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A polaron in a quantum well within an electric field

Chuan-Yu Chen†‡, Shi-Dong Liang‡ and Ming Li§

† Centre of Theoretical Physics, CCAST (World Laboratory), PO Box 8730, Beijing 100080, People's Republic of China

‡ Department of Physics, Guangzhou Teacher's College, Guangzhou 510400, People's Republic of China

§ Department of Physics, South China Normal University, Guangzhou 510631, People's Republic of China

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Abstract. The ground-state energy and effective mass of a polaron in the quantum well (QW) of an AlAs/GaAs double heterostructure (DHS) are calculated as functions of the well width and the strength of the electric field applied along the growth direction. Every type of optical phonon mode that can exist in the DHS within the continuous model is considered separately. It is found that the contribution of the interface phonon to the polaron effect is much larger than that of the confined bulk phonon modes in QWs with width $d < 50$ Å and then falls quickly with increasing well width up to $d \simeq 200$ Å. An electric field has hardly any influence on polaron effects in QWs with width $d < 50$ Å, but when $d > 50$ Å the influence becomes much stronger. The correction of the ground-state energy and effective mass originating from LO modes decreases and that from the interface modes increases with increasing strength of the electric field.

1. Introduction

In recent years there has been considerable interest in polaron effects in quasi-two-dimensional (Q2D) semiconductor systems such as heterostructures, double heterostructures (DHSs) or quantum wells (QWs), and superlattices made up of polar materials [1–6]. The interactions between electrons and longitudinal optical (LO) phonons may modify obviously the electronic properties, e.g., the mobility of electrons, in a Q2D system. These interactions were considered based on the usual bulk Frohlich interactions in calculations of electron scattering rates [7, 8], 2D polaron cyclotron resonances [9, 10] and magnetophonon effects [11, 12]. In most of these works the phonon modes were assumed to be the same as those in bulk materials. However, the optical phonon modes may be influenced by the presence of interface(s) in Q2D systems. It has been observed in experiments that the optical phonon modes in a Q2D system are rather different from those in bulk materials [13–16]. Recently, more and more attention is being paid to the interface optical phonon modes in Q2D systems [1–3, 6]. The optical phonon modes, including bulk LO modes and interface modes, in Q2D systems have been studied by several groups [1, 2]. It was found that the bulk LO modes in a DHS are confined in the central layer or semi-confined in the side layers, and the interface modes divide into two kinds: symmetric and anti-symmetric with respect to the plane of the QW's centre, and each of them splits into two branches. The interactions between an electron and phonons with confined bulk LO modes or interface modes were also established, based on the continuum model [18]. Using these interactions calculations on QWs have been carried out [3, 17].

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In most of the relevant works a QW was treated as an infinitely high barrier well. Although this gave correct limits when well width $d \to 0$ or ∞, the well is not realistic. In fact, the barriers of some kinds of QW may be only a few hundreds of meV, not sufficiently large when compared with the lowest-level energies in a narrow QW. Thus, a finite-barrier well is more likely.

In an external magnetic field a magnetophonon resonance occurs in a QW due to the emission and absorption of optical phonons [19]. From this resonance richer information about the electron–optical-phonon interactions in heterostructures can be obtained. An external electric field also has some effects on a Q2D system. For example, the luminescence of a QW can be decreased sensitively or even completely quenched by an electric field [20, 21]. The influence of an external electric field on the polaron effects in a QW is our interest in this paper. This was considered for the bound polaron in [4]; however, only the bulk LO phonon modes were considered. In this paper, we will include the interface phonon modes as well as the confined bulk LO phonon modes and calculate the self-trapped energy and effective mass of a polaron in an AlAs/GaAs DHS within an electric field applied along the growth direction of the DHS. It is found that the contribution of the interface phonon to the polaron effect is much larger than that of the confined bulk phonon modes in a QW with width $d < 50$ Å and then falls quickly with increasing well width up to $d \simeq 200$ Å. An electric field has hardly any influence on polaron effects in a QW with width $d < 50$ Å, but when $d > 50$ Å the influence becomes much stronger. The correction of the ground-state energy and effective mass originating from LO modes decreases and that from the interface modes increases with increasing strength of the electric field.

2. Theoretical formulas

Consider an electron confined in the QW of width d in an AlAs/GaAs DHS with an electric field applied in the z direction, i.e. perpendicular to the well layers. The origin of the z axis is chosen at the centre of the well, the space for $|z| < d/2$ is filled with polar crystal 1 (GaAs) and that for $|z| > d/2$ with polar crystal 2 (AsAl). It is convenient to introduce 2D vectors $\boldsymbol{k}$ as a phonon's momentum and $\boldsymbol{\rho}$ as an electron's position so that $\boldsymbol{k}_{\mathrm{ph}} = (\boldsymbol{k}, q)$ and $\boldsymbol{r} = (\boldsymbol{\rho}, z)$. Clearly, q is quantized in the central layer of the DHS with discrete value $q = m\pi/d$, $m = 1, 2, 3, \ldots$, and remains continuous in the side layers. The Hamiltonian of the system can be written, within the framework of the effective-mass approximation, as [1, 4]

$$H = H_{\mathrm{e}} + H_{\mathrm{ph}} + H_{\mathrm{e-ph}}. \tag{1}$$

The first two terms of (1) represent the unperturbed Hamiltonian

$$H_{\mathrm{e}} = p^2/2m_{\mathrm{e}} + eFz + V(z) \tag{2}$$

$$V(z) = \begin{cases} 0 & |z| < d/2 \\ V_0 & |z| > d/2 \end{cases} \tag{2a}$$

$$H_{\mathrm{ph}} = \sum_{\sigma,\boldsymbol{k}} \hbar\omega_{\sigma}\{a_{\sigma}^{+}(\boldsymbol{k})a_{\sigma}(\boldsymbol{k}) + \tfrac{1}{2}\} \tag{3}$$

where σ denotes the phonon mode: Lν ($\nu = 1$ or 2), sj ($j = 1$ or 2) or aj ($j = 1$ or 2). Lν stands for the confined bulk LO mode in material ν, sj and aj express the symmetric and anti-symmetric interface modes of branch j, respectively. The dispersion relations of the symmetric and anti-symmetric interface modes were given by Mori and Ando in [1]. m_e is the electron band mass, which is assumed to be the same in both materials. F is the strength of the electric field. The last term of (1) stands for the Hamiltonian of electron–phonon interactions, which is given by [1]

$$H_{\text{e-ph}} = H_{\text{e-LO}} + H_{\text{e-IN}} \tag{4}$$

where

$$\begin{aligned} H_{\text{e-LO}} = -\sum_{k} \mathrm{e}^{-\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{\rho}} \Bigg\{ &\sum_{m=1,3,\ldots} B_m(k) \cos\left(\frac{m\pi z}{d}\right) \{a_m(-\boldsymbol{k}) + a_m^+(\boldsymbol{k})\} \\ &+ \sum_{m=2,4,\ldots} B_m(k) \sin\left(\frac{m\pi z}{d}\right) \{a_m(-\boldsymbol{k}) + a_m^+(\boldsymbol{k})\} \Bigg\} \end{aligned} \tag{5a}$$

$$|B_m(k)|^2 = \frac{1}{Ad} \frac{4\pi e^2 \hbar\omega_{\text{L1}}}{k^2 + (m\pi/d)^2} \left(\frac{1}{\varepsilon_{\infty 1}} - \frac{1}{\varepsilon_{01}}\right) \tag{5b}$$

for $|z| < d/2$ and

$$H_{\text{e-LO}} = -\sum_{k} \sum_{q>0} \mathrm{e}^{-\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{\rho}} B_q(k) \sin(q|z| - qd/2)\{a_q(-\boldsymbol{k}) + a_q^+(\boldsymbol{k})\} \tag{6a}$$

$$|B_q(k)|^2 = (1/AD)[4\pi e^2 \hbar\omega_{L2}/(k^2 + q^2)] \left(\frac{1}{\varepsilon_{\infty 2}} - \frac{1}{\varepsilon_{02}}\right) \tag{6b}$$

for $|z| > d/2$; here A is the area of the interface and the thickness D is assumed to be infinite at the end of the calculation; and

$$\begin{aligned} H_{\text{e-IN}} = -\sum_{k,j} \mathrm{e}^{-\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{\rho} - k(|z|-d/2)} \Big\{ &B_{sj}(k)\{a_{sj}(-\boldsymbol{k}) + a_{sj}^+(\boldsymbol{k})\} \\ &- \operatorname{sign}(z) B_{aj}(k)\{a_{aj}(-\boldsymbol{k}) + a_{aj}^+(\boldsymbol{k})\} \Big\} \end{aligned} \tag{7a}$$

for $|z| > d/2$ and

$$\begin{aligned} H_{\text{e-IN}} = -\sum_{k,j} \mathrm{e}^{-\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{\rho}} \Bigg\{ &B_{sj}(k) \frac{\cosh(kz)}{\cosh(kd/2)} \{a_{sj}(-\boldsymbol{k}) + a_{sj}^+(\boldsymbol{k})\} \\ &- B_{aj}(k) \frac{\sinh(kz)}{\sinh(kd/2)} \{a_{aj}(-\boldsymbol{k}) + a_{aj}^+(\boldsymbol{k})\} \Bigg\} \end{aligned} \tag{7b}$$

for $|z| > d/2$; here

$$|B_{sj}(k)|^2 = (\pi e^2/Ak)\hbar\omega_{sj}(k)/[\bar{\varepsilon}_1 \tanh(kd/2) + \bar{\varepsilon}_2] \tag{8a}$$

$$|B_{aj}(k)|^2 = (\pi e^2/Ak)\hbar\omega_{aj}(k)/[\bar{\varepsilon}_1 \coth(kd/2) + \bar{\varepsilon}_2] \tag{8b}$$

and $\operatorname{sign}(z) = 1$ for $z > 0$ or -1 for $z < 0$. $\bar{\varepsilon}_\nu(\omega)$ is defined as

$$1/\bar{\varepsilon}_\nu(\omega) = 1/[\varepsilon_\nu(\omega) - \varepsilon_{0\nu}] - 1/[\varepsilon_\nu(\omega) - \varepsilon_{\infty\nu}] \qquad \nu = 1, 2 \tag{9}$$

with the dielectric functions of materials ν given by

$$\varepsilon_\nu(\omega) = \varepsilon_{\infty\nu}(\omega_{L\nu}^2 - \omega^2)/(\omega_{T\nu}^2 - \omega^2). \tag{10}$$

We now proceed to calculate the ground-state energy and effective mass of the polaron by perturbation theory. The unperturbed ground-state wave function is taken to be

$$|\Psi_0\rangle = |K, 0, n_{k\sigma}\rangle = (1/\sqrt{A})e^{iK_\parallel \cdot \rho} f(z) V(\beta) |n_{k\sigma}\rangle \tag{11}$$

where $K_\parallel$ is the momentum in the xy plane, and $V(\beta)$ is the variational function $\exp\{-\beta(z + d/2)\}$, introduced by Mendez *et al* [21], with a variational parameter β, and $f(z)$ is the electron ground-state wave function in the z direction:

$$f(z) = \begin{cases} C(\beta)\cos(k_z d/2)e^{-k_z'(|z|-d/2)} & |z| > d/2 \\ C(\beta)\cos(k_z z) & |z| < d/2 \end{cases} \tag{12}$$

where $C(\beta)$ is the normalization constant, which is given by

$$C(\beta)^2 = 2\beta\{\beta k_z'\{\cos(k_z d) + 1\}\cosh(\beta d) + k_z'^2\{\sinh(\beta d) + X\} \\ + \beta^2\{\cos(k_z d) + 1\}\sinh(\beta d)\}^{-1}(k_z'^2 - \beta^2)e^{\beta d} \tag{13}$$

where

$$X = \beta^2\{\cos(k_z d)\sinh(\beta d) + (k_z/\beta)\sin(k_z d)\cosh(\beta d)\}(\beta^2 + k_z^2)^{-1}. \tag{13a}$$

k_z and k_z' in (12) are related to the electron subband energy E_0:

$$k_z = \sqrt{2m_e E_0}/\hbar \qquad k_z' = \sqrt{2m_e(V_0 - E_0)}/\hbar. \tag{14}$$

Here E_0 is the lowest of the subband energy levels determined by the transcendental equation

$$E_l = V_0\cos^2[(d/2)\sqrt{2m_e E_l}/\hbar] \qquad l = 0, 1, 2, \ldots. \tag{15}$$

In equation (11) $|n_{k\sigma}\rangle$ expresses an n-phonon state with wave vector k and phonon mode σ. The expectation value $\langle\Psi_0|H_e + H_{ph}|\Psi_0\rangle$ gives the unperturbed energy of the ground state:

$$E_{Ke}^{(0)}(\beta) = \frac{\hbar^2(K_\parallel^2 + K_\perp^2)}{2m_e} + \langle eFz\rangle + \sum_{k,\sigma}(n + \tfrac{1}{2})\hbar\omega_{k\sigma} \tag{16}$$

where

$$K_\perp^2(\beta) = (k_z^2 - \beta^2) + C(\beta)^2\cos^2\left(\frac{k_z d}{2}\right)k_z'\beta\left\{\frac{1}{k_z' - \beta} - \frac{e^{-2\beta d}}{k_z' + \beta}\right\} \\ - C(\beta)^2 e^{-\beta d}\frac{k_z\beta}{k_z^2 + \beta^2}\{k_z\cos(k_z d)\sinh(\beta d) - \beta\sin(k_z d)\cosh(\beta)\} \tag{16a}$$

$$\langle eFz\rangle = \langle H_F\rangle_+ + \langle H_F\rangle_- + \langle H_F\rangle_0 \tag{16b}$$

$$\langle H_F\rangle_\pm = [\pm eFC(\beta)^2\cos^2(k_z d/2)/4(k_z' \pm \beta)^2]\{(k_z' \pm \beta)d + 1\} \tag{16c}$$

$$\langle H_F\rangle_0 = \tfrac{1}{4}eFC(\beta)^2 e^{-\beta d}\{(1/\beta^2)\sinh(\beta d) - (d/\beta)\cosh(\beta d) + Q\} \tag{16d}$$

with

$$Q = \left\{ \frac{k_z^2 - \beta^2}{k_z^2 + \beta^2} \cos(k_z d) \sinh(\beta d) - \frac{2k_z \beta}{k_z^2 + \beta^2} \sin(k_z d) \cosh(\beta d) + k_z d \sin(k_z d) \sinh(\beta d) + \beta d \cos(k_z d) \cosh(\beta d) \right\} \frac{1}{-2(k_z^2 + \beta^2)}$$

where $K_\perp$ is the electron momentum in the z-axis direction and $\langle eFz \rangle$ is the potential energy. By minimizing $E_{Ke}^{(0)}(\beta)$ we obtain a parameter $\beta_{\min}$ at the minimum value of $E_{Ke}^{(0)}(\beta)$ for each electric field strength. In the weak-coupling limit, the electron–phonon interaction energy can be calculated by second-order perturbation theory. Thus,

$$E_{Ke} - E_{Ke}^{(0)} = \sum_{k,\sigma} \frac{|\langle K_\parallel - k, 0, 1_{k\sigma} | H_{\text{e-ph}} | K_\parallel, 0, 0_{k\sigma} \rangle|^2}{E_{K_\parallel}^{(0)} - E_{K_\parallel - k}^{(0)} - \hbar\omega_{k\sigma}} = \Delta E_1 + \Delta E_2 + \Delta E_3 \tag{17}$$

where ΔE_1 and ΔE_2 are the energy corrections due to the electron interaction with the bulk LO modes confined in the well and in the barrier, respectively. ΔE_3 is the contribution from the interface phonon modes. More explicitly, these are given by

$$\Delta E_1 = -\alpha_{L1} \hbar\omega_{L1} - (\hbar^2 K^2 / 2m_e) \gamma_{L1} \tag{18a}$$

$$\Delta E_2 = -\alpha_{L2} \hbar\omega_{L2} - (\hbar^2 K^2 / 2m_e) \gamma_{L2} \tag{18b}$$

$$\Delta E_3 = -\sum_{\sigma j} \int_0^\infty \mathrm{d}k\, \alpha_{\sigma j}(k) \hbar\omega_{\sigma j}(k) - \frac{\hbar^2 K^2}{2m_e} \gamma_{\text{IN}} \qquad \sigma = \text{s, a.} \tag{18c}$$

The α and γ parameters are the Frohlich-type coupling constants with phonon modes indicated by their subscripts. They are given by

$$\alpha_{L1} = \frac{4\alpha_{F1}}{K_{p1} d} \sum_{m=1}^{\infty} G_m^2 \ln\left(\frac{m\pi}{K_{p1} d}\right) \left\{ \left(\frac{m\pi}{K_{p1} d}\right)^2 - 1 \right\}^{-1} \tag{19a}$$

$$\gamma_{L1} = \frac{2\alpha_{F1}}{K_{p1} d} \sum_{m=1}^{\infty} G_m^2 \left\{ 1 - \left(\frac{m\pi}{K_{p1} d}\right)^2 \right\}^{-3} \left[4 \left(\frac{m\pi}{K_{p1} d}\right)^2 \ln\left(\frac{m\pi}{K_{p1} d}\right) + 1 - \left(\frac{m\pi}{K_{p1} d}\right)^4 \right] \tag{19b}$$

$$\alpha_{L2} = \frac{2\alpha_{F2}}{\pi K_{p2}} \int_0^\infty \mathrm{d}q\, \{G_+(q) + G_-(q)\}^2 \ln\left(\frac{q}{K_{p2}}\right) \left[\left(\frac{q}{K_{p2}}\right)^2 - 1 \right]^{-1} \tag{20a}$$

$$\gamma_{L2} = \frac{\alpha_{F2}}{\pi K_{p2}} \int_0^\infty \mathrm{d}q\, \{G_+(q) + G_-(q)\}^2 \left[\left(\frac{q}{K_{p2}}\right)^2 - 1 \right]^{-3} \times \left[\left(\frac{q}{K_{p2}}\right)^4 - \left(\frac{q}{K_{p2}}\right)^2 \ln\left(\frac{q}{K_{p2}}\right) - 1 \right] \tag{20b}$$

$$\alpha_{\sigma j}(k) = \alpha_{F\sigma j}(k) \frac{P_\sigma^2(k)}{K_{p\sigma j}(k)} \left[1 + \left(\frac{k}{K_{p\sigma j}(k)}\right)^2 \right]^{-1} \qquad \sigma = \text{s, a} \tag{21a}$$

$$\gamma_{\text{IN}} = 2 \int_0^\infty \mathrm{d}k \sum_{\sigma j} P_\sigma^2(k) \frac{\alpha_{F\sigma j}(k) k^2 K_{p\sigma j}^3(k)}{\{k^2 + K_{p\sigma j}^2(k)\}^3} \qquad \sigma = \text{s, a} \tag{21b}$$

where

$$K_{p\nu} = \sqrt{2m_e\omega_{L\nu}/\hbar} \qquad \nu = 1, 2 \tag{22a}$$

$$K_{p\sigma j}(k) = \sqrt{2m_e\omega_{\sigma j}(k)/\hbar} \qquad \sigma = s, a \tag{22b}$$

$$\alpha_{F\nu} = (e^2 K_{p\nu}/2\hbar\omega_{L\nu})(1/\varepsilon_{\infty\nu} - 1/\varepsilon_{0\nu}) \qquad \nu = 1, 2 \tag{23a}$$

$$\alpha_{F\sigma j}(k) = \frac{e^2 K_{p\sigma j}(k)}{2\hbar\omega_{\sigma j}(k)} \frac{1}{\bar{\varepsilon}_1 \mathrm{tc}(kd/2) + \bar{\varepsilon}} \qquad \mathrm{tc}(x) = \begin{cases} \tanh(x) & \text{for } \sigma = s \\ \coth(x) & \text{for } \sigma = a \end{cases}. \tag{23b}$$

The functions G_m, $G_\pm(k)$ and $P_\sigma(k)$ are expressed as

$$G_{2n} = \tfrac{1}{2}C(\beta)^2 e^{-\beta d}(S_1 + A_1^+ - A_1^-) \tag{24a}$$

$$G_{2n+1} = \tfrac{1}{2}C(\beta)^2 e^{-\beta d}(S_2 + A_2^+ + A_2^-) \tag{24b}$$

$$G_\pm(q) = \frac{C(\beta)^2 \cos(k_z d/2)}{q^2 + 4(k_z' \pm \beta)^2} q e^{-(\beta\pm\beta)d} \tag{24c}$$

$$P_\sigma(k) = Q_-(k) \pm Q_+(k) \pm \{R_-(k) \pm R_+(k)\} \begin{cases} 1/\cosh(kd/2) & \text{'+'}, \sigma = s \\ 1/\sinh(kd/2) & \text{'−'}, \sigma = a \end{cases} \tag{24d}$$

where

$$S_1 = \{(m\pi/d)^2 + 4\beta^2\}^{-1}(2m\pi/d)\cos(m\pi/2)\sinh(\beta d) \tag{25a}$$

$$S_2 = \{(m\pi/d)^2 + 4\beta^2\}^{-1}(2m\pi/d)\sin(m\pi/2)\cosh(\beta d) \tag{25b}$$

$$A_1^\pm = \tfrac{1}{2}(b_\pm^2 + \beta^2)^{-1}\{b_\pm \cos(b_\pm d)\sinh(\beta d) - \beta\sin(b_\pm d)\cosh(\beta d)\} \tag{25c}$$

$$A_2^\pm = \tfrac{1}{2}(b_\pm^2 + \beta^2)^{-1}\{b_\pm \sin(b_\pm d)\cosh(\beta d) - \beta\cos(b_\pm d)\cosh(\beta d)\} \tag{25d}$$

(here $b_\pm = k_z \pm m\pi/2d$)

$$Q_\pm = \{C(\beta)^2\cos^2(k_z d/2)/[2(k_z' \pm \beta) + k]\}e^{-(\beta\pm\beta)d} \tag{25e}$$

$$R_\pm(k) = \tfrac{1}{4}C(\beta)^2 e^{-\beta d}\left\{\frac{\sinh(\beta_\pm d)}{\beta_\pm} + \frac{1}{\beta_\pm^2 + k_z^2}\{\beta_\pm \sinh(\beta_\pm d)\cos(k_z d) + k_z\cosh(\beta_\pm d)\sin(k_z d)\}\right\}. \tag{25f}$$

(here $\beta_\pm = \beta \pm k/2$). By combining (16)–(18) the ground-state energy of the polaron in a QW within an electric field can be written as

$$E_{Ke} = \frac{\hbar^2 K^2}{2m_e^*} + \frac{\hbar^2 K_\perp^2}{2m_e} + \sum_{k,\sigma}(n + \tfrac{1}{2})\hbar\omega_{k\sigma} + \langle eFz\rangle - \Delta E. \tag{26}$$

$$m_e^* = m_e/(1 - \gamma_{L1} - \gamma_{L2} - \gamma_{IN}) \tag{27}$$

is called the effect mass of the polaron. Since the mass in the z direction (m_e) is different from that in the xy plane (m_e^*), it is indicated that the effective mass of this Q2D system is of non-isotropic nature due to the electron–phonon interaction. The last term $-\Delta E$ in (26) is the self-energy, i.e. the correction of the system's ground-state energy; the self-trapped energy is given by

$$\Delta E = \alpha_{L1}\hbar\omega_{L1} + \alpha_{L2}\hbar\omega_{L2} + \sum_{\sigma,j}\int_0^\infty \mathrm{d}k\, \alpha_{\sigma j}(k)\hbar\omega_{\sigma j}(k). \tag{28}$$

3. Conclusions and discussion

Now we use the equations (27) and (28) to calculate the effective mass and self-trapped energy of a polaron in the GaAs QW of a GaAs/AlAs DHS within an electric field. Throughout this paper, the parameters used in our numerical work are as given in table 1. We take the effective mass of an electron as $m_e = 0.067m_0$ in the DHS, where m_0 is the mass of a free electron.

Table 1.

	$\hbar\omega_{LO}$ (meV)	$\hbar\omega_{TO}$ (meV)	ε_0	ε_∞	V_0 (eV)
GaAs	36.2	33.3	12.5	10.06	1.3
AlAs	50.1	44.8	10.6	8.16	

We plot the values of the self-trapped energies due to the confined bulk LO phonon and interface phonon versus the electric field strength for several different QW widths $d = 30$ Å, 50 Å, 80 Å, 100 Å, 200 Å, 400 Å and 600 Å, as shown in figures 1 and 2, respectively. The effects of the electron–phonon interaction on the total effective mass correction are brought out respectively by LO and interface phonon mode contributions. We also plot the LO and SO effective mass correction versus the electric field strength for these different well widths as shown in figures 3 and 4.

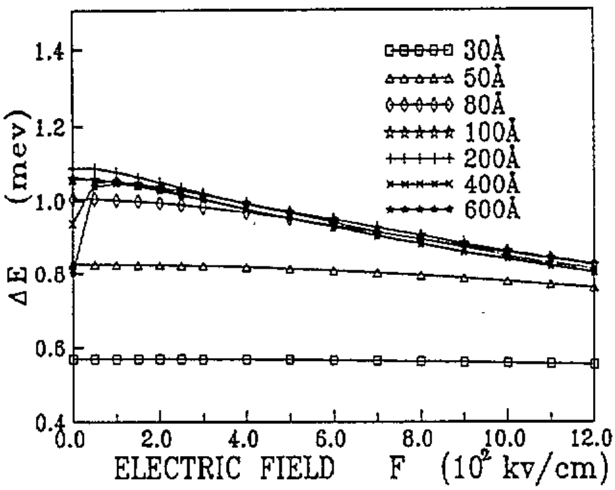

Figure 1. The self-trapped energy from confined bulk LO phonons in a QW with different well widths versus the strength of the electric field.

Intuitively, the electric field shifts the wave peak of an electron in the DHS from position $z = 0$ to $z < 0$. The electron wave function tends to concentrate near the left barrier of the DHS. The most important factor is that the interaction between an electron and interface phonons is enhanced since the overlap of the electron wave on the interface increases. Thus, the contributions of interface phonons increase with strengthening electric field. For a narrow QW, as much of the electron wave spreads into the left barrier, and the contribution of material 2 becomes large enough, a small value of β is sufficient to balance the electric

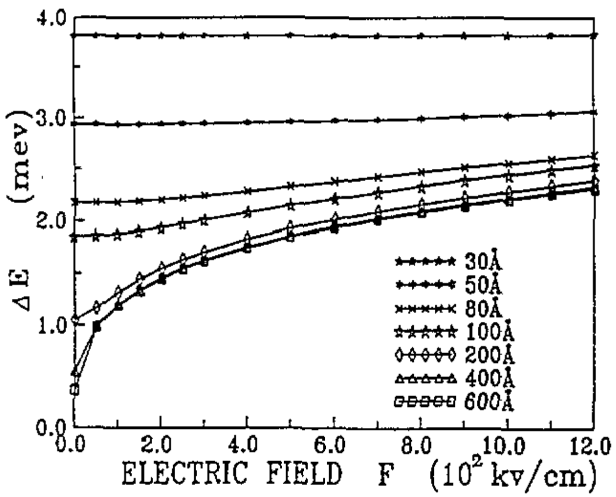

Figure 2. The self-trapped energy from interface phonons in a QW with different well widths versus the strength of the electric field.

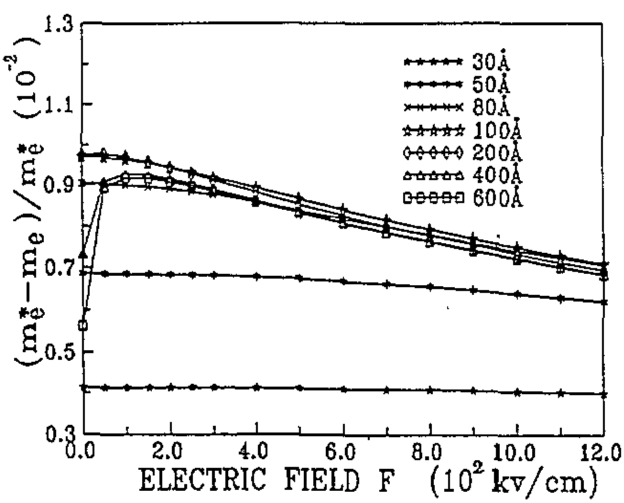

Figure 3. The effective mass correction from confined bulk LO phonons in a QW with different well widths versus the strength of the electric field.

field. Thus, an electric field has hardly any influence on the electron–phonon interactions in a narrow QW. From figures 1–4 we can see that for a narrow QW with well width $d < 50$ Å the self-trapped energy and effective mass scarcely vary even at a considerable electric field strength, and for a wide QW; however, the variation becomes more and more obvious with increasing electric field strength. Within a more intense electric field (e.g., 1200 kV cm^{-1}) the contribution of the interface mode may be much larger than that of the confined LO mode. The most interesting thing is that the lines for the QW with well width larger than about 100 Å in each figure are almost overlapping except for their beginning parts. It seems that the polaron effects in QWs with different well widths (> 100 Å) are nearly the same.

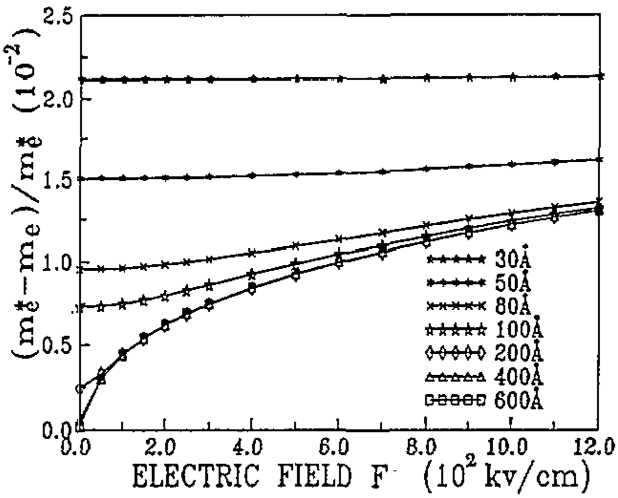

Figure 4. The effective mass correction from interface phonons in a QW with different well widths versus the strength of the electric field.

Meanwhile, we also find that the beginning part of each line for a QW with more than 100 Å width rises faster and faster as the well width increases and the rest of the line becomes more and more horizontal.

It should be pointed out that recently Ridley [22] studied the interaction between electron and LO, TO and IP (interface polariton) modes in the QW and discussed the mixing effect of LO, TO and IP modes. He proposed to use a new quasiparticle, the hybridon, to describe the mixing effect, and studied the electron–hybridon scattering problem in the infinite-depth QW. In this paper, considering both LO and IO (interface optical) phonon modes in the finite-depth QW we study the influence of electric field on the electron character in the ground state. Certainly, to discuss this system in more detail, we may study the influence of both the mixing effect of various modes and the electric field on the electron character in the ground state based on Ridley's work [22] and this work, which may be done in a future publication.

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