Properties of a magnetopolaron at the interface of polar-polar crystals

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The effective Hamiltonian of an interface magnetopolaron is derived and its properties are discussed by using Larsen's perturbational method. The results show that in the case of a repulsive image potential, a "dead" layer of magnetopolarons is formed, the thickness of which decreases when the strength of the magnetic field increases. In the case of an attractive image potential, there is a stable magnetopolaron near the interface. However, its ground-state energy increases with the increase of the magnetic field. The reduction of its binding energy and the increase of the distance between the polaron and the interface are both as expected.

I. INTRODUCTION

With the development of solid device science and technique, the interest in the properties of a surface or an interface polaron in crystals is aroused, especially regarding the influence on its properties from a magnetic field.

The mass of a polaron is usually determined by means of cyclotron resonance experiments.[1] In a polar semiconductor or an ionic crystal, the cyclotron resonance frequency $\omega_c^* = eB/m_p^* c$ is affected by the interaction of the electron with modulus of longitudinal-optical (LO) vibration. The mass of a polaron m_p^* and the band mass of an electron m_b are also influenced by the electron-phonon interactions. Thus the cyclotron resonance frequency is certainly related to magnetic field.

Generally Rayleigh-Schrödinger perturbation theory, (RSPT) Wigner-Brillouin perturbation theory (WBPT) or improved Wigner-Brillouin perturbation theory (IMBPT) is applied to obtain the relation between the mass of a magnetopolaron and the strength of the magnetic field by calculating the effect of the electron-phonon interactions on Landau energy levels.[2]

Devreese and Peeters[3] exactly calculated the effective mass of a magnetopolaron on the basis of the extended[4] Feynman[5] polaron model. Interest is not restricted only to the cyclotron resonance; the properties of a bound magnetopolaron also attract attention. The first studies of the properties of a hydrogenlike impurity were made by Yafet, Keyes, and Adams (YKA).[6] After that, many authors improved on YKA's observations. All these studies not only enlarged the magnetic field, but also obtained very exact results.[7–11] Hollox and co-workers[12] calculated the ground-state and bound energies of a conductive electron in a magnetic field. The conductive electron was attracted by a positive electric impurity in the heteroside of a heterostructure.

There have been many scientists who studied the cyclotron resonance of a two-dimensional (2D) polaron. Usually, one first applies a perturbational theory to second order,[13–16] then a more exact method.[17–20] Indeed, the electron gas in an inverse layer of a metal-oxide-semiconductor field-effect transistor (MOSFET) or a thin quantum well with a high-energy barrier can be considered to be a 2D case. However, for an electron moving near an interface or a surface, we cannot describe its behavior as that in a real 2D case. There exists an image potential, so the electron is not completely limited near the interface or the surface sometimes. Gu *et al.*[21,22] investigated this problem and learned that the properties of an electron near a surface are tremendously influenced by the electron-phonon interactions. Their results show that when it is far away from the surface only the electron–bulk-longitudinal-optical-phonon interaction needs consideration; but the electron–surface-optical- (SO) phonon interaction does not. When the distance between the electron and the surface is much less than the radius of the polaron, the electron–SO-phonon interaction has to be taken into account; when the distance is comparable with the radius of the polaron, both electron–SO-phonon and electron–LO-phonon interactions must be considered.

Following the previous studies,[19] Larsen[23] further studied the 2D magnetopolaron and obtained a pithy expression of ground and lower excited states energies of a polaron in a free magnetic field.

In this paper we use Larsen's method[23] to discuss surface and interface magnetopolarons in polar crystals, considering that an electron interacts not only with the image potential, which is from the polar electron cloud of the ions on the surface or the interface, but also with both bulk LO and SO phonons. Under different circumstances, we analyze the influence of a magnetic field on the thickness of the "dead" layer of an interface magnetopolaron and on the ground-state and binding energies of a stable interface magnetopolaron, respectively.

II. HAMILTONIAN

Now we discuss an interface magnetopolaron in polar-polar crystals. There are polar crystals 1 and 2 in

the $z>0$ and $z<0$ semispace, respectively. The x-y plane is their interface. The static uniform magnetic field $\mathbf{B}$ is along the z direction. An electron moves in crystal 1, i.e., the $z>0$ side, so there is a barrier from crystal 2 to it. Without loss of universality, we suppose that the barrier is infinitely high; therefore, the electron is restricted within crystal 1. In this paper we simultaneously take the interactions from bulk LO and SO phonons into account. In addition, we consider an image potential which is from polarizing the electron cloud of interface ions. For the sake of easiness, we suppose that the band mass of the electron is of the same numerical value in spite of the direction of its moving. Under the isotropic effective-mass approximation, the Hamiltonian of this system can be written as[24,25]

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

$$H_0=\frac{1}{2m_b}(p_x-\beta^2 y/4)^2+\frac{1}{2m_b}(p_y+\beta^2 x/4)^2+\frac{p_z^2}{2m_b}+\frac{e^2(\epsilon_{\infty 1}-\epsilon_{\infty 2})}{4z\epsilon_{\infty 1}(\epsilon_{\infty 1}+\epsilon_{\infty 2})}+\sum_{\mathbf{k}}\hbar\omega_L a_{\mathbf{k}}^{\dagger}a_{\mathbf{k}}+\sum_{\mathbf{q}}\hbar\omega_S b_{\mathbf{q}}^{\dagger}b_{\mathbf{q}} , \tag{1b}$$

$$H_{e\text{-LO}}=\sum_{\mathbf{k}}[V_{\mathbf{k}}^{*}\sin(zk_z)\exp(-i\boldsymbol{\rho}\cdot\mathbf{k}_{\parallel})a_{\mathbf{k}}^{\dagger}+\text{H.c.}] , \tag{1c}$$

$$H_{e\text{-SO}}=\sum_{\mathbf{q}}[C_{\mathbf{q}}^{*}e^{-qz}\exp(-i\boldsymbol{\rho}\cdot\mathbf{q})b_{\mathbf{q}}^{\dagger}+\text{H.c.}] , \tag{1d}$$

where

$$\beta^2=\frac{2eB}{c} , \tag{1e}$$

$$V_{\mathbf{k}}^{*}=\frac{i}{|\mathbf{k}|}\left[\frac{4\pi e^2\hbar\omega_L}{\epsilon V}\right]^{1/2} , \tag{1f}$$

$$C_{\mathbf{q}}^{*}=i\left[\frac{e^2\pi\hbar\omega_S}{\epsilon^{*}A|\mathbf{q}|}\right]^{1/2} , \tag{1g}$$

$$\frac{1}{\epsilon}=\frac{1}{\epsilon_{\infty 1}}-\frac{1}{\epsilon_{01}} , \tag{1h}$$

$$\frac{1}{\epsilon^{*}}=\frac{2}{\epsilon_{\infty 1}+\epsilon_{\infty 2}}-\frac{2}{\epsilon_{01}+\epsilon_{02}} , \tag{1i}$$

where x, y, and z are the position coordinates of the electron, $\boldsymbol{\rho}=(x,y,0)$ is the projection of the electron position vector on the x-y plane. $\mathbf{p}=(p_x,p_y,p_z)$ is the momentum of the electron. $\mathbf{k}=(k_x,k_y,k_z)$ is the wave vector of the bulk LO phonon, $\mathbf{k}_{\parallel}=(k_x,k_y,0)$ is the projection of $\mathbf{k}$ on the x-y plane. $\mathbf{q}$ is the 2D wave vector of SO phonons. $a_{\mathbf{k}}^{\dagger}$ ($a_{\mathbf{k}}$) and $b_{\mathbf{q}}^{\dagger}$ ($b_{\mathbf{q}}$) are the creation (annihilation) operators of bulk LO and SO phonons, respectively. ω_L is the frequency of bulk LO phonons, ω_S is the frequency of SO phonons.

The first two terms in (1b) are the kinetic energy of 2D motion of the electron in the x-y plane, which form Landau energy levels. The third is the kinetic energy of the electron in the z direction. The fourth is the image potential energy. The fifth and sixth are the energies of bulk LO and SO phonons, respectively. Equations (1c) and (1d) express the energy of electron–bulk-LO- and –bulk-SO-phonon interactions, respectively. V and A in (1f) and (1g) are the volume of crystal 1 and the interface area, respectively. $\epsilon_{\infty 1}$ ($\epsilon_{\infty 2}$) and ϵ_{01} (ϵ_{02}) are the optical and static dielectric constants of crystal 1 (crystal 2), respectively.

The relation between the frequency of bulk transverse-optical phonons ω_T and the one of bulk LO phonons ω_L is determined by the equation

$$\frac{\omega_L^2}{\omega_T^2}=\frac{\epsilon_0}{\epsilon_{\infty}} ,$$

and the frequencies of bulk LO and SO phonons satisfy the equality

$$\omega_S^2=\omega_L^2\left[\frac{\epsilon_{\infty 1}(\epsilon_{01}+\epsilon_{02})}{\epsilon_{01}(\epsilon_{\infty 1}+\epsilon_{\infty 2})}\right] .$$

The coupling constant of the electron with bulk LO phonons[25] is

$$\alpha_L=\frac{e^2}{2\hbar^2}\left[\frac{2m_b}{\hbar\omega_L}\right]^{1/2}\frac{1}{\epsilon} ,$$

and that of the electron with SO phonons[25] is

$$\alpha_S=\frac{e^2}{2\hbar^2}\left[\frac{2m_b}{\hbar\omega_S}\right]^{1/2}\frac{1}{\epsilon^{*}} .$$

The cyclotron resonance frequency of a free electron in a magnetic field is

$$\omega_c=\frac{\beta^2}{2m_b}=\frac{eB}{m_b c} .$$

We define

$$\lambda_L^2=\frac{\omega_c}{\omega_L},\quad \lambda_S^2=\frac{\omega_c}{\omega_S} .$$

After the introduction of the well-known 1D harmonic-oscillator operators

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

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

which satisfy the following commutation relations:

$$[A,A^{\dagger}]=[B,B^{\dagger}]=1,\quad [A,B]=[A,B^{\dagger}]=0 .$$

We can rewrite (1b), (1c), and (1d) as

$$H_0=\frac{\hbar\beta^2}{2m_b}(A^{\dagger}A+\tfrac{1}{2})+p_z^2/2m_b+e^2(\epsilon_{\infty 1}-\epsilon_{\infty 2})/[4z\epsilon_{\infty 1}(\epsilon_{\infty 1}+\epsilon_{\infty 2})]+\sum_{\mathbf{k}}\hbar\omega_L a_{\mathbf{k}}^{\dagger}a_{\mathbf{k}}+\sum_{\mathbf{q}}\hbar\omega_S b_{\mathbf{q}}^{\dagger}b_{\mathbf{q}} , \tag{3a}$$

$$H_{e\text{-LO}}=\sum_{\mathbf{k}}[V_{\mathbf{k}}^{*}\sin(zk_z)L_{\mathbf{k}}M_{\mathbf{k}}a_{\mathbf{k}}^{\dagger}+\text{H.c.}] , \tag{3b}$$

$$H_{e\text{-SO}}=\sum_{\mathbf{q}}(C_{\mathbf{q}}^{*}e^{-qz}L_{\mathbf{q}}M_{\mathbf{q}}b_{\mathbf{q}}^{\dagger}+\text{H.c.}) , \tag{3c}$$

where

$$L_{\mathbf{k}}=\exp\left[\frac{\sqrt{\hbar}}{\beta}(k_x+ik_y)A-\frac{\sqrt{\hbar}}{\beta}(k_x-ik_y)A^{\dagger}\right] ,$$

$$M_{\mathbf{k}}=\exp\left[\frac{\sqrt{\hbar}}{\beta}(k_x-ik_y)B-\frac{\sqrt{\hbar}}{\beta}(k_x+ik_y)B^{\dagger}\right] ,$$

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

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

III. EFFECTIVE HAMILTONIAN

The adiabatic approximation[21,22,24] is applied to treat the motion of the electron. We first seek the energy of the part of Hamiltonian in the x-y plane that depends on the parameter z. Then, add it to the part of Hamiltonian in the z direction to obtain the effective Hamiltonian.

The system Hamiltonian can be divided into two parts:

$$H=H_{\parallel}+H_{\perp} ,$$

where

$$H_{\perp}=\frac{p_z^2}{2m_b}+\frac{e^2(\epsilon_{\infty 1}-\epsilon_{\infty 2})}{4z\epsilon_{\infty 1}(\epsilon_{\infty 1}+\epsilon_{\infty 2})} ,$$

$$H_{\parallel}=\frac{\hbar\beta^2}{2m_b}(A^{\dagger}A+\tfrac{1}{2})+H_{e\text{-ph}} ,$$

and

$$H_{e\text{-ph}}=H_{e\text{-LO}}+H_{e\text{-SO}} .$$

In the first step, we treat $H_{\parallel}$. When

$$\alpha_L\lambda_L\rightarrow 0, \quad \alpha_S\lambda_S\rightarrow 0 ,$$

the perturbation calculation is a simple method. We shall take $(\hbar\beta^2/2m_b)(A^{\dagger}A+\frac{1}{2})$ as the unperturbed Hamiltonian and $H_{e\text{-ph}}$ as the perturbational one.

At zero temperature, the ground state of the phonons can be thought as a vacuum state $|0\rangle=|0\rangle_{\mathbf{k}}|0\rangle_{\mathbf{q}}$. Having introduced the operators A and B, we can represent the Landau levels as products of two independent 1D harmonic oscillator states, which are

$$(n!)^{-1/2}(A^{\dagger})^n|0\rangle_A \quad \text{and} \quad (M!)^{-1/2}(B^{\dagger})^M|0\rangle_B ,$$

n is the Landau quantum number, M is the z angular momentum quantum number.

Under 0 K condition, the unperturbed state can be restricted to the $n=0$ Landau state. Thus, the unperturbed ground-state wave function is

$$|0\rangle_A|0\rangle_B|0\rangle .$$

However, it is not possible to restrict the intermediate states to the $n=0$ Landau levels. Therefore, we have the intermediate wave functions as follows:

$$(n!M!)^{-1/2}(A^{\dagger})^n|0\rangle_A(B^{\dagger})^M|0\rangle_B a_{\mathbf{k}_1}^{\dagger}a_{\mathbf{k}_2}^{\dagger}\cdots a_{\mathbf{k}_N}^{\dagger}|0\rangle_{\mathbf{k}}b_{\mathbf{q}_1}^{\dagger}b_{\mathbf{q}_2}^{\dagger}\cdots b_{\mathbf{q}_N}^{\dagger}|0\rangle_{\mathbf{q}} .$$

Then we obtain the perturbated ground-state expected value to the second order,

$$\Delta E^{(2)}=\sum_{\mathbf{k}}V_{\mathbf{k}}^2\sin^2(zk_z)|\langle\mathbf{k}|a_{\mathbf{k}}^{\dagger}|0\rangle|^2{}_B\langle M|M_{\mathbf{k}}^{-1}M_{\mathbf{k}}|M\rangle_B\sum_n{}_A\langle 0|L_{\mathbf{k}}^{-1}|n\rangle_A\frac{1}{E_p-n\hbar\omega_c-\hbar\omega_L}{}_A\langle n|L_{\mathbf{k}}|0\rangle_A$$

$$+\sum_{\mathbf{q}}C_{\mathbf{q}}^2e^{-2qz}|\langle\mathbf{q}|b_{\mathbf{q}}^{\dagger}|0\rangle|^2{}_B\langle M|M_{\mathbf{q}}^{-1}M_{\mathbf{q}}|M\rangle_B\sum_n{}_A\langle 0|L_{\mathbf{q}}^{-1}|n\rangle_A\frac{1}{E_p-n\hbar\omega_c-\hbar\omega_S}{}_A\langle n|L_{\mathbf{q}}|0\rangle_A .$$

From WBPT, we know that $E_p=0$, when we calculate the energy only to the second order. After a tedious calculation, one can obtain

$$\Delta E^{(2)}=V_{e\text{-LO}}+V_{e\text{-SO}} , \tag{4}$$

where

$$V_{e\text{-LO}}=-(\sqrt{\pi}/2)(\hbar\omega_L)^2\alpha_L\lambda_L\left[\int_0^{\infty}dt\,e^{-\hbar\omega_L t}/(1-e^{-\hbar\omega_c t})^{1/2}-\int_0^{\infty}dt\,e^{-\hbar\omega_L t}\exp[(\beta^2z^2/\hbar)/(1-e^{-\hbar\omega_c t})]\right.$$

$$\left.\times\{1-\phi(\beta z/[\hbar(1-e^{-\hbar\omega_c t})]^{1/2})\}/(1-e^{-\hbar\omega_c t})^{1/2}\right] .$$

$$V_{e\text{-SO}}=-(\sqrt{\pi}/2)(\hbar\omega_S)^2\alpha_S\lambda_S\int_0^{\infty}dt\,e^{-\hbar\omega_S t}/(1-e^{-\hbar\omega_c t})^{1/2}\exp[(\beta^2z^2/\hbar)/(1-e^{-\hbar\omega_c t})]\{1-\phi(\beta z/[\hbar(1-e^{-\hbar\omega_c t})]^{1/2})\} .$$

The unperturbed ground-state expected value

$$\langle 0|_A\langle 0|_B\langle M|\frac{\hbar\beta^2}{2m_b}(A^\dagger A+\tfrac{1}{2})|M\rangle_B|0\rangle_A|0\rangle = \tfrac{1}{2}\hbar\omega_c \ . \quad (5)$$

Finally, we have the effective Hamiltonian, i.e., the summation of (4), (5), and $H_\perp$:

$$H_{\rm eff}=\tfrac{1}{2}\hbar\omega_c+\frac{p_z^2}{2m_b}+\frac{e^2(\epsilon_{\infty 1}-\epsilon_{\infty 2})}{4z\epsilon_{\infty 1}(\epsilon_{\infty 1}+\epsilon_{\infty 2})}+V_{e\text{-LO}}+V_{e\text{-SO}} \ . \quad (6)$$

In expression (6), the first term is the ground-state energy of Landau levels, the second is the kinetic energy of the electron in the z direction, the third is the image potential energy of the electron, which is represented with $V_{\rm im}$, the last two terms are the energies from the electron–bulk-LO-phonon ($V_{e\text{-LO}}$) and SO-phonon interactions ($V_{e\text{-SO}}$).

The effective potential is

$$V_{\rm eff}=V_{\rm im}+V_{e\text{-LO}}+V_{e\text{-SO}} \ . \quad (7)$$

Now we discuss the properties of an interface magnetopolaron in the following three cases.

A. The outside is vacuum i.e., $\epsilon_{02}=\epsilon_{\infty 2}=1$

When the $z>0$ semispace is filled with polar crystal 1 and $z<0$ is vacuum, the image potential is repulsive. For example, crystal 1 is GaAs, then the image potential can be written as

$$V_{\rm im}=\frac{e^2(10.9-1)}{4z\times 10.9(10.9+1)}>0 \ .$$

Figure 1 shows $V_{\rm im}$, $V_{e\text{-LO}}$, $V_{e\text{-SO}}$, and $V_{\rm eff}$ for the different z. From it we see that $V_{\rm eff}>0$, when $z<69$ Å. In this region the electron is repelled, so it hardly appears in it. This surface layer is called a polar free-surface layer (PFSL) or "dead" layer.

In order to discuss the influence of the magnetic field on a surface polaron, the potentials are evaluated when $B=10$, 10^2, 10^3, 10^4, 10^5 G. The curves of $V_{\rm im}$, $V_{e\text{-LO}}$, $V_{e\text{-SO}}$, and $V_{\rm eff}$ for $B=10^2$ and 10^5 G are shown in Fig. 2. They show that when B increases, the absolute values of $V_{e\text{-LO}}$ and $V_{e\text{-SO}}$ increase. It means that the electron–bulk-LO- and SO-phonon interactions are both strengthened with increasing magnetic field. Because $V_{\rm im}$ is independent of B, $V_{\rm eff}$ is weakened with increasing B. Thus, the thickness of the "dead" layer D_0 will decrease in this case. For example, $D_0\simeq 67$ Å when $B=10^2$ G but $D_0\simeq 61$ Å when equal to 10^5 G.

B. In the case of $\epsilon_{\infty 2}<\epsilon_{\infty 1}$ and $\epsilon_{\infty 2}\simeq\epsilon_{\infty 1}$

The image potential is weakly repulsive. For instance, take GaAs as crystal 1 and $Ga_{1-0.3}Al_{0.3}As$ as crystal 2, then

$$V_{\rm im}=\frac{e^2(10.9-10.7)}{4z\times 10.9(10.9+10.7)}>0 \ .$$

In Fig. 3 the curves of $V_{\rm im}$, $V_{e\text{-LO}}$, $V_{e\text{-SO}}$, and $V_{\rm eff}$ are drawn for $B=10^2$ and $B=10^5$ G. In addition to the conclusions given above, due to $\epsilon_{\infty 2}\simeq\epsilon_{\infty 1}$ the "dead" layer of the interface magnetopolaron is very thin.

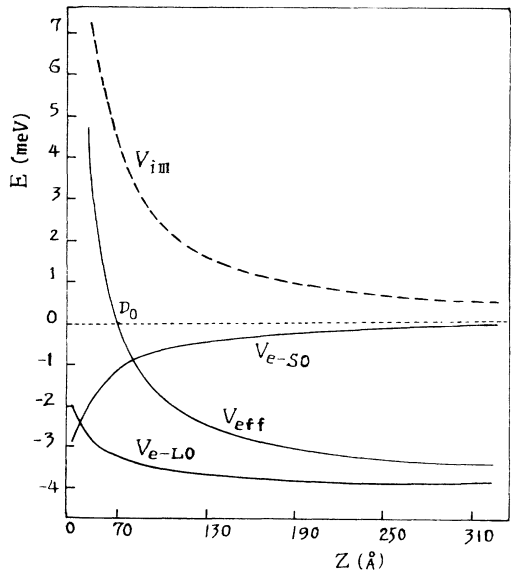

FIG. 1. $V_{\rm im}$, $V_{e\text{-LO}}$, $V_{e\text{-SO}}$, and $V_{\rm eff}$ vs z for the magnetic field $B=0$. D_0 is the thickness of its dead layer. GaAs occupies the $z>0$ semispace and $z<0$ is vacuum.

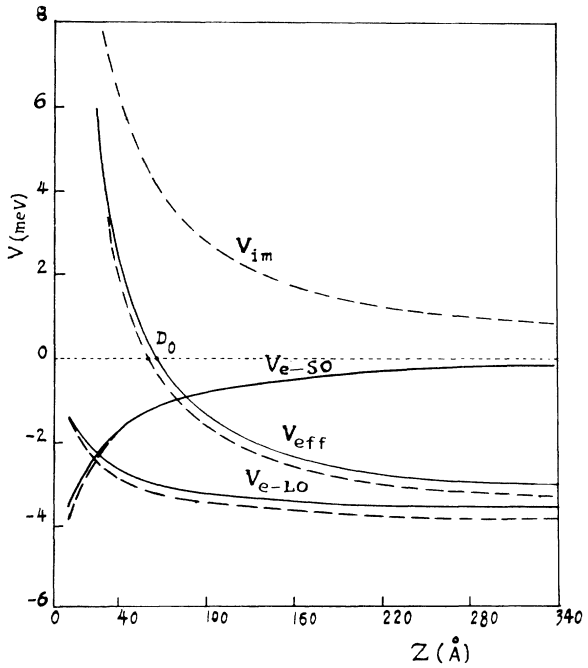

FIG. 2. $V_{\rm im}$, $V_{e\text{-LO}}$, $V_{e\text{-SO}}$, and $V_{\rm eff}$ vs z for $B=10^2$ and 10^5 G. $V_{\rm im}$ is independent of B. The solid lines are for $B=10^2$ G. The dotted lines are for $B=10^5$ G.

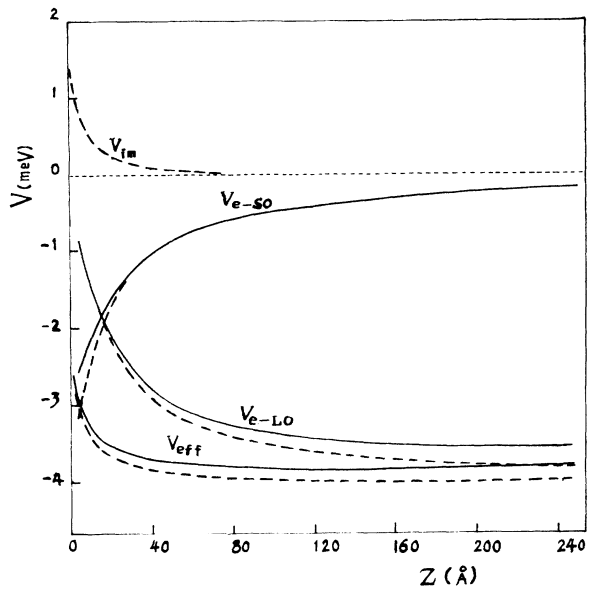

FIG. 3. Same as Fig. 2, but the $Ga_{1-0.3}Al_{0.3}As$ occupies $z<0$ semispace. Here, however, is the case of the interface magnetopolaron.

C. $\epsilon_{\infty 2}>\epsilon_{\infty 1}$, i.e., the image potential is attractive

For instance, the GaAs is taken as crystal 1 and GaSb as crystal 2. Then

$$V_{\rm im}=\frac{e^2(10.9-14.44)}{4z\times 10.9(10.9+14.44)}<0\ .$$

We make the evaluations for $B=10$, 10^2, 10^3, 10^4, and 10^5 G. The curves of $V_{\rm im}$, $V_{e\text{-LO}}$, $V_{e\text{-SO}}$, and $V_{\rm eff}$ are drawn in Fig. 4. They show that the absolute values of $V_{e\text{-LO}}$ and $V_{e\text{-SO}}$ increase with increasing B. That implies that the electron–bulk-LO- and SO-phonon interactions are both strengthened with increasing B. $V_{\rm im}$, $V_{e\text{-LO}}$, and $V_{e\text{-SO}}$ are all attractive. Therefore, the magnetopolaron is found to be bound near the interface, so that a stable interface magnetopolaron is formed. In the next section we shall discuss it in detail.

IV. INTERFACE MAGNETOPOLARON

Expression (6) shows that the effective Hamiltonian of an interface magnetopolaron is a function of magnetic field B and the coordinate z of the electron. Using variational method, we evaluate the ground-state and binding energies. From $V_{\rm eff}$ we know that the available trial wave function is

$$\phi(z)=2\xi^{3/2}ze^{-\xi z}\ , \tag{8}$$

where ξ is the variational parameter.

The expected value of $H_{\rm eff}$ for the z-direction wave function is

$$\begin{aligned}\overline{H_{\rm eff}}&=\langle\phi(z)\,|\,H_{\rm eff}\,|\,\phi(z)\rangle\\
&=\tfrac{1}{2}\hbar\omega_c+\frac{\hbar^2\xi^2}{2m_b}+\frac{\xi e^2(\epsilon_{\infty 1}-\epsilon_{\infty 2})}{4\epsilon_{\infty 1}(\epsilon_{\infty 1}+\epsilon_{\infty 2})}\\
&\quad-(\hbar\omega_L)^2\alpha_L\lambda_L\left[(\sqrt{\pi}/2)\int_0^\infty dt\,e^{-\hbar\omega_L t}/(1-e^{-\hbar\omega_c t})^{1/2}\right.\\
&\qquad\left.-\xi\frac{\sqrt{\hbar}}{\beta}\int_0^\infty dt\,e^{-\hbar\omega_L t}\int_0^\infty\frac{du}{(1+u)^3}\exp\left[-\frac{\hbar\xi^2u^2}{\beta^2}(1-e^{-\hbar\omega_c t})\right]\right]\\
&\quad-(\hbar\omega_S)^2\alpha_S\lambda_S\int_0^\infty dt\,e^{-\hbar\omega_S t}\frac{\sqrt{\hbar}\xi}{\beta}\int_0^\infty\frac{du}{(1+u)^3}\exp\left[\frac{\hbar\xi^2u^2}{\beta^2}(1-e^{-\hbar\omega_c t})\right]\ .\end{aligned}\tag{9}$$

ξ in the expression (9) is determined by the equation

$$\begin{aligned}\frac{\delta\overline{H_{\rm eff}}}{\delta\xi}&=\frac{\hbar^2\xi}{m_b}+\frac{e^2(\epsilon_{\infty 1}-\epsilon_{\infty 2})}{4\epsilon_{\infty 1}(\epsilon_{\infty 1}+\epsilon_{\infty 2})}+(\hbar\omega_L)^2\alpha_L\lambda_L\int_0^\infty dt\,e^{-\hbar\omega_L t}\left[\int_0^\infty\frac{du}{(1+u)^3}\exp\left[-(1-e^{-\hbar\omega_c t})\frac{\hbar\xi^2u^2}{\beta^2}\right]\right.\\
&\qquad\left.\times\left[\frac{\sqrt{\hbar}}{\beta}-\frac{2\sqrt{\hbar}\hbar\xi^2}{\beta^3}u^2(1-e^{-\hbar\omega_c t})\right]\right]\\
&\quad-(\hbar\omega_S)^2\alpha_S\lambda_S\int_0^\infty dt\,e^{-\hbar\omega_S t}\left[\int_0^\infty\frac{du}{(1+u)^3}\exp\left[-(1-e^{-\hbar\omega_c t})\frac{\hbar\xi^2u^2}{\beta^2}\right]\right.\\
&\qquad\left.\times\left[\frac{\sqrt{\hbar}}{\beta}-\frac{2\sqrt{\hbar}\hbar\xi^2}{\beta^3}u^2(1-e^{-\hbar\omega_c t})\right]\right]=0\ .\end{aligned}\tag{10}$$

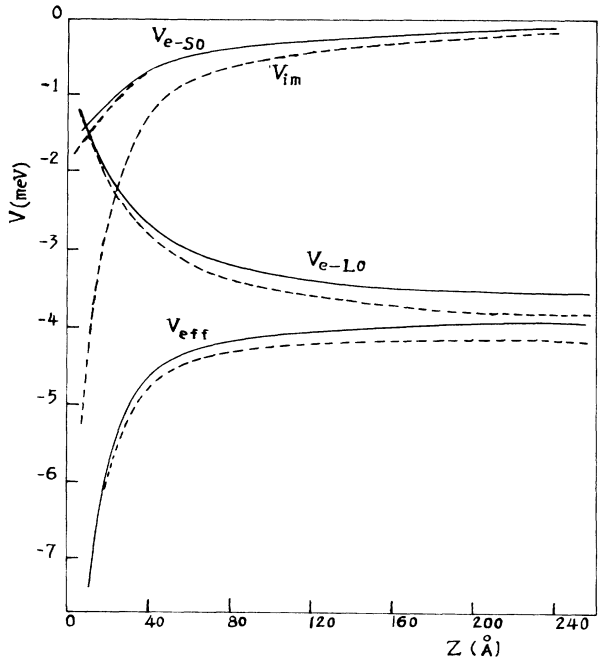

FIG. 4. Same as Fig. 2, but the GaAs occupies $z<0$ semispace.

From Eq. (10) we can obtain the solution $\xi=\xi(B)$ which represents the reverse radius of the interface polaron. Putting the numerical value of ξ in expression (9), we can come to the ground-state energy E_G. The binding energy E_G^B is defined as follows:

$$E_G^B=V_{\text{eff}}(z\to\infty)-E_G \ . \tag{11}$$

As an example, GaAs is taken as crystal 1 and GaSb as crystal 2.

The results (when $B=10^2$ G, $\xi=27\,234.1$ cm^{-1} and when $B=10^5$ G, $\xi=27\,233.6$ cm^{-1} show that the value of ξ decreases slightly when the magnetic field increases. The values of ξ, E_G, and E_G^B are listed in Table I.

TABLE I. The ground-state energy, binding energy, and parameter ξ of the interface magnetopolaron for the different magnetic fields B.

B (G)	ξ (cm^{-1})	E_G (meV)	E_G^B (meV)
10^2	27 234.11	−3.648	$3.442\,881\times10^{-3}$
10^3	27 234.10	−3.571	$3.442\,878\times10^{-3}$
10^4	27 234.00	−2.799	$3.442\,847\times10^{-3}$
10^5	27 233.56	4.839	$3.442\,633\times10^{-3}$

V. CONCLUSION AND DISCUSSION

Using Larsen perturbational method, we obtained the effective Hamiltonian of a magnetopolaron at the interface of polar-polar crystals, which is available in the weak-coupling and free-magnetic-field cases. We studied the relation between the behavior of an interface magnetopolaron and magnetic field.

When the dielectric constant of medium 2 is smaller than that of polar crystal 1, i.e., the electron at the interface or surface comes to be repelled by the interface or surface image potential, it is difficult for a polaron to appear near the interface or the surface, then the PFSL (or called "dead" layer) is formed near the interface or the surface. The thickness of the "dead" layer is dependent of magnetic field of strength, and the stronger the magnetic field is, the thinner the thickness. If the dielectric constant of medium 2 is larger than that of polar crystal 1, i.e., the electron comes to be attracted by the interface image potential, the polaron will form a stable state near the interface. The ground-state energies of the interface magnetopolaron increase with the increase of the magnetic field, but the reduction of its binding energies and the increase of the distances between the magnetopolaron and the interface are both not notable.

Larsen perturbational method applied in this paper is available in the weak-coupling cases, but the magnetic field can be free. The perturbational terms were only expanded to the second order. Because GaAs is taken as an example for numerical computation, whose coupling constants ($\alpha_L=0.067$, $\alpha_S=0.044$) are very small, it is needless to calculate the higher-order terms.

[1]See, e.g., *Polarons and Excitons*, edited by C. G. Kuper and G. D. Whitfield (Oliver and Boyd, Edinburgh, 1963); *Polarons in Ionic Crystals and Polar Semiconductors*, edited by J. T. Devreese (North-Holland, Amsterdam, 1972); *Polarons and Excitons in Polar Semiconductors and Ionic Crystals*, edited by J. T. Devreese and F. M. Peeters (Plenum, New York, 1984).
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