Magnetopolarons in lateral surface superlattices with periodically structured interfaces

1993 J. Phys.: Condens. Matter 5 7451

(http://iopscience.iop.org/0953-8984/5/40/019)

J. Phys.: Condens. Matter 5 (1993) 7451–7460. Printed in the UK

Magnetopolarons in lateral surface superlattices with periodically structured interfaces

Hong Sun

Centre of Theoretical Physics, CCAST (World Laboratory), Beijing 100080, People's Republic of China and Department of Physics and Institute of Condensed Matter Physics, Shanghai Jiao Tong University, Shanghai 200030, People's Republic of China

Received 14 April 1993, in final form 30 June 1993

Abstract. An improved Wigner–Brillouin theory, developed by Peeters and Devreese, was extended to the case of lateral surface superlattices (LSSLs), where lateral periodic potentials due to the periodically structured interfaces are introduced into two-dimensional electronic systems. Numerical calculations were carried out for the energy dispersions of the magnetopolarons in GaAs LSSL structures with periodic potentials along the lateral x direction, such as those produced by deposition of AlAs and GaAs fractional layers on (001) vicinal GaAs substrates. The degenerate magnetopolaron energy levels become dependent on the electron wave vector k_y in the y direction, forming magnetopolaron bands. For weak magnetic fields ($\omega_c \ll \omega_{LO}$), the effect of the lateral periodic potentials on the electron–LO-phonon interaction energy is negligible, the effect being mainly on the Landau levels of the free electrons in the LSSLs. For strong magnetic fields ($\omega_c \geqslant \omega_{LO}$), the effect of the lateral periodic potentials on the electron–LO-phonon interaction energy must be considered in order to obtain correct magnetopolaron energy dispersions. The transition energy between the first two magnetopolaron energy levels forms a wide band which will broaden greatly the cyclotron resonance absorption spectra of the LSSLs.

1. Introduction

Polaron states in two-dimensional electronic systems with magnetic fields applied normally to the systems, known as two-dimensional magnetopolarons, have attracted much attention from both theorists and experimentalists of condensed matter physics in the last ten years. Extensive theoretical works on the problem have been performed by Lindemann *et al* [1], Das Sarma [2], Peeters and Devreese [3] and Larsen [4], for example. Very recently, Müller *et al* have experimentally observed the resonant inter-Landau-level tunnellings in a GaAs/AlAs superlattice with the applied magnetic field as high as 18 T [5]. With the rapid development of crystal growth technology, it is now possible to fabricate quasi-two-dimensional electron systems with periodic potentials in the lateral directions [6–14]. If the lateral confining potentials are not strong enough to prevent electrons from travelling in the lateral directions, the electronic systems are referred to as lateral surface superlattices (LSSLs), otherwise parallel quantum well wires or quantum dot arrays are formed.

A variety of interesting phenomena associated with the electronic energy band structures of LSSLs have been observed experimentally. Magnetoresistance oscillations with magnetic fields in high-mobility GaAs/$Ga_{1-x}Al_xAs$ LSSLs have been observed, which show a different behaviour from the well-known Shubnikov–de Hass oscillations [15–17]. Strong anisotropies in the ratios of the electron–light-hole-exciton peak intensities to the electron–heavy-hole-exciton peak intensities of photoluminescence excitation spectra have been

0953-8984/93/407451+10$07.50

reported in GaAs/AlAs LSSLs [18–20]. Exciton–polariton localizations in GaAs/$Ga_{1-x}Al_xAs$ LSSLs with the width of the lateral interface structures below 150 nm have been demonstrated [21]. In addition, a series of peaks in the differential resistance of LSSLs fabricated on modulation doped GaAs/$Ga_{1-x}Al_xAs$ heterostructures have been observed in the region where the Fermi energy is of the same order of magnitude as the potential modulation, as a consequence of an oscillatory density of electronic states in the LSSLs [22].

Of the various structures proposed for the LSSLs, the one produced by deposition of AlAs and GaAs fractional layers on (001) vicinal GaAs substrates seems to offer the greatest potential for wide application in microelectronics and electrooptics [18, 19, 23]. This is because of the possibility of producing large periodic structures on LSSL interfaces with periods $L_x = 100 \sim 300$ Å comparable with the electron wavelengths and the large band offsets between GaAs and AlAs materials which enhance the effect of interface structures on the motions of the electrons in the LSSLs. Indeed, electronic wave interference devices using the GaAs/AlAs LSSLs have been designed [24].

The magnitude of the lateral periodic potential caused by the fluctuation of the width of the LSSL can be estimated simply by [25]

$$\Delta V_0 = (\hbar^2/2m_e)(\pi/L_z)^2(2\Delta L_z/L_z)$$

where L_z and ΔL_z are the averaged width and width fluctuation due to the periodically structured interfaces of the LSSL, respectively. For GaAs LSSLs, when $\Delta L_z/L_z$ is $10\% \sim 20\%$, ΔV_0 is about $10 \sim 20$ meV which is not negligibly small compared with the LO phonon energy $\hbar\omega_{LO} = 36.8$ meV in GaAs. In the resonant region, where $\omega_c = eB/m_e c \approx \omega_{LO}$, we have for the GaAs LSSLs $B \approx 20$ T, which corresponds to a magnetic length $l_0 = \sqrt{\hbar c/eB} = 57$ Å. For a sine-shaped lateral periodic potential

$$\Delta V(x) = \Delta V_0 \sin(2\pi x/L_x)$$

our calculation shows that the energy spacing of the first two Landau levels of the free electrons in the GaAs LSSLs are mostly affected by the periodic potential $\Delta V(x)$ when $(\pi l_0/L_x) \approx 1$, which requires $L_x \approx 200$ Å just in the possible region of the lateral periods of the GaAs/AlAs LSSLs. It is expected that the lateral periodic potentials will change significantly the Landau levels of the free electrons and so change the electron–LO-phonon interaction energies in GaAs LSSLs in the resonant region. In this paper, we report the calculated magnetopolaron energies in GaAs LSSLs with periodically structured interfaces.

2. Theory

In a previous paper [26], we proposed a method of calculating the electronic subband structures in LSSLs with periodically structured interfaces, where a coordinate transformation was found to transform the LSSL structures to quantum wells with planar interfaces plus a lateral periodic potential which takes a rather complicated form. In this paper, to facilitate the calculation we adopt a much simpler treatment. For a model calculation, the LSSL is described by a quantum well with planar interfaces and an additional periodic potential in the lateral x direction due to the width fluctuations caused by the periodically structured interfaces. The lateral periodic potential is given approximately by [25]

$$\Delta V(x) = -\frac{\hbar^2}{2m_e}\left(\frac{\pi}{L_z}\right)^2 \frac{2\Delta L_z(x)}{L_z} = \Delta V_0 \sin(2\pi x/L_x) \quad (1)$$

where $\Delta L_z(x)$ is the width fluctuation of the LSSL caused by the periodically structured interfaces. In the effective mass approximation, the Hamiltonian of the electron–LO-phonon interaction system in a normally applied static magnetic field $\boldsymbol{B} = (0, 0, B)$ reads

$$H = H_0 + H_1 \tag{2}$$

with

$$H_0 = E_z + \frac{1}{2m_e}\left(\boldsymbol{p} + \frac{e}{c}\boldsymbol{A}(x)\right)^2 + \Delta V(x) + \sum_{\boldsymbol{k}} \hbar\omega_{\mathrm{LO}} a_{\boldsymbol{k}}^{+} a_{\boldsymbol{k}} \tag{3}$$

and

$$H_1 = \sum_{\boldsymbol{k}} [V_{\boldsymbol{k}} \exp(\mathrm{i}\boldsymbol{q}\cdot\boldsymbol{\rho}) a_{\boldsymbol{k}} + V_{\boldsymbol{k}}^{*} \exp(-\mathrm{i}\boldsymbol{q}\cdot\boldsymbol{\rho}) a_{\boldsymbol{k}}^{+}] \tag{4}$$

where m_e is the electron band mass in GaAs, $\boldsymbol{\rho} = (x, y)$ and $\boldsymbol{p} = -\mathrm{i}\hbar(\partial/\partial x, \partial/\partial y)$ are the in-plane components of the electron position and momentum operators, $\boldsymbol{A}(x) = (0, Bx, 0)$ is the vector potential of the magnetic field in the Landau gauge, $\boldsymbol{k} = (\boldsymbol{q}, k_z)$ is the LO-phonon wave vector with $\boldsymbol{q}$ the in-plane component, and the electron–LO-phonon interaction coefficient is given by

$$V_{\boldsymbol{k}} = \hbar\omega_{\mathrm{LO}}\left(\frac{4\pi\alpha}{V}\right)^{1/2}\left(\frac{\hbar}{2m_e\omega_{\mathrm{LO}}}\right)^{1/4}\frac{1}{k}\langle\varphi_0(z) \mid \mathrm{e}^{\mathrm{i}k_z z} \mid \varphi_0(z)\rangle \tag{5}$$

with α the electron–LO-phonon coupling constant and $\varphi_0(z)$ the ground-state wavefunction in the z direction determined by the following equation:

$$\left(-\frac{\hbar^2}{2m_e}\frac{\mathrm{d}^2}{\mathrm{d}z^2} + V_0(z)\right)\varphi_0(z) = E_z\varphi_0(z). \tag{6}$$

The electron potential $V_0(z)$, which confines the electron in the two-dimensional well of the LSSL, is assumed to be infinitely high on the interfaces between the well and barrier of the LSSL. In writing the Hamiltonian H, (2)–(4), we have implicitly neglected the differences of the lattice dynamical properties of the well and barrier, so a bulk LO-phonon system is considered. This approximation has been used by many authors to study magnetopolarons in two-dimensional electronic systems [2–4].

The eigen wavefunction of H_0 is taken as

$$\psi_{nk_y}(\boldsymbol{\rho}) = \frac{\exp(-\mathrm{i}k_y y)}{\sqrt{L_y}}\varphi_{nk_y}(x)|ph\rangle. \tag{7}$$

Because of the translational symmetry of H_0 in the y direction, the dependence of the wavefunction on the coordinate y is a plane wave. $|ph\rangle$ are the free LO-phonon states. $\varphi_{nk_y}(x)$ satisfies the equation

$$\left(-\frac{\hbar^2}{2m_e}\frac{\mathrm{d}^2}{\mathrm{d}x^2} + \frac{m_e\omega_c^2}{2}(x - x_{ky})^2 + \Delta V(x)\right)\varphi_{nk_y}(x) = \varepsilon_n(k_y)\varphi_{nk_y}(x) \tag{8}$$

with $x_{k_y} = \hbar c k_y/eB$. For the case where $\Delta V(x) = 0$, equation (8) is just the eigenvalue equation of a harmonic oscillator. When the periodic potential $\Delta V(x)$ is not very strong, such as $\Delta V(x)/\hbar\omega_c \ll 1$, equation (8) can be solved with the perturbation expansion theory, which gives

$$\begin{aligned}\varepsilon_n(k_y) = (n + \tfrac{1}{2})\hbar\omega_c &+ \langle\varphi_n^{(0)}(x - x_{k_y}) \mid \Delta V(x) \mid \varphi_n^{(0)}(x - x_{k_y})\rangle \\ &+ \sum_{m \neq n}^{\infty} \frac{|\langle\varphi_m^{(0)}(x - x_{k_y}) \mid \Delta V(x) \mid \varphi_n^{(0)}(x - x_{k_y})\rangle|^2}{(n - m)\hbar\omega_c} \\ &+ \ldots\end{aligned} \tag{9}$$

and

$$\begin{aligned}\varphi_{nk_y}(x) = \varphi_n^{(0)}(x - x_{k_y}) &\\ &+ \sum_{m \neq n}^{\infty} \frac{\langle\varphi_m^{(0)}(x - x_{k_y}) \mid \Delta V(x) \mid \varphi_n^{(0)}(x - x_{k_y})\rangle}{(n - m)\hbar\omega_c} \varphi_m^{(0)}(x - x_{k_y}) \\ &+ \ldots\end{aligned} \tag{10}$$

where $\varphi_n^{(0)}(x)$ is the nth eigen wavefunction of a harmonic oscillator.

In the Rayleigh–Schrödinger (RS) theory [1, 3], the energy levels of the magnetopolarons in the LSSLs are given to the second-order energy correction by

$$E_n^{\mathrm{RS}}(k_y) = E_z + \varepsilon_n(k_y) + \Delta E_n^{\mathrm{RS}}(k_y) \tag{11}$$

with

$$\Delta E_n^{\mathrm{RS}}(k_y) = -\sum_q \sum_{m=0}^{\infty} |V_q|^2 \frac{|\langle\varphi_{m,k_y+q_y}(x) \mid \mathrm{e}^{-\mathrm{i}q_x x} \mid \varphi_{n,k_y}(x)\rangle|^2}{\varepsilon_m(k_y + q_y) - \varepsilon_n(k_y) + \hbar\omega_{\mathrm{LO}}} \tag{12}$$

where $|V_q|^2 = \sum_{k_z} |V_k|^2$. For GaAs LSSLs, because of the weak couplings between the electrons and LO phonons with $\alpha = 0.07$ and $|V_q|^2 \propto \alpha$, in the calculation of $\Delta E_n^{\mathrm{RS}}(k_y)$ we can take the first-order approximation for $\varepsilon_n(k_y)$ and the zeroth-order approximation for $\varphi_{n,k_y}(x)$. From equations (9) and (10), we have

$$\begin{aligned}\varepsilon_n(k_y) &= (n + \tfrac{1}{2})\hbar\omega_c + \delta\varepsilon_n(k_y) \\ &= (n + \tfrac{1}{2})\hbar\omega_c + \langle\varphi_n^{(0)}(x - x_{k_y}) \mid \Delta V(x) \mid \varphi_n^{(0)}(x - x_{k_y})\rangle\end{aligned} \tag{13}$$

and

$$\varphi_{nk_y}(x) = \varphi_n^{(0)}(x - x_{k_y}). \tag{14}$$

RS theory works well in the weak-field region where $\omega_c \ll \omega_{\mathrm{LO}}$. But in the resonant region where $\omega_c \approx \omega_{\mathrm{LO}}$, $\Delta E_n^{\mathrm{RS}}(k_y)$ goes to infinity for $n \neq 0$. A modified Wigner–Brillouin theory [1, 3], referred to as 'improved Wigner–Brillouin' (IWB) theory in the literature, has been proposed to study the magnetopolarons in two-dimensional electronic systems with planar interfaces. Here we extend the IWB theory to the case of the LSSLs. We redivide the Hamiltonian H (2) into

$$H = H_0' + H_1' = [H_0 + \Delta E_0^{\mathrm{RS}}(-p_y/\hbar)] + [H_1 - \Delta E_0^{\mathrm{RS}}(-p_y/\hbar)] \tag{15}$$

where $\Delta E_0^{\mathrm{RS}}(k_y)$ is the ground-state electron–LO-phonon interaction energy determined by equation (12), and $p_y = -i\hbar\partial/\partial y$ is the y component of the electron momentum operator. The eigen wavefunctions of H_0' are the same as those of H_0 given by equation (7), with the eigenenergies equal to $E_z + \varepsilon_n(k_y) + \Delta E_0^{\mathrm{RS}}(k_y) + n\hbar\omega_c$. The energy levels of the magnetopolarons of the LSSLs in the extended IWB theory are given to the second-order energy correction by

$$E_n(k_y) = E_z + \varepsilon_n(k_y) + \Delta E_n(k_y) \tag{16}$$

with

$$\begin{aligned}\Delta E_n(k_y) = -\sum_q \sum_{m=0}^{\infty} (|V_q|^2 |\langle \varphi_m^{(0)}(x + x_{q_y}) \mid \mathrm{e}^{-\mathrm{i}q_x x} \mid \varphi_n^{(0)}(x)\rangle|^2) \\ \times [(m-n)\hbar\omega_c + \delta\varepsilon_m(k_y + q_y) - \delta\varepsilon_n(k_y) \\ + \hbar\omega_{\mathrm{LO}} - \Delta E_n(k_y) + \Delta E_0^{\mathrm{RS}}(k_y + q_y)]^{-1}.\end{aligned} \tag{17}$$

For all the three cases we consider, where we take $\Delta V_0/\hbar\omega_c = 0.5$, and $B = 10$ T ($\omega_c \approx 0.5\omega_{\mathrm{LO}}$), $B = 20$ T ($\omega_c \approx \omega_{\mathrm{LO}}$) and $B = 30$ T ($\omega_c \approx 1.5\omega_{\mathrm{LO}}$), respectively, numerical calculation shows that the bandwidth of $\Delta E_0^{\mathrm{RS}}(k_y + q_y)$ is much less than $\hbar\omega_c$ (see figures 1, 3 and 4(*b*)). We can safely take $\Delta E_0^{\mathrm{RS}}(k_y + q_y) \approx \Delta E_0^{\mathrm{RS}}(k_y)$ in equation (17), which makes $\Delta E_0(k_y) = \Delta E_0^{\mathrm{RS}}(k_y)$. The electron–LO-phonon interaction energy in the extended IWB theory is further simplified to

$$\begin{aligned}\Delta E_n(k_y) = -\sum_q \sum_{m=0}^{\infty} (|V_q|^2 |\langle \varphi_m^{(0)}(x + x_{q_y}) \mid \mathrm{e}^{-\mathrm{i}q_x x} \mid \varphi_n^{(0)}(x)\rangle|^2) \\ \times [(m-n)\hbar\omega_c + \delta\varepsilon_m(k_y + q_y) - \delta\varepsilon_n(k_y) \\ + \hbar\omega_{\mathrm{LO}} - \Delta E_n(k_y) + \Delta E_0(k_y)]^{-1}.\end{aligned} \tag{18}$$

For weak lateral periodic potentials $\Delta V_0/\hbar\omega_c < 1$, $\delta\varepsilon_n(k_y)$ is much less than $\hbar\omega_c$. In the summation $\sum_m$ of equation (18), when $m - n$ is larger than a particular value M_0, which depends on the required accuracy of the numerical results, we can neglect $\delta\varepsilon_m(k_y + q_y) - \delta\varepsilon_n(k_y)$ in the denominators in equation (18). So we have

$$\begin{aligned}\Delta E_n(k_y) = -\sum_q \sum_{m=0}^{\infty} \frac{|V_q|^2 |\langle \varphi_m^{(0)}(x + x_{q_y}) \mid \mathrm{e}^{-\mathrm{i}q_x x} \mid \varphi_n^{(0)}(x)\rangle|^2}{(m-n)\hbar\omega_c + \hbar\omega_{\mathrm{LO}} - \Delta E_n(k_y) + \Delta E_0(k_y)} \\ - \sum_q \sum_{m=0}^{n+M_0} \Bigg((|V_q|^2 |\langle \varphi_m^{(0)}(x + x_{q_y}) \mid \mathrm{e}^{-\mathrm{i}q_x x} \mid \varphi_n^{(0)}(x)\rangle|^2) \\ \times [(m-n)\hbar\omega_c + \delta\varepsilon_m(k_y + q_y) - \delta\varepsilon_n(k_y) \\ + \hbar\omega_{\mathrm{LO}} - \Delta E_n(k_y) + \Delta E_0(k_y)]^{-1} \\ - \frac{|V_q|^2 |\langle \varphi_m^{(0)}(x + x_{q_y}) \mid \mathrm{e}^{-\mathrm{i}q_x x} \mid \varphi_n^{(0)}(x)\rangle|^2}{(m-n)\hbar\omega_c + \hbar\omega_{\mathrm{LO}} - \Delta E_n(k_y) + \Delta E_0(k_y)} \Bigg).\end{aligned} \tag{19}$$

The first infinite summation, which is the electron–LO-phonon interaction energy in a two-dimensional electronic system with planar interfaces, can be converted into an integral, as first discovered by Peeters and Devreese [3]. The second summation, which describes the effects of the lateral periodic potential on the electron–LO-phonon interaction energy, is finite and can be calculated numerically.

3. Results and discussion

We have carried out numerical calculations of the energy levels of magnetopolarons, equations (16) and (19), below the LO-phonon continuum, for GaAs LSSLs where $m_e = 0.0665m_0$ and $\hbar\omega_{LO} = 36.8$ meV. The normally applied magnetic field strength is $B = 20$ T ($\omega_c \approx \omega_{LO}$) in the resonant region. The structural parameters of the GaAs LSSL are $L_z = 100$ Å, $L_x = 200$ Å, and the magnitude of the lateral periodic potential is fixed at $\Delta V_0 = 0.5\hbar\omega_c$. In figure 1, the energy dispersions of the magnetopolarons, $E_n(k_y) - E_z$ (the full curves in figure 1(*a*)), and electron–LO-phonon interaction energies, $\Delta E_n(k_y)$ (the full curves in figure 1(*b*)), for the first two levels ($n = 0, 1$) are given as functions of $x_{k_y} = \hbar ck_y/eB$. The broken lines in figure 1 are the corresponding $E_n(k_y) - E_z$ and $\Delta E_n(k_y)$ of the magnetopolarons in a GaAs quantum well with planar interfaces. In all the numerical calculations carried out in this paper, we fix $M_0 = 10$. By further increasing M_0, the numerical results are changed by less than 1%.

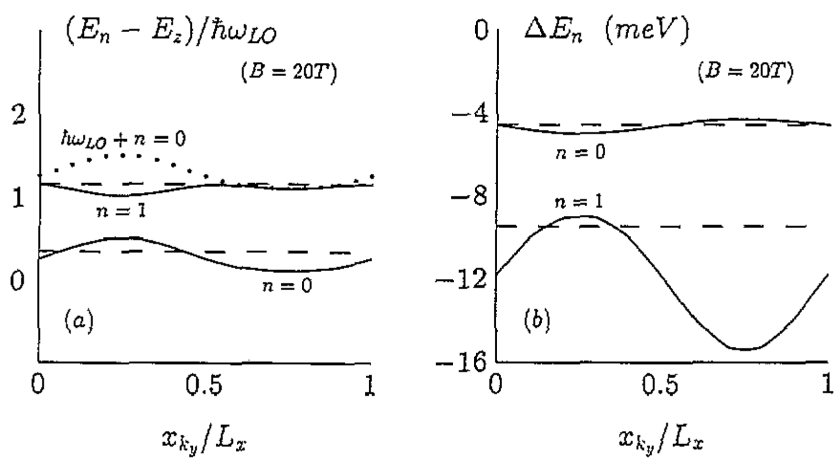

Figure 1. The energy dispersions of the magnetopolarons, $E_n(k_y) - E_z$ (the full curves in (*a*)), and electron–LO-phonon interaction energies, $\Delta E_n(k_y)$ (the full curves in (*b*)), for the first two levels ($n = 0, 1$) as functions of $x_{k_y} = \hbar ck_y/eB$ in a GaAs LSSL where $m_e = 0.0665m_0$ and $\hbar\omega_{LO} = 36.8$ meV. The normally applied magnetic field strength is $B = 20$ T ($\omega_c \approx \omega_{LO}$) in the resonant region. The structural parameters of the GaAs LSSL are $L_z = 100$ Å, $L_x = 200$ Å, with the magnitude of the sine-shaped lateral periodic potential fixed at $\Delta V_0 = 0.5\hbar\omega_c$. The broken lines are the corresponding $E_n(k_y) - E_z$ and $\Delta E_n(k_y)$ of the magnetopolarons in a GaAs quantum well with planar interfaces.

It is interesting to note that the energy dispersions of the first two magnetopolaron levels in LSSLs are out of phase, which results in a wide transition energy band, $E_1(k_y) - E_0(k_y)$, given in figure 2 as a function of x_{k_y} for $B = 20$ T (full curve). The broken line in figure 2 is the transition energy of the first two magnetopolaron levels in a GaAs quantum well with planar interfaces. It is also noteworthy from figure 1 that $E_n(k_y)$ and $\Delta E_n(k_y)$ are out of phase. The electron–LO-phonon interaction compensates for the effects of the lateral periodic potential on the free electron, especially for the $n = 1$ level. The effect of the lateral periodic potential on the electron–LO-phonon interaction energy is not negligible in the resonant region. The averaged value of the $\Delta E_1(k_y)$ curve in figure 1(*b*) is shifted downwards compared with the corresponding $n = 1$ broken line, indicating an enhancement of the electron–LO-phonon interaction energy by the lateral periodic potential. From magnetopolaron theory, it is well known that the energy levels of the magnetopolarons do not cross the energy level at $E_0(k_y) + \hbar\omega_{LO}$, which is given in figure 1(*a*) as the dotted curve. The high magnetopolaron levels ($n > 1$) lie between the $E_0(k_y) + \hbar\omega_{LO}$ and $E_1(k_y)$ curves. For the LSSL structures we have considered, in the resonant region ($B = 20$ T) all the magnetopolaron levels are pinned to the energy level at $E_0(k_y) + \hbar\omega_{LO}$ when $0.5 \leqslant x_{k_y} \leqslant 1$.

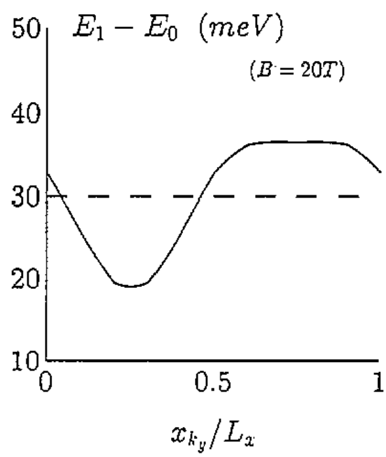

Figure 2. The transition energy, $E_1(k_y) - E_0(k_y)$ (full curve), of the first two magnetopolaron energy levels ($n = 0, 1$) in the LSSL structure described in figure 1 as a function of x_{k_y}. The broken line is the transition energy of the first two magnetopolaron levels in a GaAs quantum well with planar interfaces.

In figure 3, we give the energy dispersions of the magnetopolarons, $E_n(k_y) - E_z$ (the full curves in figure 3(*a*)), and electron–LO-phonon interaction energies, $\Delta E_n(k_y)$ (the full curves in figure 3(*b*)), for the first two levels ($n = 0, 1$) as functions of x_{k_y} in the same GaAs LSSL structure as that described in figure 1 except with $B = 10$ T and $\Delta V_0 = 0.5\hbar\omega_c$. The definition of the broken lines and dotted curves in figure 3 are the same as in figure 1. The results in figure 3(*b*) show that, in the weak magnetic field region ($\omega_c \ll \omega_{LO}$), the effect of the lateral periodic potential on the electron–LO-phonon interaction energy is very small. The effect of the lateral periodic potential on the energy dispersions of the magnetopolarons is mainly on the Landau levels of the free electron in the LSSL.

The strong magnetic field region ($\omega_c \gg \omega_{LO}$) is hard to reach in the GaAs LSSL, as it requires a magnetic field higher than 20 T. But, for a model calculation, in figure 4 we give the energy dispersions of the magnetopolarons, $E_n(k_y) - E_z$ (the full curves in figure 4(*a*)), and electron–LO-phonon interaction energies, $\Delta E_n(k_y)$ (the full curves in figure 4(*b*)), for the first two levels ($n = 0, 1$) as functions of x_{k_y} in the same GaAs LSSL structure as

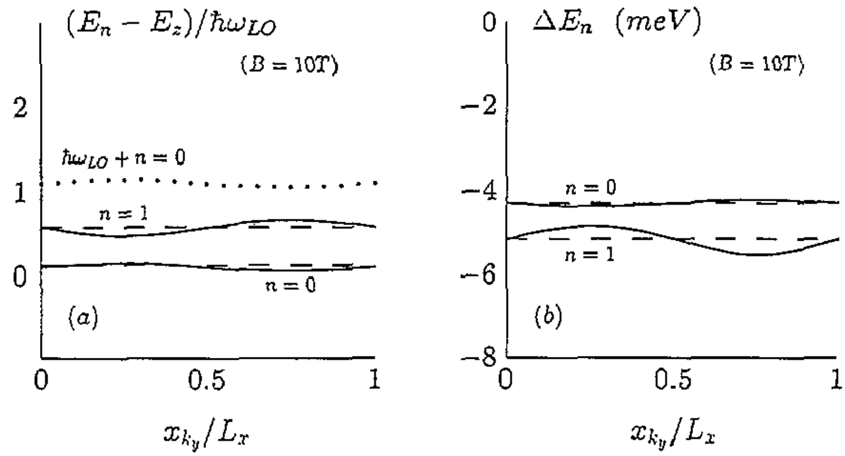

Figure 3. The energy dispersions of the magnetopolarons, $E_n(k_y) - E_z$ (the full curves in (a)), and electron–LO-phonon interaction energies, $\Delta E_n(k_y)$ (the full curves in (b)), for the first two levels ($n = 0, 1$) as functions of x_{k_y} in the same GaAs LSSL structure as that described in figure 1 except with $B = 10$ T ($\omega_c \ll \omega_{LO}$) and $\Delta V_0 = 0.5\hbar\omega_c$. The definition of the broken lines is the same as in figure 1.

that described in figure 1, except that $B = 30$ T and $\Delta V_0 = 0.5\hbar\omega_c$. The definition of the broken lines and dotted curves in figure 4 are the same as in figure 1. The results in figure 4(b) show that the electron–LO-phonon interaction energy is greatly enhanced by the lateral periodic potential for the $n = 1$ level, which shifts the magnetopolaron level $E_1(k_y)$ towards $E_0(k_y)$ (see the full curves in figure 4(a)), resulting in a wide transition energy band $E_1(k_y) - E_0(k_y)$.

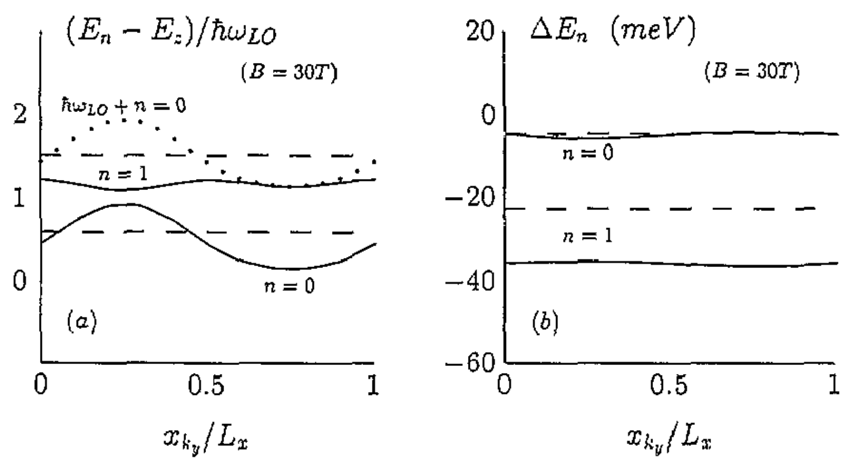

Figure 4. The energy dispersions of the magnetopolarons, $E_n(k_y) - E_z$ (the full curves in (a)), and electron–LO-phonon interaction energies, $\Delta E_n(k_y)$ (the full curves in (b)), for the first two levels ($n = 0, 1$) as functions of x_{k_y} in the same GaAs LSSL structure as that described in figure 1 except with $B = 30$ T ($\omega_c \gg \omega_{LO}$) and $\Delta V_0 = 0.5\hbar\omega_c$. The definition of the broken lines is the same as in figure 1.

To summarize, improved Wigner–Brillouin theory was extended to the case of the LSSLs, where lateral periodic potentials due to the periodically structured interfaces are introduced into the two-dimensional electronic systems. Numerical calculations were carried out for the energy dispersions of the magnetopolarons in the GaAs LSSL structures with periodic potentials along the lateral x direction, such as those produced by deposition of AlAs and GaAs fractional layers on (001) vicinal GaAs substrates. The degenerate magnetopolaron energy levels become dependent on the electron wave vector k_y in the y direction, forming magnetopolaron bands. For weak magnetic fields ($\omega_c \ll \omega_{LO}$), the effect of the lateral periodic potentials on the electron–LO-phonon interaction energy is negligible. The effect is mainly on the Landau levels of the free electron in the LSSL. For strong magnetic fields ($\omega_c \geqslant \omega_{LO}$), the effect of the lateral periodic potentials on the electron–LO-phonon interaction energy must be considered in order to obtain the correct magnetopolaron energy dispersions and the transition energies. The transition energy between the first two magnetopolaron energy levels forms a wide band which will broaden greatly the cyclotron resonance absorption spectra of the LSSL.

In our calculation, we have neglected the differences between the lattice dynamical properties of the well and barrier by considering a bulk LO-phonon system. So we have neglected the interface optical phonon modes. This approximation is fairly good for GaAs/$Ga_{1-x}Al_xAs$ LSSLs with $x < 0.3$. However, it may not be accurate enough for GaAs/GaAl LSSLs, where the interface optical phonon modes become important. But our results, as a first step calculation, do predict the general properties of magnetopolarons in LSSLs with periodically structured interfaces in strong magnetic fields ($\omega_c \geqslant \omega_{LO}$). A detailed investigation of the magnetopolaron problem in LSSLs with interface optical phonon modes and electron–electron interaction effects taken into consideration is actively under way in our group.

Acknowledgments

This work is supported by the Fok Ying Tung Education Foundation and the National Natural Science Foundation of China under Grant No 19004006.

References

[1] Lindemann G, Lassning R, Seidenbusch W and Gornil E 1983 *Phys. Rev.* B **28** 4693
[2] Das Sarma S 1984 *Phys. Rev. Lett.* **52** 859
[3] Peeters F M and Devreese J T 1985 *Phys. Rev.* B **31** 3689
[4] Larsen M D 1986 *Phys. Rev.* B **33** 799
[5] Müller W, Grahn H T, Haug R T and Ploog K 1992 *Phys. Rev.* B **46** 9800
[6] Warren A C, Antoniadis D A, Smith H I and Melngailis J 1985 *IEEE Electron. Device Lett.* **EDL-6** 294
[7] Ismail K, Chu W, Yen A, Antoniadis D A and Smith H I 1989 *Appl. Phys. Lett.* **54** 460
[8] Stormer H L, Pfeiffer L N, Baldwin K W, West K W and Spector J 1991 *Appl. Phys. Lett.* **58** 726
[9] Sakaki H, Wagatsuma K, Hamasaki J and Saito S 1976 *Thin Solid Films* **36** 497
[10] Petroff P M, Gossard A C and Wiegmann W 1984 *Appl. Phys. Lett.* **45** 620
[11] Fukui T and Saito H 1987 *Appl. Phys. Lett.* **50** 824
[12] Fukui T and Saito H 1988 *J. Vac. Sci. Technol.* B **6** 1373
[13] Gaines J M, Petroff P M, Kroemer H, Simes R J, Geels R S and English J H 1988 *J. Vac. Sci. Technol.* B **6** 1478
[14] Motohisa J, Tanaka M and Sakaki H 1989 *Appl. Phys. Lett.* **55** 1214
[15] Weiss D, Klitzing K V, Ploog K and Weimann G 1989 *Europhys. Lett.* **8** 117
[16] Gerhardts R R, Weiss D and Klitzing K V 1989 *Phys. Rev. Lett.* **62** 1173

[17] Winkler R W and Kotthaus J P 1989 *Phys. Rev. Lett.* **62** 1177
[18] Tanaka M and Sakaki H 1988 *Japan. J. Appl. Phys.* **27** L2025
[19] Tsuchiya M, Gaines J M, Yan R H, Simes R J, Holtz P O, Coldren L A and Petroff P M 1989 *Phys. Rev. Lett.* **62** 466
[20] Tanaka M and Sakaki H 1989 *Appl. Phys. Lett.* **54** 1326
[21] Kohl M, Heitmann D, Grambow P and Ploog K 1990 *Phys. Rev.* B **42** 2941
[22] Smoliner J, Rosskopf V, Berthold G, Gornik E, Böhm G and Weimann G 1992 *Phys. Rev.* B **45** 1915
[23] Sugawara H, Schulman J N and Sakaki H 1991 *J. Appl. Phys.* **69** 2722
[24] Fukui T, Tsubaki K, Saito H, Kasu M and Honda T 1992 *Surf. Sci.* **267** 588
[25] Sakaki H, Noda T, Hirakawa K, Tanaka M and Matsusue T 1987 *Appl. Phys. lett.* **51** 1934
[26] Sun H 1992 *Phys. Rev* B **46** 12371