) depending upon the confinement, growth direction, and materials of the QD and barrier matrix, the single particle hole states can have a dominantly HH, mixed HH, and LH or dominantly LH character. For instance, in the standard flat QDs grown on the (001) surface the hole states have a dominant HH character due to the strong confinement along the z direction ([001] growth direction of the QD).[1] Excitation of an electron–hole pair forms an exciton in a QD due to the strong Coulomb interactions. However, its lifetime is of the order of nanoseconds (ns) and recombination of electron–hole pair emits a photon.[2,3] Excitation of two electron–hole pairs forms a biexciton which emits two photons via the decay path of an intermediate exciton state. The charging of QDs with an extra electron or hole in addition to an electron–hole pair gives rise to the charged excitons. The ground states of biexciton and charged excitons have no fine structure splitting (FSS) because of the zero total spin of electrons or holes or both. However, the ground state of a neutral exciton has the non-zero FSS due to the non-zero spin of electron and hole states, and depending upon the symmetry of the confining potential.[4,5]

Despite extensive experimental and theoretical studies the FSS of an exciton state is not completely understood. It is the common feature of all the QDs grown along different growth directions in experiments. It varies from about 10 μeV–100 μeV in InGaAs/GaAs QDs[6–8] and GaAs/AlGaAs QDs[9,10] to 1000's μeV in InAs/AlGaAs QDs.[11] In theoretical calculations the FSS is usually found to be quite low.[4,12,13] The origin of large values of FSS in some QDs is not yet understood. The excitonic emission in experimental QDs in a given sample have fluctuations in the FSS and polarization directions with a majority of emissions along a certain direction. This fluctuation is poorly understood. The origin of large polarization anisotropy of exciton emissions measured in the growth-plane of QDs is also not understood.[14] However, it is of immense interest to understand the FSS, the polarization anisotropy, and the directions of polarization of exciton emissions of QDs for their practical uses in the various proposed potential applications.[15,16]

The FSS of an exciton state is due to the electron–hole exchange interactions. In the bulk zincblende semiconductors the ground state exciton is 8-fold degenerate in the absence of electron–hole exchange interactions because it is formed by the hole in 4-fold (including spin) degenerate valence band Γ_{8V} and 2-fold degenerate conduction band Γ_{6C}. The electron–hole exchange interactions splits the 8-fold exciton state into the 5-fold bright exciton states and 3-fold dark exciton states. In semiconductor QDs the symmetry is further reduced depend-

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ing upon the growth surface. The effect of LH–HH mixing in hole states on the optical properties of semiconductor QDs has been investigated by using different approaches.[17–19] It is still an open question whether the symmetry of a QD has any influence on the LH–HH mixing and its optical properties. In this paper, we investigate the effect of LH–HH mixing on the polarization anisotropy and the directions of polarization of exciton emissions of semiconductor QDs in detail, using a simple mesoscopic model.

We find that the LH–HH mixing does affect the FSS and polarization anisotropy of bright states in the growth-plane in the low symmetry QDs (e.g., C_{2V}, C_S, C_1), while it has no effect on the FSS in high symmetry QDs (e.g., C_{3V}, D_{2d}). When the hole ground state is pure HH or LH, the bright states in the growth-plane are normal to each other. The LH–HH mixing affects the relative intensities and directions of bright states in the growth-plane of the QD. The polarization anisotropy of exciton emissions in the growth-plane of the QD is independent of the phase angle of LH–HH mixing but strongly depends on the magnitude of LH–HH mixing in low symmetry QDs. The rest of the paper is organized as follows: in Section 2 we briefly discuss the mesoscopic model, in Section 3 we present the results and discussion, and in the last Section 4 we conclude the results.

2. Mesoscopic model

The exchange interactions between the electron and hole states are given as

$$E_{\rm ex} = \iint {\rm d}r_{\rm e}\,{\rm d}r_{\rm h}\Phi_X^{\dagger}(r_{\rm e},r_{\rm h})\left(\frac{1}{|r_{\rm e}-r_{\rm h}|}\right)\Phi_X(r_{\rm h},r_{\rm e}),$$

where Φ_X is the exciton wavefunction. $r_{\rm e}$ and $r_{\rm h}$ are the position coordinates of electron and hole, respectively. This equation requires the exciton wavefunctions for its solution. The computational approaches such as $k.p$ effective mass approximation,[20,21] tight binding approximation,[22,23] and empirical pseudopotentials[24,25] construct these wavefunctions Φ_X by calculating the single particle wavefunctions of electrons and holes. However, there is also a simple mesoscopic model for the exchange interactions based upon the theory of invariants. In this model the exchange interaction Hamiltonian between electron and hole is written as[26]

$$H_{\rm ex} = -\sum_{i=x,y,z}(a_i\sigma_i\otimes J_i + b_i\sigma_i\otimes J_i^3), \tag{1}$$

where σ_i and J_i are the spin $1/2$ and $3/2$ matrices, respectively. a_i and b_i are the spin–spin coupling constants. These constants are basically the adjustable parameters introduced in the theory of invariants. This Hamiltonian has been used to describe the exciton FSS in QDs and quantum wells,[26–29] where hole states are derived from the VBM with atomic p symmetry and the electron states originate from the CBM with atomic s symmetry.

3. Results and discussion

3.1. Energetic splitting of exciton states

In a semiconductor QD the hole states have pure HH or mixed HH-LH or pure LH character depending upon the strength of confinement, the growth direction and the material of the QD and barrier matrix. In general, the hole states can be written as

$$\Psi_{\rm h}^{\pm} = A\left|\frac{3}{2},\pm\frac{3}{2}\right\rangle + B\left|\frac{3}{2},\mp\frac{1}{2}\right\rangle, \tag{2}$$

where $|\frac{3}{2},\pm\frac{3}{2}\rangle$ and $\frac{3}{2},\mp\frac{1}{2}\rangle$ represents HH and LH states in the notation $|J,J_z\rangle$, J being the total angular momentum and J_z is its component along the z direction. A and B are the complex numbers with the condition $|A|^2+|B|^2=1$. This is a general characteristic of the superposition of two quantum states. For simplicity we can consider $|A|=\beta$, then B can be written as $\pm\sqrt{(1-\beta^2)}\exp(i\phi)$, ϕ being the relative coupling quantum phase between the HH and LH states. β is a real number such that $0\le\beta^2\le 1$. In such a case we can write hole states as

$$\Psi_{\rm h}^{\pm} = \beta\left|\frac{3}{2},\pm\frac{3}{2}\right\rangle \mp\sqrt{(1-\beta^2)}\exp(\mp i\phi)\left|\frac{3}{2},\mp\frac{1}{2}\right\rangle, \tag{3}$$

where β^2 is the weight of the HH component in the hole states and $1-\beta^2$ is the weight of its LH component. β^2 is determined by the strength of confinement, the growth direction and the material of the QD and barrier matrix. For instance, in the case of flat QDs on (001) surface, $\beta^2\approx 1$ due to the strong confinement along the z direction.[1] However, in the case of tall QDs on the (001) surface or QDs on (110) and (111) surfaces, the values of β^2 are less than 1.[12,30] The electron states are given as

$$\Psi_{\rm e}^{\pm} = \left|\frac{1}{2},\pm\frac{1}{2}\right\rangle. \tag{4}$$

The exciton basis set $\{\phi_X^i\}$ has four states as $\{\Psi_{\rm e}^{+}\Psi_{\rm h}^{+}, \Psi_{\rm e}^{+}\Psi_{\rm h}^{-}, \Psi_{\rm e}^{-}\Psi_{\rm h}^{+}, \Psi_{\rm e}^{-}\Psi_{\rm h}^{-}\}$. Using Eqs. (1)–(4) the 4×4 matrix of exchange interactions $\langle\phi_X^i|H_{\rm ex}|\phi_X^j\rangle$ of the ground state exciton, is given by

$$\begin{pmatrix} K & 0 & 0 & \delta \\ 0 & -K & \Delta & 0 \\ 0 & \Delta^{\dagger} & -K & 0 \\ \delta^{\dagger} & 0 & 0 & K \end{pmatrix},$$

where $\delta^{\dagger}(\Delta^{\dagger})$ is the complex conjugate of $\delta\,(\Delta)$ and

$$2K = a_z\left(2\beta^2-\frac{1}{2}\right)+b_z\left(\frac{7}{2}\beta^2-\frac{1}{8}\right),$$
$$2\delta = (1-\beta^2)\left[(a_x-a_y)+\frac{5}{2}(b_x-b_y)\right]-\frac{3}{4}\beta^2(b_x-b_y),$$
$$2\Delta = (1-\beta^2)\left[(a_x+a_y)+\frac{5}{2}(b_x+b_y)\right]-\frac{3}{4}\beta^2(b_x+b_y).$$

The diagonalization of the above matrix gives the four eigenvalues of the ground state exciton as $K\pm\sqrt{\delta\delta^{\dagger}}$, $-K\pm\sqrt{\Delta\Delta^{\dagger}}$ with the corresponding eigenstates as $(1/\sqrt{2})\left(\Psi_{\rm e}^{+}\Psi_{\rm h}^{+}\pm\Psi_{\rm e}^{-}\Psi_{\rm h}^{-}\right)$, $(1/\sqrt{2})\left(\Psi_{\rm e}^{+}\Psi_{\rm h}^{-}\pm\Psi_{\rm e}^{-}\Psi_{\rm h}^{+}\right)$, respectively.

Pure HH: $-\beta^2 = 1$, we have

$$K = \frac{3}{4}a_z + \frac{27}{16}b_z, \tag{5}$$

$$\delta = -\frac{3}{8}(b_x - b_y), \tag{6}$$

$$\Delta = -\frac{3}{8}(b_x + b_y). \tag{7}$$

Pure LH: $-\beta^2 = 0$, we have

$$K = -\frac{1}{4}a_z - \frac{1}{16}b_z, \tag{8}$$

$$\delta = \frac{1}{2}(a_x - a_y) + \frac{5}{4}(b_x - b_y), \tag{9}$$

$$\Delta = \frac{1}{2}(a_x + a_y) + \frac{5}{4}(b_x + b_y). \tag{10}$$

The values of K, δ, and Δ can be calculated using the computational approaches such as the $k.p$ effective mass approximation,[20,21] the tight binding approximation[22,23] or the empirical pseudopotentials.[24,25] For given values of K, δ, and Δ in the cases of pure HH and LH states, we can calculate the a_i and b_i spin–spin coupling constants.

The underlying symmetry of a QD is quite crucial for determining the excitonic structure. In the case of low symmetry QDs such as C_{2V}, C_S, C_1, the a_x, a_y and b_x, b_y belong to different representations,[26] which means $a_x \neq a_y$ and $b_x \neq b_y$. These low symmetry QDs have nonzero FSS, which is further dependent upon the LH–HH mixing. For high symmetry QDs such as C_{3V} or D_{2d}, the a_x, a_y and b_x, b_y belong to the same representation,[26] which means $a_x = a_y$ and $b_x = b_y$. In such a case $\delta = 0$, i.e., FSS of bright states polarized in the growth-plane is vanishing irrespective of LH–HH mixing. This has been confirmed for InGaAs/GaAs and GaAs/AlGaAs QDs by the numerical calculations.[12,31,32]

3.2. Oscillator strengths of excitonic transitions

Oscillator strengths of excitonic transitions is determined by the envelope function and the Bloch states of the electron and hole states in a QD. The transition probability is given by the dipole matrix elements as

$$P = |\langle \Phi_X | \hat{e} \cdot \hat{p} | 0 \rangle|^2, \tag{11}$$

where $|0\rangle$ is the vacuum state and Φ_X is the exciton state wavefunction. The $\hat{e}$ is the polarization vector of light and $\hat{p}$ is the momentum operator. In the confinement regime where the carriers are treated independently, the above equation can be rewritten as

$$P = |\langle \Psi_{\rm e} | \hat{e} \cdot \hat{p} | \Psi_{\rm h} \rangle|^2, \tag{12}$$

where $\Psi_{\rm e}$ ($\Psi_{\rm h}$) is the single particle wavefunction of the electron (hole). Since envelope functions are slowly varying in terms of r, the operator $\hat{p}$ acts only on the Bloch part of the total wavefunction. Therefore, the above equation can be simplified as

$$P = |\langle u_{\rm c} | \hat{e} \cdot \hat{p} | u_{\rm v} \rangle|^2 |\langle f_{\rm e} | f_{\rm h} \rangle|^2. \tag{13}$$

This shows that the contributions of the envelope function and the Bloch states can be separated. The darkness or brightness of a given excitonic state is determined by the Bloch states contribution unless the envelope function overlap is zero. The envelope function part only increases or decreases the intensity. We discuss here the oscillator strengths of the excitonic transitions based upon their description on the Bloch states. The Bloch states of the conduction band in a semiconductor QD with zincblende structure have orbital s symmetry[33] and are given as

$$|\Psi_{\rm e}^+\rangle = |s\rangle|\uparrow\rangle, \tag{14}$$

$$|\Psi_{\rm e}^-\rangle = |s\rangle|\downarrow\rangle, \tag{15}$$

where $|s\rangle$ is a function that is invariant under all symmetry transformations of the Bravais lattice. We define the spin states $|\uparrow\rangle$ and $|\downarrow\rangle$ along the z axis, the [001] crystal direction. The electron states have total angular momentum, $J_{\rm e} = 1/2$. The Bloch states in the valance band have orbital p symmetry. By coupling the orbital angular momentum $l = 1$ and the spin $S = 1/2$ according to the usual Clebsch–Gordan theory, the HH states are obtained as[33]

$$|\Psi_{\rm hh}^+\rangle = -\frac{1}{\sqrt{2}}|X + {\rm i}Y\rangle|\Uparrow\rangle, \tag{16}$$

$$|\Psi_{\rm hh}^-\rangle = \frac{1}{\sqrt{2}}|X - {\rm i}Y\rangle|\Downarrow\rangle, \tag{17}$$

and the LH states as

$$|\Psi_{\rm lh}^+\rangle = -\frac{1}{\sqrt{6}}\left(|X + {\rm i}Y\rangle|\Downarrow\rangle\right) + \frac{2}{\sqrt{6}}\left(|Z\rangle|\Uparrow\rangle\right), \tag{18}$$

$$|\Psi_{\rm lh}^-\rangle = \frac{1}{\sqrt{6}}\left(|X - {\rm i}Y\rangle|\Uparrow\rangle\right) + \frac{2}{\sqrt{6}}\left(|Z\rangle|\Downarrow\rangle\right). \tag{19}$$

In the above expressions, $|X + {\rm i}Y\rangle = |X\rangle + {\rm i}|Y\rangle$ and the states $|X\rangle, |Y\rangle, |Z\rangle$ are the p-type orbital parts of the Bloch states, which transform like the coordinates x, y, z. The up (down) spins of p-type orbital are denoted as $\Uparrow$ ($\Downarrow$). The E-field vector of light is given as $e = (e_x, e_y, e_z)$, where $e_x = (1/\sqrt{2})(\hat{x} + \hat{y})$, $e_y = (1/\sqrt{2})(-\hat{x} + \hat{y})$, and $e_z = \hat{z}$ for the QDs on the (001) surface, $e_x = \hat{z}$, $e_y = (1/\sqrt{2})(-\hat{x} + \hat{y})$ and $e_z = (1/\sqrt{2})(\hat{x} + \hat{y})$ for the QDs on the (110) surface, and $e_x = (1/\sqrt{2})(-\hat{x} + \hat{y})$, $e_y = (1/\sqrt{6})(\hat{x} + \hat{y} - 2\hat{z})$, and $e_z = (1/\sqrt{3})(\hat{x} + \hat{y} + \hat{z})$ for the QDs on the (111) surface. $\hat{x}$, $\hat{y}$, and $\hat{z}$ are the unit vectors along the [001], [010], and [001] crystal directions. The transition dipole moments between different electron and HH or LH states are given as

$$\langle \Psi_{\rm e}^+ | e.p | \Psi_{\rm hh}^+ \rangle = -\frac{\pi}{\sqrt{2}}(e_x + {\rm i}e_y), \tag{20}$$

$$\langle \Psi_{\rm e}^- | e.p | \Psi_{\rm hh}^- \rangle = \frac{\pi}{\sqrt{2}}(e_x - {\rm i}e_y), \tag{21}$$

$$\langle \Psi_{\rm e}^+ | e.p | \Psi_{\rm hh}^- \rangle = 0, \tag{22}$$

$$\langle \Psi_{\rm e}^- | e.p | \Psi_{\rm hh}^+ \rangle = 0, \tag{23}$$

$$\langle \Psi_{\rm e}^+ | e.p | \Psi_{\rm lh}^+ \rangle = 0 + \frac{2\pi}{\sqrt{6}}(e_z), \tag{24}$$

$$\langle \Psi_{\rm e}^+ | e.p | \Psi_{\rm lh}^- \rangle = \frac{\pi}{\sqrt{6}}(e_x - {\rm i}e_y) + 0, \tag{25}$$

$$\langle \Psi_{\rm e}^{-}|e.p|\Psi_{\rm lh}^{+}\rangle = -\frac{\pi}{\sqrt{6}}(e_x + {\rm i}e_y) + 0, \tag{26}$$

$$\langle \Psi_{\rm e}^{-}|e.p|\Psi_{\rm lh}^{-}\rangle = 0 + \frac{2\pi}{\sqrt{6}}(e_z), \tag{27}$$

where $\pi = \langle s|p_x|X\rangle = \langle s|p_x|Y\rangle = \langle s|p_z|Z\rangle$.[1]

In terms of $\Psi_{\rm hh}$ and $\Psi_{\rm lh}$ notations, the HH-LH mixed hole states can be written as

$$\Psi_{\rm h}^{\pm} = \beta\Psi_{\rm hh}^{\pm} \mp \sqrt{(1-\beta^2)}\exp(\mp {\rm i}\phi)\Psi_{\rm lh}^{\mp}. \tag{28}$$

Using the above equations, the transition dipole moments of the excitonic states in low symmetry QDs (C_1, C_S, C_{2V}) are

$$|1\rangle = \frac{1}{\sqrt{2}}\left(\Psi_{\rm e}^{+}\Psi_{\rm h}^{-} - \Psi_{\rm e}^{-}\Psi_{\rm h}^{+}\right) = {\rm i}\left(\frac{2\pi}{\sqrt{3}}\sqrt{1-\beta^2}\sin\phi\right)e_z,$$

$$|2\rangle = \frac{1}{\sqrt{2}}\left(\Psi_{\rm e}^{+}\Psi_{\rm h}^{-} + \Psi_{\rm e}^{-}\Psi_{\rm h}^{+}\right) = \left(\frac{2\pi}{\sqrt{3}}\sqrt{1-\beta^2}\cos\phi\right)e_z,$$

$$|3\rangle = \frac{1}{\sqrt{2}}\left(\Psi_{\rm e}^{+}\Psi_{\rm h}^{+} + \Psi_{\rm e}^{-}\Psi_{\rm h}^{-}\right)$$

$$= -{\rm i}\pi\left[(k\sqrt{1-\beta^2}\sin\phi)e_x + (\beta + k\sqrt{1-\beta^2}\cos\phi)e_y\right],$$

$$|4\rangle = \frac{1}{\sqrt{2}}\left(\Psi_{\rm e}^{+}\Psi_{\rm h}^{+} - \Psi_{\rm e}^{-}\Psi_{\rm h}^{-}\right)$$

$$= \pi\left[(-\beta + k\sqrt{1-\beta^2}\cos\phi)e_x - (k\sqrt{1-\beta^2}\sin\phi)e_y\right],$$

where $k = 1/\sqrt{3}$. The oscillator strengths of exciton transitions are directly proportional to the transition dipole moments of these states.

Three-dimensional polar plots of the transition dipole moments of these exciton states for $\phi = 0$ are given in Fig. 1. In this case the state $|1\rangle$ is always dark and not shown in Fig. 1. The state $|2\rangle$ is dark for pure HH but becomes bright with the direction of polarization along the growth direction of the QD for nonzero LH character in the hole states. Its intensity increases with an increase in the LH character in the hole states. The states $|3\rangle$ and $|4\rangle$ are equally bright for the pure HH. However, for the nonzero LH character in the hole state these exciton states have different intensities. The intensity of $|3\rangle$ decreases, while that of $|4\rangle$ increases with an increase in the LH in hole states. For 75% LH character in the hole state the exciton state $|3\rangle$ becomes dark. For pure LH states the $|3\rangle$ and $|4\rangle$ states are equally bright in the growth-plane and $|2\rangle$ is strongly bright along the growth direction of the QD.

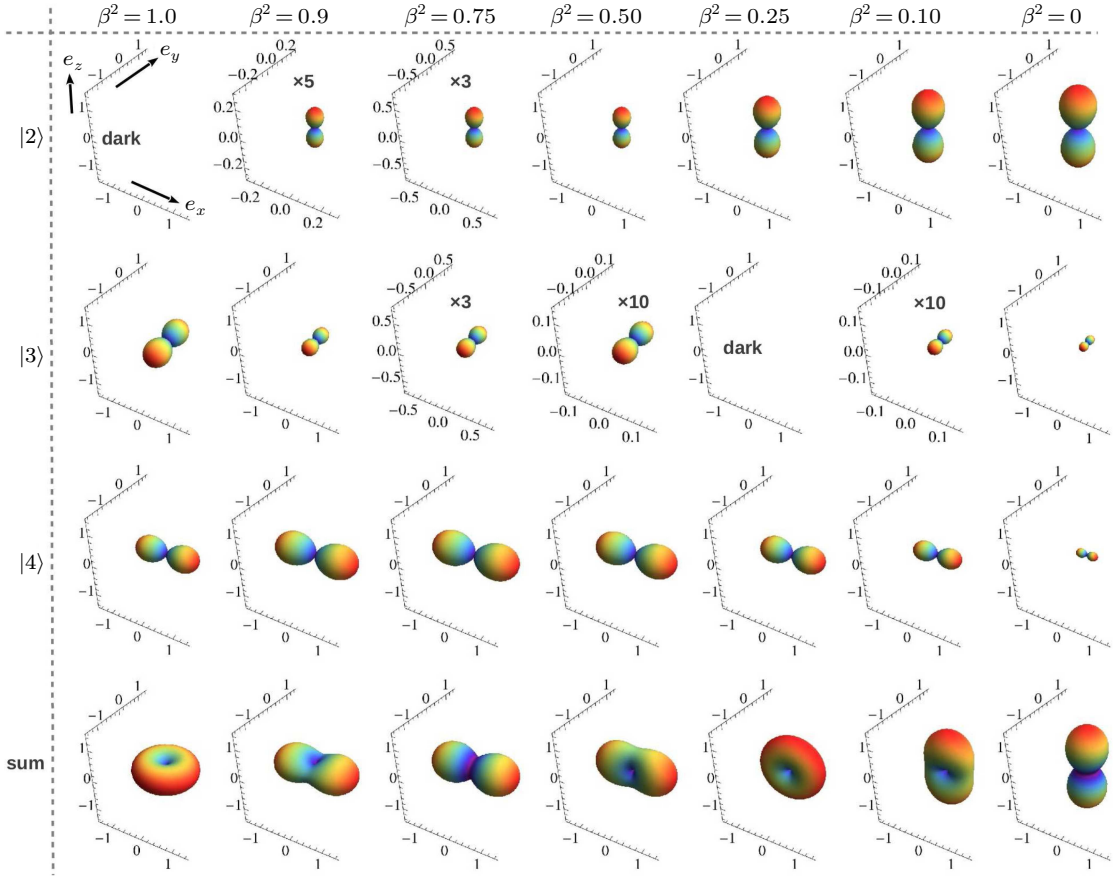


Fig. 1. (color online) Three-dimensional polar plots of transition dipole moments of different states of ground state exciton for different magnitudes of LH–HH mixing with $\phi = 0$ in a QD. The sum represents the total oscillator strength of $|1\rangle + |2\rangle + |3\rangle + |4\rangle$. In the case of $\phi = 0$, the state $|1\rangle$ is always dark for any value of β^2. The state $|2\rangle$ becomes bright with the direction of polarization along the growth direction of the QD for a nonzero value of the LH character in the hole states. The intensity of this state increases with an increase in the LH character. The $|3\rangle$ and $|4\rangle$ states are equally bright with directions of polarization in the growth-plane of the QD. However, for the case of mixed HH-LH hole states the intensity of $|3\rangle$ decreases, while that of $|4\rangle$ increases with an increase in LH character. Eventually the $|3\rangle$ becomes dark for about 75% LH character in the hole state. For further increase in LH character the $|3\rangle$ becomes bright again. For pure LH the $|3\rangle$ and $|4\rangle$ are equally bright with the directions of polarization normal to each other in the growth-plane, while $|2\rangle$ is strongly bright with the direction of polarization along the growth direction of the QD.

For $\phi \neq 0$, no exciton state is dark for nonzero magnitude of LH–HH mixing. Three-dimensional polar plots of the transition dipole moments of ground state exciton for $\phi = 30^\circ$ are given in Fig. 2. The exciton states $|1\rangle$ and $|2\rangle$ are bright with different intensities along the growth direction of the QD. The intensities of these states increase with an increase in the LH character in the hole states. The $|3\rangle$ and $|4\rangle$ are bright with different intensities in the growth-plane of the QD.

The states $|1\rangle$ and $|2\rangle$ are always dark in the growth-plane for any value of LH–HH mixing. These states become bright only along the growth direction of the QD in the presence of LH–HH mixing. The states $|3\rangle$ and $|4\rangle$ are bright in the growth-plane but always dark along the growth direction of the QD. Thus, $|1\rangle$ and $|2\rangle$ have no contribution to the polarization anisotropy

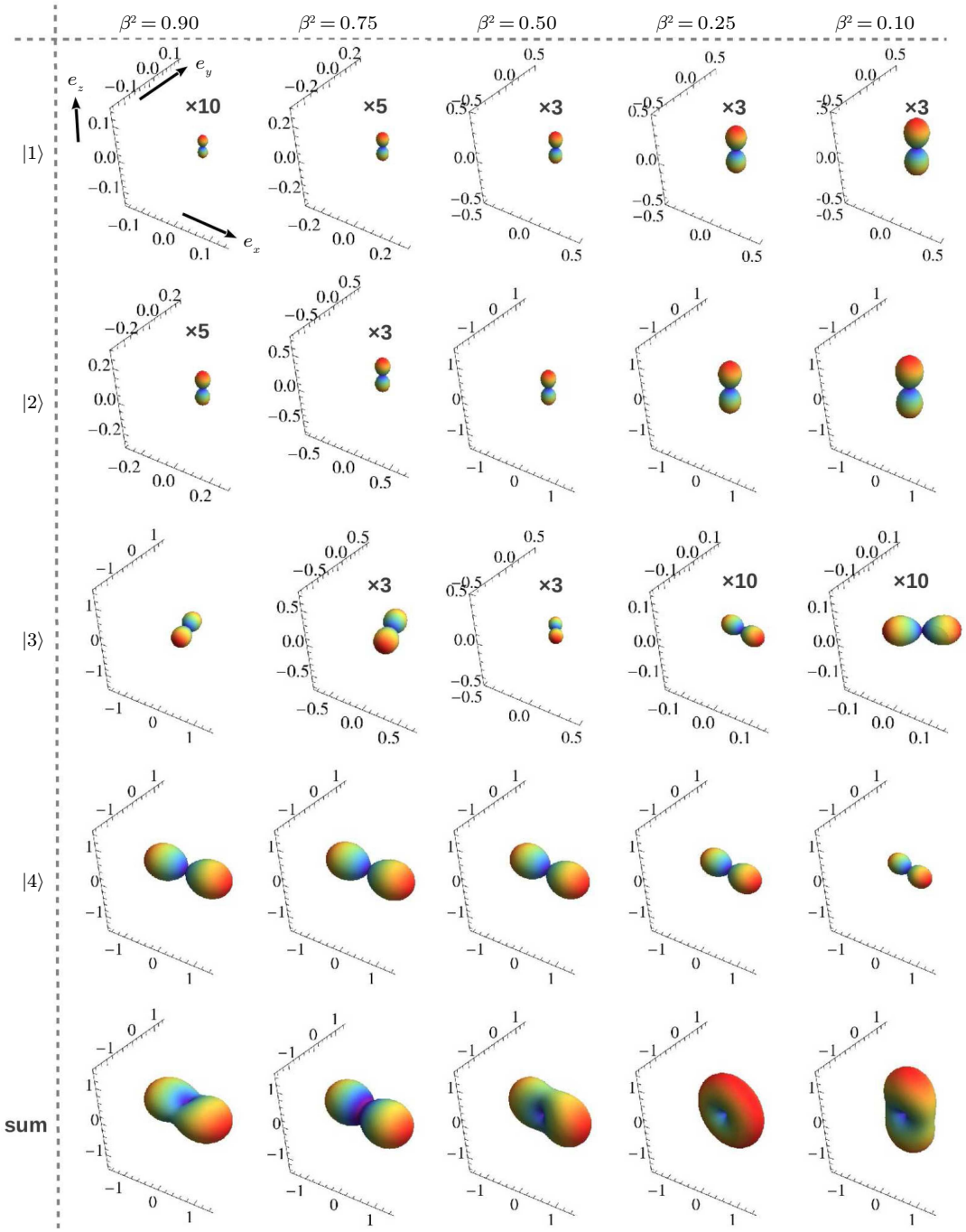


Fig. 2. (color online) Three-dimensional polar plots of transition dipole moments of different energy states of ground state exciton for different magnitudes of LH–HH mixing and $\phi = 30^\circ$. The sum represents the total oscillator strength of $|1\rangle + |2\rangle + |3\rangle + |4\rangle$. For $\phi \neq 0$, no exciton state is completely dark for nonzero magnitude of LH–HH mixing. The $|1\rangle$ and $|2\rangle$ are bright with different intensities and directions of polarization along the growth direction of the QD. The intensities of these states increase with an increase in the LH character of the hole state. The $|3\rangle$ and $|4\rangle$ are bright with different intensities and directions of polarization not normal to each other in the growth-plane of the QD. The intensity of $|3\rangle$ decreases with increasing LH character and becomes completely dark when the LH character is 75%. Thereafter it increases again with an increase in LH character in the hole state. The intensity of $|4\rangle$ decreases with an increase in LH character in the hole state.

of exciton emissions measured in experiments in the growth-plane of QD samples.[14,34] Two-dimensional polar plots of dipole moments of exciton states $|3\rangle$ and $|4\rangle$ are given in Fig. 3. In the case of pure HH or LH, the $|3\rangle$ and $|4\rangle$ are bright equally along the e_y and e_x directions, respectively. However, in the presence of LH–HH mixing the $|3\rangle$ and $|4\rangle$ have different intensities. The fluctuations in the directions of polarization of emitted exciton photons[10] can be attributed to the fluctuation in the LH–HH mixing in the QDs in a given sample. The polarization anisotropy of exciton emissions measured in experiments is defined as the degree of linear polarization (DLP) given as

$$\mathrm{DLP} = \frac{I_{\max} - I_{\min}}{I_{\max} + I_{\min}} = \frac{2k\beta\sqrt{1-\beta^2}}{\beta^2 + k^2(1-\beta^2)}, \tag{29}$$

which is independent of ϕ but depends upon the magnitude of LH–HH mixing. Thus, DLP is strongly dependent upon the HH-LH mixing. For 75% LH character in the hole state the DLP in the growth-plane becomes 100%. It is due to the fact that for 75% LH character in the hole state there is only one bright exciton transition in the growth-plane, while the other bright transition is along the growth direction of the QD. The large values of polarization anisotropy of excitonic emissions in experimental QDs[14] can be attributed to the LH–HH mixing in hole states.

For the high symmetry QDs (C_{3V}, D_{2d}), the transition dipole moments of the excitonic states are

$$\begin{aligned}
|1\rangle &= \frac{1}{\sqrt{2}}\left(\Psi_{\mathrm{e}}^{+}\Psi_{\mathrm{h}}^{-} - \Psi_{\mathrm{e}}^{-}\Psi_{\mathrm{h}}^{+}\right) = \mathrm{i}\left(\frac{2\pi}{\sqrt{3}}\sqrt{1-\beta^2}\sin\phi\right)e_z,\\
|2\rangle &= \frac{1}{\sqrt{2}}\left(\Psi_{\mathrm{e}}^{+}\Psi_{\mathrm{h}}^{-} + \Psi_{\mathrm{e}}^{-}\Psi_{\mathrm{h}}^{+}\right) = \left(\frac{2\pi}{\sqrt{3}}\sqrt{1-\beta^2}\cos\phi\right)e_z,\\
|3\rangle &= \Psi_{\mathrm{e}}^{+}\Psi_{\mathrm{h}}^{+}\\
&= \frac{\pi}{\sqrt{2}}\left(-\beta - k(\sqrt{1-\beta^2})\mathrm{e}^{-\mathrm{i}\phi}\right)e_x\\
&\quad - \mathrm{i}\frac{\pi}{\sqrt{2}}\left(\beta - k(\sqrt{1-\beta^2})\mathrm{e}^{-\mathrm{i}\phi}\right)e_y,\\
|4\rangle &= \Psi_{\mathrm{e}}^{-}\Psi_{\mathrm{h}}^{-}\\
&= \frac{\pi}{\sqrt{2}}\left(\beta - k(\sqrt{1-\beta^2})\mathrm{e}^{\mathrm{i}\phi}\right)e_x\\
&\quad - \mathrm{i}\frac{\pi}{\sqrt{2}}\left(\beta + k(\sqrt{1-\beta^2})\mathrm{e}^{\mathrm{i}\phi}\right)e_y.
\end{aligned}$$

The states $|1\rangle$ and $|2\rangle$ have polarizations along the growth direction of the QD, while $|3\rangle$ and $|4\rangle$ states have polarization in the growth-plane. The DLP as defined above for these high symmetry QDs will be

$$\mathrm{DLP} = \frac{I_{\max} - I_{\min}}{I_{\max} + I_{\min}} = 0.$$

It means the polarization anisotropy in the growth-plane in these high symmetry QDs is vanishing irrespective of any amount of LH–HH mixing. It has also been shown by the atomistic empirical pseudo-potential approach.[12] Thus, for high symmetry QDs the FSS and polarization anisotropy is vanishing, irrespective of any amount of LH–HH mixing.

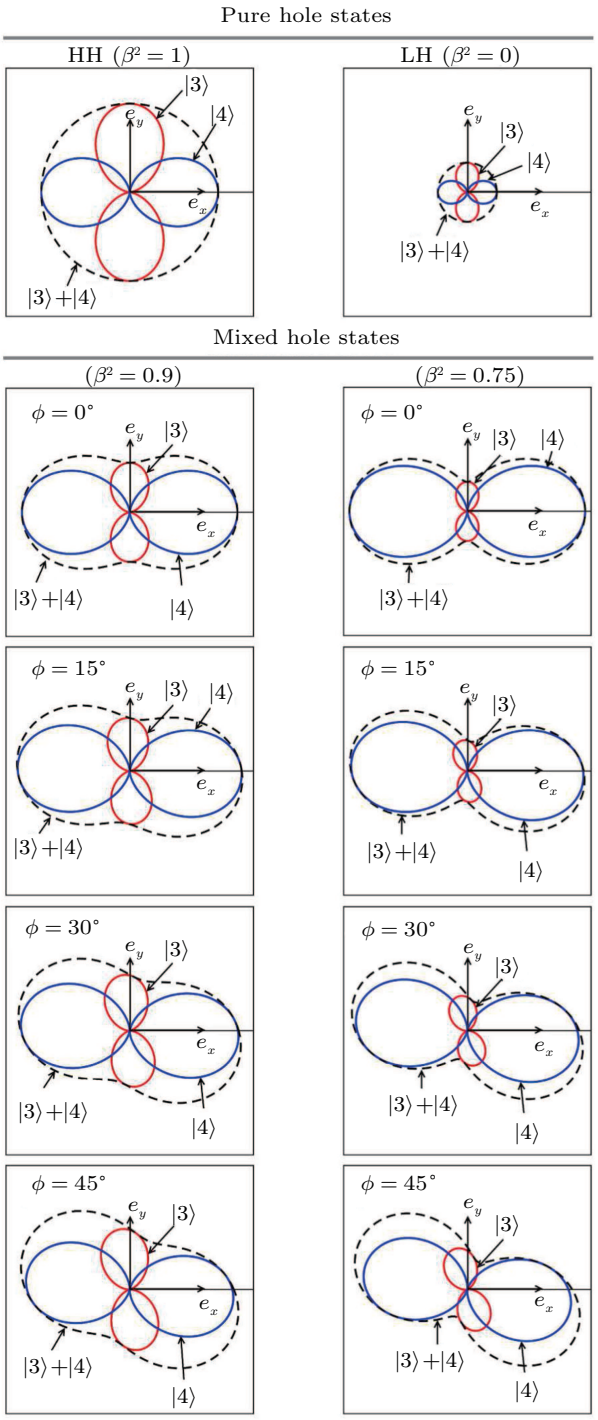


Fig. 3. (color online) Two-dimensional polar plots of bright exciton transitions in the growth-plane of a QD. The intensity of the bright exciton states for pure HH states is 3-times larger than that for pure LH states. For pure HH or LH hole states, the exciton states are along e_x and e_y in the growth-plane. However, in the case of the mixed HH-LH hole states the exciton transitions no longer have directions of polarization along e_x and e_y but rather are rotated by some angle with respect to e_x and e_y.

4. Conclusions

We investigated the effect of HH-LH mixing on the FSS, polarization anisotropy, and directions of polarization of ex-

citon emissions in semiconductor QDs using a simple mesoscopic model. We find that FSS, polarization anisotropy, and directions of polarization of exciton emissions in low symmetry QDs are sensitive to the LH–HH mixing. In high symmetry QDs, the FSS and polarization anisotropy in the growth-plane are vanishing irrespective of LH–HH mixing. For pure HH states, two higher energy states of the ground state exciton are equally bright in the growth-plane of the QD and other two lower energy states are completely dark. For pure LH states, the lowest energy state of the ground state exciton is dark, the other two intermediate energy states are equally bright in the growth-plane and the highest energy state is strongly bright along the growth direction of the QD. In the case of low symmetry QDs, for mixed LH–HH states with nonzero phase angle (ϕ) of mixing, two exciton states are bright with different intensities in the growth-plane and the other two are bright with different intensities along the growth direction of the QD. For $\phi = 0$, the lowest energy state of the ground state exciton is completely dark for any magnitude of LH–HH mixing, while other intermediate energy states become dark for 75% LH character in hole states. The polarization anisotropy of the ground state exciton emissions in the growth-plane is independent of the phase angle of LH–HH mixing but strongly depends on the magnitude of LH–HH mixing.

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