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Acoustic surface plasmons in type-II semiconducting superlattices

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The spectrum of electronic collective excitations of a semi-infinite type-II semiconducting superlattice contains two novel surface polariton branches, one of which is an acoustic surface plasmon. As a result of the quantized nature of the electronic motion within the layers, this new type of surface polariton does not suffer Landau damping. The acoustic branch is of particular interest because the group velocity of this mode can be controlled by sample preparation, making these systems potentially useful for surface-wave devices.

The spectrum of the collective excitations of infinite type-I and type-II semiconducting superlattices[1] has recently received some attention both from a theoretical[2,3] as well as from an experimental[4,5] viewpoint. More recently, the surface plasmons of a semi-infinite type-I superlattice have also been studied.[6–8]

In this paper we investigate the dispersion relation of surface plasmons for a model semi-infinite type-II semiconducting superlattice consisting of a periodic array of alternate quasi-two-dimensional electron and hole layers embedded in a background characterized by a dielectric constant ϵ_s, and terminated at the interface with an insulating medium of dielectric constant ϵ_0. In the electric quantum limit and at low temperatures this model gives a reasonable description of the intrasubband excitations[9] of a typical type-II superlattice like the GaSb/InAs system.

Within the random-phase approximation the dispersion relation for the bulk plasmons of a type-II superlattice in the electrostatic limit is readily found to be[2,3]

$$\omega_{\pm}(q,k)=\left\{\frac{\omega_P^{(e)2}+\omega_P^{(h)2}}{2\epsilon_s}S(qa,ka)\pm\left[\left(\frac{\omega_P^{(e)2}+\omega_P^{(h)2}}{2\epsilon_s}S(qa,ka)\right)^2+\frac{\omega_P^{(e)2}\omega_P^{(h)2}}{\epsilon_s}\tilde{S}(qa,ka)\right]^{1/2}\right\}^{1/2}, \quad (1)$$

where we have defined

$$S(x,y)=\frac{x\sinh 2x}{\cosh 2x-\cos 2y} \quad (2)$$

and

$$\tilde{S}(x,y)=\frac{x^2(1-\cosh 2x)}{\cosh 2x-\cos 2y} \quad . \quad (3)$$

Here q and k are the wave-vector components parallel and perpendicular to the layers, respectively. The electric field associated with the excitations is taken to be proportional to $\exp(iqy+ikz-i\omega t)$, where z is normal to the layers. In (1),

$$\omega_P^{(e),(h)}=(4\pi n_{e,h}e^2/2am_{e,h}^*)^{1/2}$$

is the equivalent three-dimensional plasma frequency for the electron (hole) system alone. $n_{e,h}$ and $m_{e,h}^*$ are, respectively, the number density per unit area and the effective mass of the electron- (hole-) type layers. As k varies between 0 and $\pi/2a$, these excitations form two bands of acoustic modes which remain separated provided that the ratio $\omega_R=\omega_P^{(h)}/\omega_P^{(e)}$ is different from unity.[10] This situation is illustrated in Fig. 1, a plot of ω vs q, in which the bulk plasmon bands are represented by the two upper shaded regions. The single-particle continuum appears as the lower shaded region in Fig. 1.

Following the procedure discussed in Ref. 6 leads in a straightforward way to the surface-mode condition for a semi-infinite type-II superlattice. In the case of an electron-type first layer, in the nonrelativistic limit we ob-

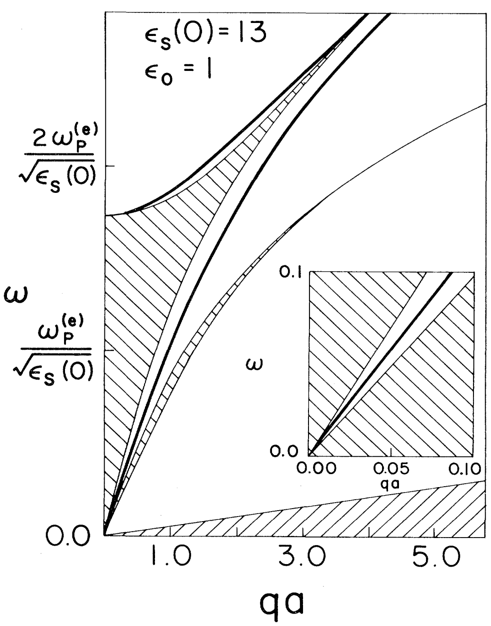

FIG. 1. Dispersion relation ω vs q for the surface plasmons of a type-II superlattice. The spectrum consists of two branches (thick lines). The inset contains an expanded plot of the small qa region for the acoustic branch. The two upper shaded (left to right) regions constitute the bulk-plasmon-band continuum. The lower shaded (right to left) region is the single-particle continuum. The values of the parameters are as follows: $\epsilon_s=13.0$, $\epsilon_0=1.0$, $\omega_R^2=2.0$, $\omega_P^{(e)}=1.1\times10^{14}$ Hz, $v_F^e=3.3\times10^7$ cm/sec, $a=500$ Å.

tain[11]

$$[\epsilon_s+\epsilon_0-2v(q)\chi(q,\omega)][\epsilon_s+2v(q)\chi(q,\omega)+(e^{2\alpha a+qa}-\cosh qa)(\sinh qa)^{-1}]$$
$$-[\epsilon_s-\epsilon_0+2v(q)\chi(q,\omega)][\epsilon_s-2v(q)\chi(q,\omega)-(e^{2\alpha a-qa}-\cosh qa)(\sinh qa)^{-1}]=0\ , \quad (4)$$

where $v(q)=2\pi e^2/q$, and $\chi(q,\omega)$ is the density response function of a two-dimensional electron gas.[12] In (4), q has the same meaning as before, whereas α is a complex quantity whose real part is positive and represents the inverse penetration depth of the surface mode. More precisely, we have assumed that if E_n is the electric field at the nth electronic layer of the array, the following relation holds: $E_n=\exp(-2\alpha a)E_{n-1}$. The value of α to be used in (4) is determined by the requirement that the bulk-mode condition be satisfied for the same values of ω and q. This leads to the following relationship:

$$\alpha=\frac{1}{2a}\text{arccosh}\left[\cosh 2qa-\frac{\omega_P^{(e)2}+\omega_P^{(h)2}}{\epsilon_s\omega^2}qa\sinh 2qa-\left(\frac{\omega_P^{(e)}\omega_P^{(h)}}{\epsilon_s\omega^2}\right)^2(qa)^2(1-\cosh 2qa)\right]\ . \quad (5)$$

A direct inspection of Eq. (5) makes it apparent that α is real everywhere in the plane ω,q except inside the bulk plasmon bands where it is purely imaginary and coincides with k, and in the gap region between these two bands where it acquires an imaginary part equal to $\pi/2a$. Finally, the dispersion relation for the surface plasmons can be obtained with use of (5) in (4).

We have carried out an extensive numerical analysis of the dispersion relation (4) for the longitudinal surface plasmons in a type-II superlattice, for both the cases of a rigid and a polarizable dielectric background. This has been done introducing a plasmon-optical-phonon coupling through the use of a frequency-dependent dielectric constant $\epsilon_s(\omega)$. While a more complete discussion will be deferred to a forthcoming publication,[8] we report here the main results and conclusions of our study with particular emphasis on the case of a rigid dielectric background.

Type-II superlattices admit, in general, two branches of surface plasmons whose features and location in the ω,q plane crucially depend upon the values of the dielectric constants ϵ_s and ϵ_0, and the parameter ω_R. Similarly to the corresponding case in type-I superlattices, surface plasmons can exist above or below the upper and lower edges of the bulk-plasmon bands (BPB's). The similarity also extends to the absence of intrinsic Landau damping which is a peculiar feature of the surface modes of all types of semiconducting superlattices. This remarkable property, which is due to the quantization of the motion along the z direction, makes the surface modes of these systems long-lived excitations.

A novel feature of the present case is the existence of a surface-plasmon branch lying within the gap between the two BPB's. It turns out that a branch of acoustic surface plasmons always exists between the two band edges

$$\omega\pm(q,\pi/2a)=\omega_P^{(e),(h)}(qa\tanh qa/\epsilon_s)^{1/2}$$

for all values of q. In fact, for small values of qa in this region, Eq. (4) can be solved analytically, leading to the following acoustic plasmon dispersion relation:

$$\omega_{SP}(q)=[2/\epsilon_s(1+\omega_R^2)]^{1/2}\omega_P^{(h)}qa+O(q^2a^2)\ . \quad (6)$$

This behavior is illustrated in the inset of Fig. 1. This acoustic surface mode never merges into the bulk-plasmon continuum.

The location of the second surface-plasmon branch in the ω,q plane strongly depends upon the value of the ratio $\epsilon_R=\epsilon_s/\epsilon_0$, and the situation must be analyzed case by case. We first notice that this branch, which directly corresponds to the one appearing in type-I superlattices,[6,7] can only exist either above the upper edge of the upper BPB continuum or under the lower edge of the lower one. In both cases α is purely real. Furthermore, except for extremely limiting situations, these modes never exist for small values of qa. At some finite value of q they intersect the BPB continuum. The critical wave vector q^* where this occurs also corresponds to the value of α going to zero as the surface mode turns into a bulk mode. q^* is readily evaluated from the mode condition, Eq. (4), and is found to be

$$q^*=\frac{1}{2a}\ln\left[\frac{(\epsilon_R+1)(\epsilon_R+\omega_R^2)}{(\epsilon_R-1)(\epsilon_R-\omega_R^2)}\right]\ . \quad (7)$$

Notice that q^* is infinite for ϵ_R between 1 and ω_R^2. This implies that for these values of ϵ_R only the acoustic branch is present.

For $\epsilon_R>\omega_R^2$, as illustrated in Fig. 1, the second surface-plasmon branch lies above the BPB continuum and merges into it from above at the value of q^* as given by Eq. (7). Making use of Eq. (5) it is possible to calculate for each branch the wave-vector-dependent inverse penetration depth. For the case corresponding to the parameters used in Fig. 1, this quantity is shown in Fig. 2. Finally, for $\epsilon_R<1$, the second branch lies below the BPB continuum,

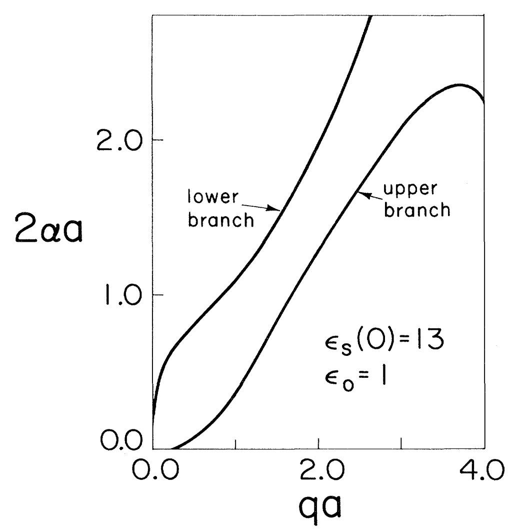

FIG. 2. Inverse penetration depth $2\alpha a$ vs q for the two surface-plasmon branches of Fig. 1.

and merges into it from below at a value of q^* again given by Eq. (7).

It is interesting to notice that for large values of qa the frequency of one of the surface-plasmon branches coincides with that of the two-dimensional plasmon of the first layer which is in this case given by $2^{1/2}\omega_P^{(e)}qa/(\epsilon_s+\epsilon_0)$.

In systems in which the dielectric background is polarizable, as is the case in polar semiconductors, it is necessary to allow for a frequency-dependent ϵ_s. A suitable form is provided by

$$\epsilon_s(\omega)=\epsilon(\infty)(\omega^2-\omega_L^2)(\omega^2-\omega_T^2)^{-1} ,$$

where $\epsilon(\infty)$ is the high-frequency dielectric constant, and ω_L and ω_T are the longitudinal and the transverse-optical-phonon frequencies, respectively. In this case the spectrum of the longitudinal excitations consists of four branches of coupled surface plasmon-optic-phonon modes.[6] This coupling causes the bulk longitudinal phonon frequency to split into two separated phonon bands. If ω_L is substantially larger than both $\omega_P^{(e)}$ and $\omega_P^{(h)}$ the lower two branches correspond to the surface plasmon, whereas the higher-frequency ones are essentially phononlike. A detailed discussion of this problem will be reported in a forthcoming paper.[8]

In conclusion, we have shown that the longitudinal spectrum of type-II semiconducting superlattices is characterized by an acoustic branch of surface plasmons. These modes could be observed by means of the familiar methods of the attenuated total reflection, resonant Raman and Brillouin scattering, or electron energy loss. It is important to emphasize here that, unlike the case of type-I superlattices, it is not necessary to achieve in the experiment a relatively large in-plane wave-vector transfer as these surface modes are well defined as q goes to zero.

As mentioned above these peculiar acoustic surface plasmons are long lived because they do not suffer Landau damping. This property, together with the fact that their group velocity can be adjusted by sample preparation, makes these modes potentially useful for surface wave devices.

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[1]L. Esaki, in *Novel Materials and Techniques in Condensed Matter,* edited by G. W. Crabtree and P. Vashishta (North-Holland, New York, 1982), p. 1.

[2]A. L. Fetter, Ann. Phys. (N.Y.) 81, 367 (1973); S. Das Sarma and J. J. Quinn, Phys. Rev. B 25, 7603 (1982); W. L. Bloss and E. M. Brody, Solid State Commun. 43, 523 (1982).

[3]A. C. Tselis, G. Gonzalez De La Cruz, and J. J. Quinn, Solid State Commun. 46, 779 (1983); G. Gonzalez De La Cruz, A. C. Tselis, and J. J. Quinn, J. Chem. Phys. Solids 44, 807 (1983); A. C. Tselis and J. J. Quinn (unpublished).

[4]D. Olego, A. Pinczuk, A. C. Gossard, and W. Wiegmann, Phys. Rev. B 25, 7867 (1982).

[5]Z. J. Tien, J. M. Worlock, C. H. Perry, A. Pinczuk, R. L. Aggerwal, H. L. Stormer, A. C. Gossard, and W. Weigmann, Surf. Sci. 113, 89 (1982).

[6]G. F. Giuliani and J. J. Quinn, Phys. Rev. Lett. 51, 919 (1983).

[7]G. F. Giuliani, G. Qin, and J. J. Quinn, *Proceedings of the 5th International Conference on the Electronic Properties of Two-Dimensional Systems, Oxford 1983,* [Surf. Sci. (in press)].

[8]G. F. Giuliani, G. Qin, and J. J. Quinn (unpublished).

[9]A more realistic model is discussed in Ref. 3; the surface plasmons associated with intersubband excitations will be considered in a subsequent paper.

[10]The dispersion relation of Eq. (1) is valid under assumptions that $q \ll p_F^{e,h}$, $p_F^{e,h}$ being the Fermi wave vector of the two-dimensional electron or hole gas in each layer, and $\omega \ll qp_F^{e,h}\hbar/m_{e,h}$.

[11]The alternative case of a hole-type first layer can be easily discussed by following the very same procedure.

[12]F. Stern, Phys. Rev. Lett. 18, 546 (1967).