Cyclotron resonance of two-dimensional interface polarons

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A perturbation method has been used to study the interface polaron in a magnetic field with a two-dimensional approximation. The results show that there are two splittings of the polaron cyclotron mass occurring due to the coupling of the electron with two branches of interface phonon modes. In general, one splitting caused by the phonons with higher frequencies will be more important experimentally. However, both of the two splittings might be observable for some semiconductor heterostructures, for instance GaAs/AlAs.

I. INTRODUCTION

Polaronic properties are influenced dramatically in polar-polar semiconducting heterojunctions due to the presence of their interfaces. As usual in polaron problems, the investigation of the cyclotron resonance (CR) can be performed by using the one-electron approximation at sufficiently low density of electrons.[1] The electron in a heterojunction is generally confined to a thin layer near the interface by the energy-band bending within the interface region (confinement potential). It has two dimensional (2D) character.[2] Some authors[3–5] discussed the interface polaron CR by considering the interaction between the electron and 2D bulk longitudinal-optical (LO) phonons to simplify the real electron-phonon coupling. In those works, the influence of the LO phonons was overexaggerated because of polaron dimensionality effects.

Taking into account the influences of the bulk LO phonons and a single effective interface-optical (IO)-phonon branch, Gu *et al.* discussed the ground state of an interface magnetopolaron[6] and the polaron CR (Refs. 7 and 8) in such a heterojunction. They considered the extension of the electron wave function in a direction perpendicular to the interface. Some interesting results are given to understand the fundamental properties of the IO polarons. However, their introduction of the image potential without considering the energy bending at the interface, which for heterojunctions is more important, weakens the real effects of the interface phonons. On the other hand, the single effective branch approximation adopted for the IO phonons is not satisfactory, as pointed out in our previous paper.[9]

In recent years, the real contribution from the interaction between an electron and IO phonons in heterostructures of polar-polar semiconductors has been investigated.[10–12] The results of a free polaron[9,13] and a bound polaron[14] indicated that more rigorous consideration about the interaction between the electron and two branches of IO-phonon modes should be adopted for a single heterojunction.

Several years ago, experimental investigations of $Ga_{1-x}In_xAs/InP$ ($Al_xIn_{1-x}As$) heterojunctions[15] and InAs/GaSb quantum wells[16] observed resonant splitting around the frequencies of the TO phonons of the systems. Some authors[17,18] interpreted this puzzling observation as the effect of electron–IO-phonon coupling. On the other hand, there was other opinion in contrast to this explanation.[16] More detained theoretical study and carefully arranged experimental observation, therefore, should be done.

Since the electron–IO-phonon coupling is weak enough in most III-V and II-VI semiconductive compounds which form the heterojunctions, here we adopt a perturbation method proposed by Larsen[5,7] to investigate the CR of an interface polaron in a semiconductive heterojunction, including the influence of both the two branches of IO-phonon modes at zero temperature. Without losing generality, we neglect many-particle effects such as occupation and screening effects, and assume the confined electron is a 2D one moving on the interface of the heterojunction. The polaron CR mass shows that there are two splittings around both IO-phonon frequencies. One splitting at the higher frequency, with which the electron coupling is stronger, is wider than the other one around the lower frequency. However the two splittings may seem remarkable according to carefully performed experiments for some heterojunctions such as GaAs/AlAs, InP/GaP, etc.

II. HAMILTONIAN AND PERTURBATION CALCULATION

Let us suppose a heterojunction consisting of two semi-infinite polar semiconductors labeled 1 and 2, re-

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spectively. The electron in the conductive band of material 1 moves on the interface (x-y plane) interacting with two branches of IO-phonon modes.[9,12] A static uniform magnetic field with a symmetrical gauge transformation[7] $\mathbf{B}=(0,0,B_M)$ is applied perpendicular to the interface. The Hamiltonian H of this system can be written as

$$H=H_0+H_{e\text{-IO}} , \tag{1a}$$

with

$$H_0=\frac{1}{2m}[p_x-\beta^2(y/4)]^2+\frac{1}{2m}[p_y+\beta^2(x/4)]^2 + \sum_{q\xi}\hbar\omega_\xi b^+_{q\xi}b_{q\xi} \tag{1b}$$

and

$$H_{e\text{-IO}}=\sum_{q\xi}[G_\xi(q)e^{iq\cdot\rho}b_{q\xi}+\text{H.c.}] . \tag{1c}$$

Equation (1b) is the non-interaction term of the system, and (1c) the electron–IO-phonon interaction contribution. $\mathbf{p}=(p_x,p_y,0)$ $[\boldsymbol{\rho}=(x,y,0)]$ is the 2D momentum (position) operator of the electron, $b^+_{q\xi}$ ($b_{q\xi}$) the creation (annihilation) operator of an IO phonon of the ξth branch ($\xi=+,-$) with wave vector $\mathbf{q}$ and energy $\hbar\omega_\xi$. m is the band mass of the electron. β is defined by

$$\beta^2=\frac{2eB_M}{c} . \tag{2}$$

The electron-phonon interaction coefficient in Eq. (1c) is given by

$$G_\xi(q)=-i\left[\frac{2\pi\hbar^2\omega_\xi}{Sq}\alpha_\xi\right]^{1/2}\left[\frac{\hbar\omega_\xi}{2m}\right]^{1/4} . \tag{3}$$

In Eq. (3), the dimensionless electron–IO-phonon coupling constant is

$$\alpha_\xi=\frac{e^2}{(\delta_1^2+\delta_2^2)(\hbar\omega_\xi)^2}\left[\frac{2m\hbar}{\omega_\xi}\right]^{1/2} , \tag{4a}$$

in which

$$\delta_\zeta=(\epsilon_{0\zeta}-\epsilon_{\infty\zeta})^{1/2}\omega_{T\zeta}/(\omega_{T\zeta}^2-\omega_\xi^2) \quad (\zeta=1,2) , \tag{4b}$$

where S denotes the area of the interface. $\epsilon_{0\zeta}$ ($\epsilon_{\infty\zeta}$) is the static (optical) dielectric constant, and $\omega_{T\zeta}$ is the transverse-optical (TO)-phonon frequency of the ζth material. The frequency ω_ξ of the ξth branch of phonon modes can be obtained by following equation:

$$\omega_\pm^2=\frac{b\pm\sqrt{b^2-4ac}}{2a} , \tag{4c}$$

where

$$a=\epsilon_{\infty1}+\epsilon_{\infty2} ,$$

$$b=\epsilon_{\infty1}(\omega_{L1}^2+\omega_{T2}^2)+\epsilon_{\infty2}(\omega_{L2}^2+\omega_{T1}^2) ,$$

and

$$c=\epsilon_{\infty1}\omega_{L1}^2\omega_{T2}^2+\epsilon_{\infty2}\omega_{L2}^2\omega_{T1}^2 .$$

In Eqs. (4), $\omega_{L\zeta}$ denotes the bulk LO-phonon frequency of the ζth material.

We introduce the one-dimensional (1D) harmonic oscillator operators

$$A=\frac{1}{\sqrt{\hbar}\beta}\left[\left[p_x-\beta^2\frac{y}{4}\right]-i\left[p_y+\beta^2\frac{x}{4}\right]\right] \tag{5a}$$

and

$$B=A^+-i\frac{\beta}{2\sqrt{\hbar}}(x+iy) , \tag{5b}$$

which satisfy the following commutation relations

$$[A,A^+]=[B,B^+]=1 \quad \text{and} \quad [A,B]=[A,B^+]=0 .$$

Hamiltonian (1) can be rewritten as

$$H_0=\frac{\hbar\beta^2}{2m}(A^+A+\tfrac{1}{2})+\sum_{q\xi}\hbar\omega_\xi b^+_{q\xi}b_{q\xi} \tag{6a}$$

and

$$H_{e\text{-IO}}=\sum_{q\xi}[G_\xi(q)L_q^{-1}M_q^{-1}b_{q\xi}+\text{H.c.}] , \tag{6b}$$

where

$$L_k=\exp\left[\frac{\sqrt{\hbar}}{\beta}(q_x+iq_y)A-\frac{\sqrt{\hbar}}{\beta}(q_x-iq_y)A^+\right] ,$$

$$M_k=\exp\left[\frac{\sqrt{\hbar}}{\beta}(q_x-iq_y)B-\frac{\sqrt{\hbar}}{\beta}(q_x+iq_y)B^+\right] .$$

The CR frequency of a free electron in a magnetic field is given by

$$\omega_c=\frac{\beta^2}{2m}=\frac{eB_M}{mc} . \tag{7}$$

We choose H_0 as the unperturbed Hamiltonian and $H_{e\text{-IO}}$ as the perturbed one. In the situation of weak electron-phonon coupling $\lambda_\xi\alpha_\xi\ll1$, the Landau levels are given by

$$\psi_{Mn}=(n!M!)^{-1/2}(A^+)^n|0\rangle_A(B^+)^M|0\rangle_B=|n\rangle_A|M\rangle_B , \tag{8}$$

where A^+ lowers the z angular momentum quantum number M and raises the Landau quantum number n by one unit. B^+ raises M by one unit but has no effect on n. The phonon vacuum state is described by

$$|0\rangle=\prod_\xi|0\rangle_\xi . \tag{9}$$

The unperturbed ground-state energy for the nth Landau level can be obtained by using Eqs. (6a), (8), and (9):

$$E_n^{(0)}=\langle0|_B\langle M|_A\langle n|H_0|n\rangle_A|M\rangle_B|0\rangle=\hbar\omega_c(n+\tfrac{1}{2}) . \tag{10}$$

The intermediate wave functions in the perturbation calculation are the products of ψ_{Mn} with a phonon state and can be written as

$$|n\rangle_A|M\rangle_B \prod_\xi b^+_{q_1\xi} b^+_{q_2\xi}\cdots b^+_{q_N\xi}|0\rangle_\xi \ . \tag{11}$$

Considering the perturbation of Eq. (6b), the energy correction for the nth Landau level to the second order can be derived:

$$\Delta E_n^{(2)} = \sum_{q\xi} G^2_\xi(q)|_\xi\langle q|b^+_{q\xi}|q\rangle_\xi|^2{}_B\langle M|M_q^{-1}M_q|M\rangle_B \times \sum_{n'} \frac{|_A\langle n'|L_q^{-1}|n\rangle_A|^2}{E_n^{(0)}-E_{n'}^{(0)}-\hbar\omega_\xi} \ . \tag{12}$$

From (10) and (12), the second-order-corrected polaron energy for the state with Landau quantum number n is

$$E_n = E_n^{(0)} + \Delta E_n^{(2)} \ . \tag{13}$$

For most heterojunctions consisting of III-V and II-VI semiconductive compounds,[9] the electron–IO-phonon coupling is weak enough so that the higher-order correction of the polaron energy is negligible.

We now focus our attention on the most remarkable polaron CR due to the electron transition between the lowest two Landau states $n=0$ and 1 caused by a single-phonon process. The polaron CR frequency can be obtained as $\omega_c^* = (E_1 - E_0)/\hbar$. The corresponding polaron cyclotron mass is $m^* = eB_M/\omega_c^* c$. From Eqs. (10), (12) and (13), by defining $\lambda_\xi^2 = \omega_c/\omega_\xi$, the polaron ground-state energy for the $n=0$ Landau level is obtained by a calculation similar to Refs. 5–7:

$$E_0 = \tfrac{1}{2}\hbar\omega_c - \frac{\sqrt{\pi}}{2}\sum_\xi \hbar\omega_\xi\alpha_\xi\lambda_\xi \int_0^\infty dt\, e^{-t}(1-e^{-\lambda_\xi^2 t})^{-1/2} \ . \tag{14}$$

In the high-field limit defined by the conditions $\lambda_\xi^2 \to \infty$ and $\lambda_\xi\alpha_\xi \to 0$, the electron states can be restricted to the $n=0$ Landau state. The polaron energy has the simple form

$$E_0 = \tfrac{1}{2}\hbar\omega_c - \frac{\sqrt{\pi}}{2}\sum_\xi \hbar\omega_\xi\alpha_\xi\lambda_\xi \ . \tag{15}$$

Let us discuss the polaron energy of $n=1$ Landau state for the following several situations.

A. Weak magnetic field

When $\omega_c < \omega_\xi$ is satisfied, for any intermediate Landau state n' the denominator in Eq. (12) is reduced to

$$E_1^{(0)} - E_{n'}^{(0)} - \hbar\omega_\xi = (1-n')\hbar\omega_c - \hbar\omega_\xi < 0 \ .$$

For the $n=1$ Landau state the polaron energy (13) can be written as

$$E_1 = \tfrac{3}{2}\hbar\omega_c - \frac{\sqrt{\pi}}{4}\sum_\xi \hbar\omega_\xi\alpha_\xi \int_0^\infty dt\, e^{-(1-\lambda_\xi^2)t}(1+e^{-\lambda_\xi^2 t}) \times (1-e^{-\lambda_\xi^2 t})^{-1/2} \ , \tag{16}$$

We can easily obtain the polaron cyclotron frequency ω_c^*:

$$\frac{\omega_c^*}{\omega_c} = 1 - \frac{\sqrt{\pi}}{4}\sum_\xi \frac{\alpha_\xi}{\lambda_\xi}\int_0^\infty dt\, e^{-(1-\lambda_\xi^2)t}(1-e^{-\lambda_\xi^2 t})^{1/2} \ . \tag{17}$$

The polaron cyclotron mass m^* can be derived by

$$\frac{m^*}{m} = 1 - \frac{\sqrt{\pi}}{4}\sum_\xi \frac{\alpha_\xi}{\lambda_\xi}\int_0^\infty dt\, e^{-(1-\lambda_\xi^2)t}(1-e^{-\lambda_\xi^2 t})^{-1/2} \ . \tag{18}$$

For the low-magnetic-field limit ($\omega_c \ll \omega_\xi$ or $\lambda_\xi^2 \ll 1$), Eqs. (17) and (18), respectively, are reduced to

$$\frac{\omega_c^*}{\omega_c} = 1 - \frac{\pi}{8}\sum_\xi \alpha_\xi - \frac{9\pi}{64}\sum_\xi \alpha_\xi\lambda_\xi^2 \tag{19}$$

and

$$\frac{m^*}{m} = 1 + \frac{\pi}{8}\sum_\xi \alpha_\xi + \frac{9\pi}{64}\sum_\xi \alpha_\xi\lambda_\xi^2 \ . \tag{20}$$

Comparing (19) and (20) with the results of bulk 2D limit,[7] the different influences between the interface and bulk phonons are easily seen. When $\lambda_\xi^2 = 0$, the polaron mass agrees with what was obtained for a free interface polaron in our previous paper.[9]

B. Strong magnetic field

This is the case of $\omega_c \gg \omega_\xi$, for which one can restrict the intermediate Landau state n' to be 1 in Eq. (12). The energy, cyclotron frequency, and cyclotron mass of the polaron can be derived as the following simplest forms:

$$E_1 = \tfrac{3}{2}\hbar\omega_c - \tfrac{3}{8}\sqrt{\pi}\sum_\xi \hbar\omega_\xi\alpha_\xi\lambda_\xi \ , \tag{21}$$

$$\frac{\omega_c^*}{\omega_c} = 1 + \frac{\sqrt{\pi}}{8}\sum \frac{\alpha_\xi}{\lambda_\xi} \ , \tag{22}$$

and

$$\frac{m^*}{m} = 1 - \frac{\sqrt{\pi}}{8}\sum \frac{\alpha_\xi}{\lambda_\xi} \ . \tag{23}$$

The above two limiting results of Eqs. (19), (20), (22), and (23) indicate that the polaron cyclotron mass is pinned at the free polaron effective mass for zero field, whereas it is pinned at its band mass for high field. This property is same as that of a bulk magnetopolaron.

C. Resonant electron–IO-phonon coupling

Around $\omega_c \cong \omega_\xi$, the two states of the system given by an electron in the zeroth Landau level plus one real IO phonon, and only an electron in the first Landau level, are originally degenerate. However, a weak electron-phonon coupling lifts the degeneracy and thus splits the first Landau level into two levels. Replacing $E_1^{(0)} - E_{n'}^{(0)}$ by $E_1 - E_{n'}$ in Eq. (12), with a procedure similar to that in Ref. 7, one can easily obtain the polaron energy for the first Landau level. Since there are two branches of IO-phonon modes interacting with the electron, we now dis-

TABLE I. Parameters of some III-V compound semiconductors. The electron band mass m is in units of a free electron. α_B is the dimensionless electron bulk LO-phonon coupling constant. LO- and TO-phonon frequencies ω_{LO} and ω_{TO} are given in units of meV.

Materials	ε_0	ε_∞	m	α_B	$\hbar\omega_{LO}$	$\hbar\omega_{TO}$
GaP (Ref. 9)	10.28	8.46	0.336	0.200	49.96	45.32
AlAs (Ref. 19)	10.06	8.16	0.109	0.126	50.09	45.11
InP (Ref. 9)	14.00	10.70	0.07	0.105	42.00	36.72
GaAs (Ref. 19)	12.83	10.90	0.0655	0.0681	36.70	33.83
GaSb (Ref. 7)	15.69	14.44	0.0467	0.0255	29.80	28.59
InAs (Ref. 8)	14.61	11.80	0.0342	0.0640	30.20	27.14

cuss the following two cases.

(i) When the resonant field satisfies the condition $|(\omega_c-\omega_-)/(\omega_c-\omega_+)|\ll 1$ and $\omega_c\cong\omega_-$, the ω_+ branch can be neglected. In this situation, we obtain

$$E_1=\tfrac{3}{2}\hbar\omega_c+\sum_{\mathbf{q}} G_-^2(q)\frac{{}_A\langle 1|L_q^{-1}L_q|1\rangle_A}{E_1-E_0-\hbar\omega_-}$$
$$=\tfrac{3}{2}\hbar\omega_c+\frac{\sqrt{\hbar}}{4}\frac{(\hbar\omega_-)^2\alpha_-\lambda_-}{E_1-E_0-\hbar\omega_-}\ , \tag{24}$$

where E_0 is given by Eq. (14). From the two solutions $E_{1,\pm}$ of Eq. (24), we derive a twofold splitting of the polaron CR frequencies:

$$\omega^*_{c,\pm}=(E_{1,\pm}-E_0)/\hbar\ . \tag{25}$$

The splitting of the polaron cyclotron mass is given by

$$\frac{m^*_\pm}{m}=\frac{\hbar\omega_c}{E_{1,\pm}-E_0}\ . \tag{26}$$

(ii) The other twofold splitting will occur at the condition of $|(\omega_c-\omega_+)/(\omega_c-\omega_-)|\ll 1$ and $\omega_c\cong\omega_+$. In this case Eq. (24) becomes

$$E_1=\tfrac{3}{2}\hbar\omega_0+\frac{\sqrt{\hbar}}{4}\frac{(\hbar\omega_+)^2\alpha_+\lambda_+}{E_1-E_0-\hbar\omega_+}\ , \tag{27}$$

and Eqs. (25) and (26) also hold.

Since there is an obvious difference between the higher and lower frequencies of IO-phonon modes,[9] the trifold splitting cannot occur for general heterojunctions. In special cases, for example in the heterojunction of InAs/GaSb ($\omega_+/\omega_-\cong 1.07$; see the parameters listed in Tables I and II of Sec. III), the trifold splitting might occur around $\omega_c\cong\omega_+\cong\omega_-$. The degenerated polaron energy $E_{1,i}$ ($i=1,2,3$) for the $n=1$ Landau level can be obtained by solving

$$E_1=\tfrac{3}{2}\hbar\omega_c+\frac{\sqrt{\hbar}}{4}\frac{(\hbar\omega_+)^2\alpha_+\lambda_+}{E_1-E_0-\hbar\omega_+}+\frac{\sqrt{\hbar}}{4}\frac{(\hbar\omega_-)^2\alpha_-\lambda_-}{E_1-E_0-\hbar\omega_+}\ . \tag{28}$$

The polaron cyclotron frequency and mass, respectively, are given by

$$\omega^*_{c,i}=(E_{1,i}-E_0)/\hbar \tag{29}$$

and

$$\frac{m_i^*}{m}=\frac{\hbar\omega_c}{E_{1,i}-E_0}\ . \tag{30}$$

However, the electron-phonon coupling due to IO phonons of lower frequency in InAs/GaSb is so small ($\alpha_+/\alpha_-=6.33$) that the observable effect in experiment is weakened.

TABLE II. Calculated parameters for some heterojunctions ($\frac{1}{2}$). α_s is the coupling constant between the electron and IO phonons with frequency ω_s obtained by single IO-phonon mode approximation (Ref. 6). α_+ (α_-) is the coupling constant corresponding to the IO phonon mode, with frequency ω_+ (ω_-) obtained by using our formulas.

Heterojunctions	α_s	$\hbar\omega_s$	α_+	$\hbar\omega_+$	α_-	$\hbar\omega_-$
InP/GaP	0.106	41.33	0.0629	48.08	0.0391	38.97
GaAs/AlAs	0.0862	37.07	0.0496	47.56	0.0308	35.17
InAs/GaSb	0.0408	29.17	0.0350	29.91	0.0055	27.87
GaAs/GaSb	0.0439	35.89	0.0333	35.22	0.0121	29.13

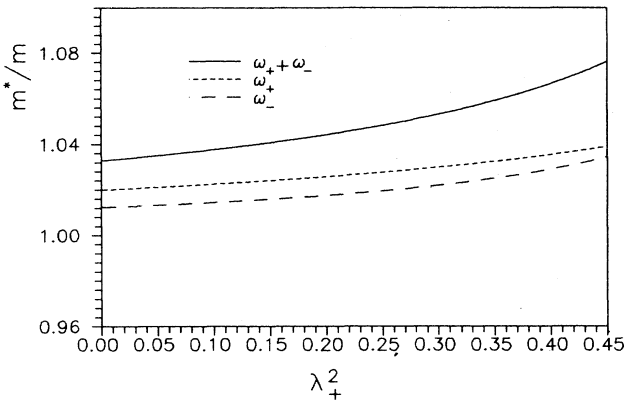

FIG. 1. The polaron cyclotron mass influenced by the two IO-phonon modes ω_+ and ω_- as a function of the magnetic-field intensity (weak-field case).

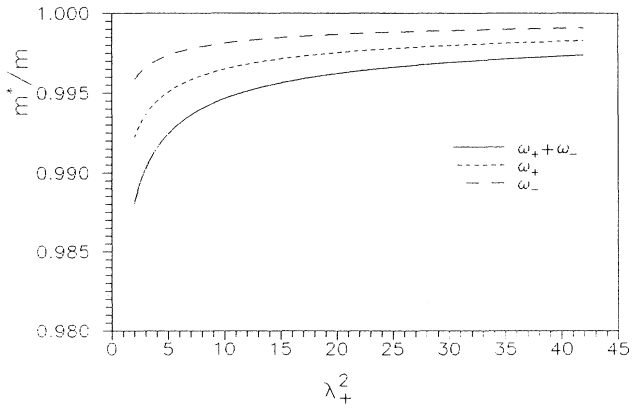

FIG. 2. The polaron cyclotron mass influenced by the two IO-phonon modes ω_+ and ω_- as a function of the magnetic-field intensity (strong-field case).

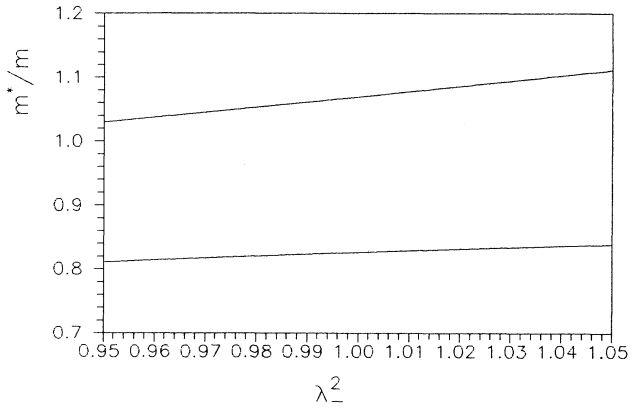

FIG. 3. The polaron cyclotron mass splitting around the cyclotron frequency ω_c approaching the IO-phonon frequency ω_-.

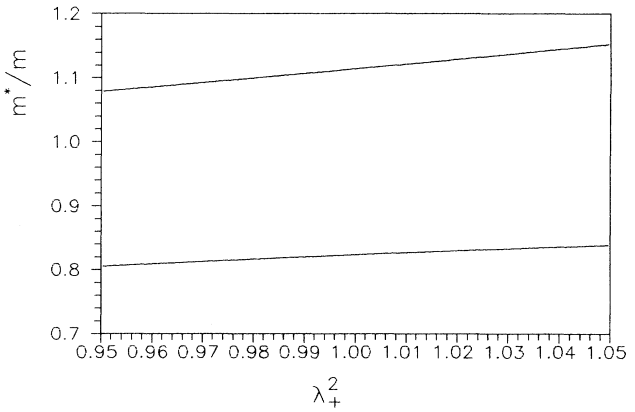

FIG. 4. The polaron cyclotron mass splitting around the cyclotron frequency ω_c approaching the IO-phonon frequency ω_+.

III. DISCUSSION AND CONCLUSION

The parameters for some heterojunctions of III-V semiconductive compounds used in our discussion have been listed in Tables I and II. It is easily seen from Table II that the single branch approximation of IO-phonon modes could be improved by considering the contribution of the two exact IO-phonon modes.

As an example, we have computed numerically the polaron cyclotron mass for the GaAs/AlAs heterojunction to show the twofold splitting properties. For convenience, we choose λ^2 to represent the magnetic-field intensity (here $\lambda_+^2/\lambda_-^2=\omega_-/\omega_+$). In the weak magnetic-field region ($\omega_c<\omega_-<\omega_+$), the polaron cyclotron mass as a function of λ_+^2 is given in Fig. 1. The polaron mass beginning from its normalized effective value increases with increasing λ_+^2. The result in the strong magnetic-field region ($\omega_c>\omega_+$) has been plotted in Fig. 2. At this time, the mass from its value below the electron band mass increases with increasing field. At the high-field limit, its value approaches the electron band mass. The contributions from the two branches of IO-phonon modes, which have been given in Figs. 1 and 2, are comparable. It indicates that not only the IO phonons with higher frequency ω_+ but also those with lower frequency ω_- should be considered.

Figures 3 and 4 denote the mass splittings near the resonant magnetic-field intensities around $\omega_c\cong\omega_-$ ($\lambda_-^2\cong 1$) and $\omega_c\cong\omega_+$ ($\lambda_+^2\cong 1$). For this kind of twofold splitting, one branch is above whereas the other branch is below the electron band mass. Besides the difference of the splitting positions for the magnetic-field intensity, the splitting widths are different since they are proportional to the square root of the electron-phonon coupling strength.

Wei, Kong, and Gu[7] proposed theoretically a trifold splitting of polaron cyclotron mass for GaAs/GaSb by taking into account of the single virtual IO phonons and bulk LO phonons. Our accurate consideration for the IO phonons has given rise to a similar splitting around $\omega_\pm\cong\omega_c\cong\omega_{\rm LO}$.

It is seen from the results in Table II that for some heterojunctions of III-V compounds there is at least one IO-phonon mode near to and above the frequencies of TO modes in material 1. We believe that the IO phonons may play an important role in the explanation of the experimental observations. In fact, if one includes the many-particle effects such as the screening, which slightly lower IO-phonon frequencies, this makes our explanation seem more promising.

In conclusion, we have studied the CR of the interface polarons in semiconductive heterojunctions considering the influence of the entire two IO-phonon modes. There are two splittings around the frequencies of the IO phonons, of which the one caused by phonons with higher frequency is more important. However, the splitting due to the other mode cannot be neglected in, e.g., GaAs/AlAs and InP/GaP structures. In the two IO-phonon modes, at least one is near the TO-phonon frequency. This indicates that further investigation both in theory and experiment is needed to clear up the role of the IO-phonon affect on the interface polarons in semiconductive heterojunctions.

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[1]D. M. Larsen, Phys. Rev. B **30**, 4595 (1984).
[2]L. Esaki, IEEE J. Quantum Electron. **QE-22**, 1611 (1986).
[3]S. Das Sarma, Phys. Rev. Lett. **52**, 859 (1984).
[4]F. M. Peeters, Xiaoguang Wu, and J. T. Devreese, Phys. Rev. B **34**, 1160 (1986).
[5]D. M. Larsen, Phys. Rev. B **33**, 799 (1986).
[6]S. W. Gu, X. J. Kong, and C. W. Wei, Phys. Rev. B **36**, 7977 (1987).
[7]C. W. Wei, X. J. Kong, and S. W. Gu, Phys. Rev. B **38**, 8390

(1988).
[8]Ze Hu, Y. T. Wang, S. W. Gu, T. C. Au-Yeung, and Y. Y. Yeung, Solid State Commun. **81**, 325 (1992); J. Phys. Condensed Matter **4**, 5087 (1992).
[9]S. L. Ban, X. X. Liang, and R. S. Zheng, Phys. Lett. A **192**, 110 (1994).
[10]L. Wendler and R. Haupt, Phys. Status Solidi B **143**, 487 (1987).
[11]N. Mori and T. Ando, Phys. Rev. B **40**, 6175 (1989).
[12]X. X. Liang and S. G. Davison, Surf. Sci. **298**, 225 (1993); X. X. Liang, J. Phys. Condensed Matter **4**, 9769 (1992).
[13]M. H. Degani and G. A. Farias, Phys. Rev. B **41**, 3572 (1990).
[14]G. A. Farias, M. H. Degani, and O. Hipólito, Phys. Rev. B **43**, 4113 (1991).
[15]R. J. Nicholas, L. C. Brunel, S. Huant, K. Karrai, J. C. Portal, M. A. Brummell, M. Razeghi, K. Y. Cheng, and A. Y. Cho, Phys. Rev. Lett. **55**, 883 (1985).
[16]M. Ziesmann, D. Heitmann, and L. L. Chang, Phys. Rev. B **35**, 4541 (1987).
[17]S. Das Sarma, Phys. Rev. Lett. **57**, 651 (1986).
[18]G. Z. Zhao, X. Wang, and X. X. Liang, Acta Sci. Nat. Univ. NeiMenggu **25**, 399 (1994).
[19]G. Q. Hai, F. M. Peeters, and J. T. Devreese, Phys. Rev. B **42**, 11 063 (1991).