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PLASMONS DISPERSION RELATION IN LAYERED SUPERCONDUCTORS AT FINITE TEMPERATURES

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Abstract

The electronic excitation spectrum for layered conductors has been studied within the random phase approximation (RPA). The plasmon dispersion relation of the layered system at $T = 0$ has been reported and discussed in earlier works. In this paper we report the effect of the temperature dependence and the inclusion of the leading higher order in wave vector q on the dispersion relation. We show that at low temperature limit the slope of the acoustic plasmon increases significantly due to the inclusion of the higher order in q while the thermal enhancement is nearly negligible. For a high temperature limit the plasmon dispersion relation is shown to be valid only in a very narrow range of the wave vector near zero and agrees quite well with the known classical limit of a two dimensional (2D) system.

Keywords: Acoustic plasmon, layered superconductors, dispersion relation, retarded polarization, finite temperatures

Introduction

A plasmon, or plasma excitation, is an elementary excitation of any charged particle system such as an electron system. It is a collective normal mode of charge density oscillation in the free carrier system. The collective behavior in an electron plasmon of lower dimensionality is quite different from that in three dimensional (3D) systems. The differences occur because the electric fields remain 3D while the induced charge densities have reduced dimensionality. The plasmon frequency ω_p , is exactly known only at long wavelengths ($q \rightarrow 0$) where the f -sum rule fixes

it to be given by classical electrodynamics. In a 2D system, $\omega_p = (2\pi n e^2/m)^{1/2} q^{1/2}$ and $\omega_p = (4\pi n e^2/m)^{1/2}$ in a 3D system where n is the 2D (3D) electron density and e^2 is interpreted as e^2/ϵ where ϵ is the background lattice dielectric constant. Therefore, the dispersion relation for the 2D electron system does not have a gap at $q \rightarrow 0$. At finite q , away from the long wavelength limit, there are several corrections to the plasmon dispersion arising from mechanisms relevant to the specific electron system being studied.

In a layered electron gas, the free charges are constrained to move on parallel planes spaced

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by a distance L . The plasmon frequency is predicted (Fetter, 1974a) to be

$$\omega_p(Q) = \left[\left(\frac{2\pi n e^2}{\epsilon m} q \right) \frac{\sinh qL}{\cosh qL - \cos q_z L} \right]^{1/2} \quad (1)$$

where q and q_z are the in-plane and normal components of the wave vector. Therefore, q and $Q = (q, q_z)$ are the 2D and 3D wave vectors respectively. Eqn. (1) exemplified the transition behavior between regimes of different dimensionality. For large separation between the planes ($qL \gg 1$), the dispersion reduces to that of a 2D plasmon with the characteristic square-root dependence. For long in-plane wavelengths ($qL \gg 1$) and with all the planes oscillating in phase ($q_z = 0$), the dispersion is similar to that of a 3D system. However, when $q_z \neq 0$ the contributions from induced electric fields in different planes tend to cancel and Eqn. (1) will be of the linear dependence form

$$\omega_p = q \left[\frac{2\pi n e^2}{\epsilon m} \frac{L}{1 - \cos q_z L} \right]^{1/2} \quad (2)$$

In this case ($qL \gg 1$ and $q_z \neq 0$) the dispersion is different from that in a 2D and 3D plasmon.

It is well known from experimental facts that all of high temperature superconductors (HTS) have been discovered in a number of families of compounds. These compounds have a layered structure characterized by copper-oxygen (CuO_2) layers in which the charge carriers are concentrated. It is also known that the layered superconductors exhibit a whole collection of unusual properties which has led to the various theoretical models to explain the high T_c 's observed. Although the various theoretical models so far differ in detail from each other, they seem to fall into two principal categories: one depending on the spin fluctuations of the charge carriers and the other on their charge fluctuations. Both seem to have achieved limited success as far as explaining the available experimental data. From the second

point of view, several non-phonon mechanisms have been proposed. Among the several non-phonon mechanisms is the plasmon exchange mechanism. It is assumed that the CuO_2 layer forms a two dimensional electron gas (2DEG) and that two electrons in a given layer can interact attractively by plasmon exchange, either within that layer or via the various neighboring layers. An isolated layer has only one optical plasmon mode with a dispersion relation $\omega_p \propto q^{1/2}$, while the interacting layers have additional modes of an acoustic nature or $\omega_p \propto q$ which enhance prospects for pairing.

The idea of acoustic plasmons goes back to the past (Pines, 1956), and, following they have been considered as a mechanism for superconductivity in the transition metals in the past (Fröhlich, 1968). The idea is that in the presence of two carrier species with very different effective masses, the light carriers can act to screen the Coulomb repulsion between the heavy carriers, and thereby lower the plasma frequency of the heavy carriers. The resulting plasmon mode will have a linear dispersion $\omega_p \propto q$ or acoustic plasmon. The acoustic plasmon mode is Landau-damped by decay into the electron-hole continuum of the light carriers, but for small wave vectors q the mode will be located within the lower range of the electron-hole continuum, so that the damping may be weak, and the acoustic plasmon may exist as a well-defined collective mode. Since there were no experimental evidences of acoustic plasmons to be confirmed at that time and the usual phonon mechanism was the only one able to explain superconductivity in the transition metals, interest in the topic began to wane in the late 1980's.

A layered electron-gas model for the cuprate HTS has been studied (Kresin and Morawitz, 1988). The interlayer Coulomb interaction leads to a noticeable modification of the pure 2D dispersion relation, namely, to the formation of plasmon bands. Anisotropy results in a picture which differs in a striking way from the isotropic 3D case. In addition, the maxima of the plasmon density of states are shifted toward the boundaries of the plasmon

branches, which enhances the electron-plasmon coupling considerably. The plasmon bands are restricted to lie between the upper and lower branches. These branches correspond to $q_z = 0$ and $q_z = \pi/L$ for the upper-most and the lowest branches, respectively. There is a crossover from 3D to 2D behavior in the region $q_z \sim 1/L$. This crossover corresponds to a maximum in the density of states (Landau and Lifshitz, 1976). The influence of acoustic plasmon modes in layered superconductors has been studied within the strong-coupling phonon-plasmon scheme (Bill *et al.*, 2000). Only the lowest acoustic branch at $q_z = \pi/L$ and the optical branch at $q_z = 0$, for which the density of states (DOS) is peaked, are considered. The optical branch has a smaller attractive interaction than the acoustic branch. Thus the presence of acoustic plasmons was shown to enhance noticeably the T_c of layered superconductors. The discovery of superconductivity in MgB_2 (Nagamatsu *et al.*, 2001) and strong evidence for the existence of an acoustic plasmon mode in MgB_2 (Voelker *et al.*, 2004), which leads to an attractive contribution, also confirm the significant role of the acoustic plasmon mode to enhance the superconducting transition temperature.

We would like to stress that the high T_c in cuprate HTS is due not to the plasmon mechanism by itself, but to its coexistence with strong electron-phonon coupling. There is experimental evidence that the electron-phonon interaction is very important, but it is not sufficient to lead to the observed high T_c . This coexistence concept has received various experimental support.

The aim of the presence paper is to evaluate the plasmon dispersion relation in layered conductors. The finite-thickness effect of a layered system, the finite-wave-vector nonlocal (higher order in q) effect and the temperature effect have been considered. We model the layered conductors as a system consisting of stacks of conducting sheets a long the z -axis spaced a distance L apart and separated by dielectric spacers with dielectric constant ϵ_M . The electrons are moving within the conducting sheets and can be considered as

2DEG. They interact with each other within the same layer, as well as from layer to layer via an effective Coulomb interaction which will be treated, in this paper, in the framework of the RPA. The finite-thickness effect on the electronic excitation spectrum comes from the effective Coulomb interaction. The finite-wave-vector nonlocal effect and the temperature effect come from the 2D nature of the polarization propagator.

Retarded Polarization as a Function of Temperatures

The plasmon dispersion can be calculated by finding poles of the density-density correlation function. Within the random phase approximation (RPA), the retarded polarization at finite wave vectors and finite temperatures, $\Pi^R(q, \omega; T)$ is given by, (Fetter and Walecka, 1971),

$$\Pi^R(q, \omega; T) = \frac{\Pi^{\text{OR}}(q, \omega; T)}{1 - V(q)\Pi^{\text{OR}}(q, \omega; T)} \quad (3)$$

where

$$\Pi^{\text{OR}}(q, \omega; T) = -\lim_{\eta \rightarrow 0} 2 \int \frac{d^2p}{(2\pi)^2} \left[\frac{f(p+q) - f(p)}{\hbar\omega + i\eta - \epsilon_{p+q} + \epsilon_p} \right] \quad (4)$$

is the non-interacting irreducible polarizability and $f(p)$ is the Fermi-Dirac distribution function. The factor 2 in Eqn. (4) arises from the spin degeneracy.

Since the retarded polarization has poles at the exact excitation energy of collective states of the interacting system, the singularities of $\Pi^R(q, \omega; T)$ in the complex ω -plane will determine the resonant frequencies of the system. The plasmon modes at finite wave vectors and temperatures are given by the zeros of the dielectric function

$$\epsilon(q, \omega; T) = 1 - V_c(Q)\text{Re}\Pi^{\text{OR}}(q, \omega; T) \quad (5)$$

where $V_c(Q)$ is the Coulomb interaction for conducting layered electron gas separated by spacers of dielectric constant ϵ_M . The Coulomb potential $V_c(Q)$ is the Fourier transform of the 3D Coulomb interaction $V_c(\mathbf{r}) = e^2/\epsilon_M|\mathbf{r}|$ and

can be written in the form

$$V_c(Q) = \frac{2\pi e^2}{\epsilon_M q} R(Q) \quad (6)$$

and

$$R(Q) \equiv R(q, q_z) = \frac{\sinh(qL)}{\cosh(qL) - \cos(q_z L)} \quad (7)$$

where $R(Q)$ is the dimensionless function, the layer form factor, introduced by Fetter (1974a) in the hydrodynamic treatment of the layer plasmon dispersion. This function is suitable for our model since, as mentioned

before, for large separation between the planes ($qL \gg 1$) the function $R(Q) \rightarrow 1$ and $V_c(Q)$ reduces to that of a 2D Coulomb potential. For long in-plane wavelengths ($qL \ll 1$) and $q_z = 0$, $V_c(Q)$ reduces to the 3D system. The form of $R(Q)$ given by Eqn. (7) has been applied to various layered systems such as metal-intercalated halide nitrides, layered organic materials and high- T_c oxides (Bill *et al.*, 2003), lanthanum cuprates (Varshney and Singh, 1995). Other forms of $R(Q)$ differ from Eqn. (7) have been suggested such as for a superlattice with a basis of two different kinds of metallic sheets (Griffin, 1988) which is different from our model.

The real part of $\Pi^{\text{OR}}(q, \omega; T)$ in Eqn. (4) can be written as

$$\text{Re}[\Pi^{\text{OR}}(q, \omega; T)] = -\frac{2}{(2\pi)^2} P \int d^2 p f(p) \left[\frac{1}{\hbar\omega - \epsilon_p + \epsilon_{p-q}} - \frac{1}{\hbar\omega - \epsilon_p - \epsilon_{p+q}} \right] \quad (8)$$

where P is the principal value. In order to keep the expression simple we introduce the

following dimensionless variables $x = \frac{p}{k_F}$, $t = \frac{T}{T_F}$, $\frac{x^2}{t} = z$, $\alpha = \beta\mu$

$$\text{and } \xi_{\pm} = \frac{1}{t} \left(\frac{m\omega}{\hbar q k_F} \pm \frac{q}{2k_F} \right)^2 \equiv \frac{A_{\pm}^2}{t}$$

where $\beta = 1/k_B T$ and μ is the chemical potential. Inserting the above expression into Eqn. (8), we obtain the following compact form for the real part of the noninteracting irreducible polarizability

$$\text{Re}[\Pi^{\text{OR}}(q, \omega; T)] = \frac{m}{\pi \hbar^2 q} \frac{k_F t^{1/2}}{2} \left[\int_0^{\xi_+} \frac{dz}{(e^{z-\alpha} + 1)(\xi_+ - z)^{1/2}} - \int_0^{\xi_-} \frac{dz}{(e^{z-\alpha} + 1)(\xi_- - z)^{1/2}} \right] \quad (9)$$

Low Temperature Limit

Since the real part of the polarizability will be a negative quantity for $\xi_{\pm} < \alpha$, $\xi_- < \alpha$ and $\xi_+ > \alpha$, there is no plasmon mode in these latter two cases. Only the case $\xi_{\pm} > \alpha$ will be considered. We can separate the last integral in Eqn. (9) into 3 parts as

$$\int_0^{\xi_-} \frac{dz}{(e^{z-\alpha} + 1)(\xi_- - z)^{1/2}} = \int_0^{\alpha} \frac{dz}{(\xi_- - z)^{1/2}} + \int_{\alpha}^{\xi_-} \frac{dz}{(e^{z-\alpha} + 1)(\xi_- - z)^{1/2}} - \int_0^{\alpha} \frac{dz}{(e^{\alpha-z} + 1)(\xi_- - z)^{1/2}}$$

and similarly for the integration of z from 0 to ξ_+ . In the small q limit, $\frac{m\omega}{\hbar q} \gg \frac{q}{2}$ and after integration over z we obtain

$$\text{Re}[\Pi^{\text{OR}}(q, \omega; T) = -\frac{m}{\pi \hbar^2 q} k_F \left\{ \frac{q}{k_F} + \left[(A_-^2 + 1)^{1/2} - (A_+^2)^{1/2} \right] \right. \\ \left. + \left[(A_+^2 - 1)^{-3/2} - (A_-^2 - 1)^{-3/2} \right] \frac{\pi^2 t^2}{24} \right\} \quad (10)$$

which is identical to the result obtained by Stern (Stern, 1967) when $t = 0$. Inserting the value of $A_{\pm}$ into Eqn. (10) we obtain the temperature dependence of the real part of the polarizability

$$\text{Re}[\Pi^{\text{OR}}(q, \omega; T) = \frac{q^2 k_F^2}{2\pi m \omega^2} \left[1 + \left(\frac{3}{4} + \frac{\pi^2 t^2}{4} \right) \left(\frac{\hbar q k_F}{m \omega} \right)^2 \right] \quad (11)$$

High Temperature Limit

In high t limit, we approximate the Fermi factor by the Boltzmann factor and hence Eqn. (9) becomes

$$\text{Re}[\Pi^{\text{OR}}(q, \omega; T) = \frac{m}{\pi \hbar^2 q} \frac{k_F t^{1/2}}{2} e^{\alpha} \left[\int_0^{\xi_+} \frac{e^{-z} dz}{(\xi_+ - z)^{1/2}} - \int_0^{\xi_-} \frac{e^{-z} dz}{(\xi_- - z)^{1/2}} \right] \\ = \frac{q^2 k_F^2}{2\pi m \omega^2} \left[1 + \frac{3}{2} t \left(\frac{\hbar q k_F}{m \omega} \right)^2 \right] \quad (12)$$

Plasmon Spectrum of a Layered Conductor

The spectrum of collective electronic excitations at finite wave vectors and temperatures is determined by the zeros of the dielectric function given by Eqn. (5). From Eqns. (6), (7), and (11) we obtain the general expression for the layer plasmon dispersion relation at low temperature limit:

$$\omega_p^2 = \left(\frac{4e^2 \epsilon_F}{\epsilon_m L} \right) \left(\frac{L}{2\hbar^2} \right) q R(Q) \left[1 + \left(\frac{3}{4} + \frac{\pi^2 t^2}{4} \right) \frac{\hbar^2 q}{m e^2} \right] \quad (13)$$

where $\epsilon_F = \hbar^2 k_F^2 / 2m$ and k_F is the in-plane Fermi wave-vector. The layered plasmon dispersion consists of two parts. The first part, that appears as prefactor of Eqn. (13), comes from the modified Coulomb interaction given by Eqn. (6) with the layer form factor $R(Q) \equiv R(qL, q_z L)$ given by Eqn. (7). This part reflects the finite-thickness effect of a layered system. The characteristic curves of the layer form factor $R(qL, q_z L)$ as a function of qL for five values of $q_z L$ is shown in Figure 1. It is found

that $R(qL, q_z L) \rightarrow 1$ for $qL \gg 1$ and $R(qL, q_z L) \rightarrow 2q / (q^2 + q_z^2)$ for $qL \ll 1$ and shows different curves for different values of $q_z L$. Furthermore, $V_c(Q)$ diverges as $1/|Q|^2$ for $|Q| \rightarrow 0$ in agreement with the 3D Coulomb interaction. Therefore, the prefactor of Eqn. (13) at $q_z = 0$ and $qL \ll 1$ reduces to the value $4\pi n e^2 / \epsilon_m L m$ of an optical plasmon. On the other hand, for $q_z L = \pi$ and, $qL \ll 1$, the prefactor of Eqn. (13) becomes a function of q^2 and the plasmon dispersion reduces to the acoustic plasmon ($\omega_p \propto q$). Therefore

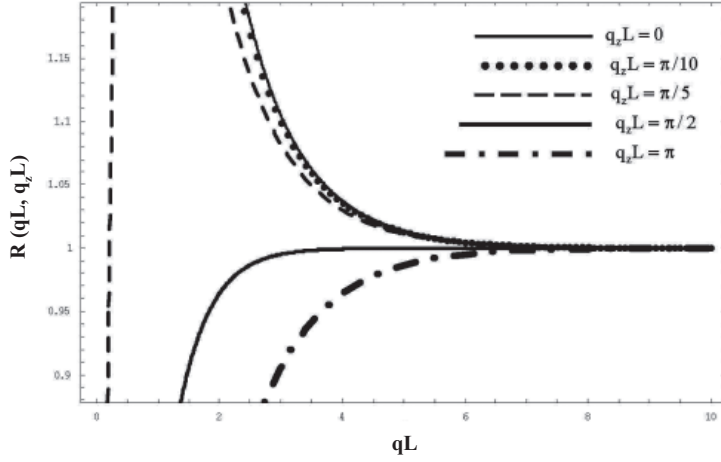


Figure 1. Characteristic curves of the layer from factor $R(qL, q_z L)$ given by Eqn. (7)

the plasmon band is confined between the upper branch with $q_z L = 0$ and the lower branch with $q_z L = \pi$. Finally, for large separation between the planes ($qL \gg 1$), $V_c(Q)$ becomes 2D Coulomb interaction and the prefactor of Eqn. (13) reduced to the value $2\pi e^2 q / \epsilon_m m$ of 2D plasmon.

The second part of the layer plasmon dispersion at low temperature limit is the terms in bracket of Eqn. (13). This part comes from the 2D nature of the polarization propagator $\Pi^{OR}(q, \omega, T)$ given by Eqn. (11). It contains the finite-wave-vector nonlocal (higher order in q) effect and the temperature effect. At $T = 0K$, the effect of the leading higher order in q , which is $(3/4) \hbar^2 q / me^2$, in agreement with the term $(3/4)(q/q_{TF})$ reported by Hwang and Sarma (2001) and the previous work (Rajagopal, 1977). The electronic excitation spectrum for the layered system at $T = 0$ is shown in Figure 2(b) and shows similar curves that reported by Bill *et al.* (2003). The finite-temperature effect, which is $(\pi^2 t^2 / 4) \hbar^2 q / me^2$, is quite small and negligible for $t \equiv T/T_F \ll 1$ as seen in Figure 2(c)-(f) and confirmed by Hwang and Sarma (2001)

From Eqns. (6), (7), and (12) we obtain the plasmon dispersion at high temperature limit:

$$\omega_p^2 = \left(\frac{4e^2 \epsilon_F}{\epsilon_m L} \right) \left(\frac{L}{2\hbar^2} \right) q R(Q) \left[1 + \frac{3}{2} t \frac{\hbar^2 q}{me^2} \right] \quad (14)$$

Since the finite-temperature effect comes from the 2D nature of $\Pi^{OR}(q, \omega, T)$ which is the term $(3/2)t \hbar^2 q / me^2$, in the bracket of Eqn. (14). For $t \gg 1$, this term is in agreement with the term $3q/k_D$ calculated by Fetter (1974b) for the classical electron surface layer ($k_D = 2\pi e^2 / k_B T$ is the 2D analog of the Debye-Hückel-screening constant). This term is also in agreement with the term $(3/2\sqrt{2})(T/T_F)(q/r_s)$ reported by Hwang and Sarma (2001) for a very-low-density electron system on the surface of liquid helium. It must be noted that this result is valid only in a very narrow range of q due to the condition $\xi_{\pm} > 1$ or $(me^2 / \hbar q t) < 1$. Since $t \gg 1$, this expression limits the range of q values very near zero. If $\xi_{\pm} > 1$ is not satisfied we do not get any plasmon mode as may be easily seen by the explicit calculation.

Discussion

Collective modes of a layered electron gas are characterized by the presence of acoustic plasmons in addition to the usual optical plasmon branch. The inter-layer Coulomb interaction induces such low-energy collective modes in layered conductors. The plasmon bands are confined between the optical branch at $q_z L = 0$ and the lower branch at $q_z L = \pi$. The layer form factor $R(Q)$ given by Eqn. (7) is an essential part of the Coulomb interaction. For

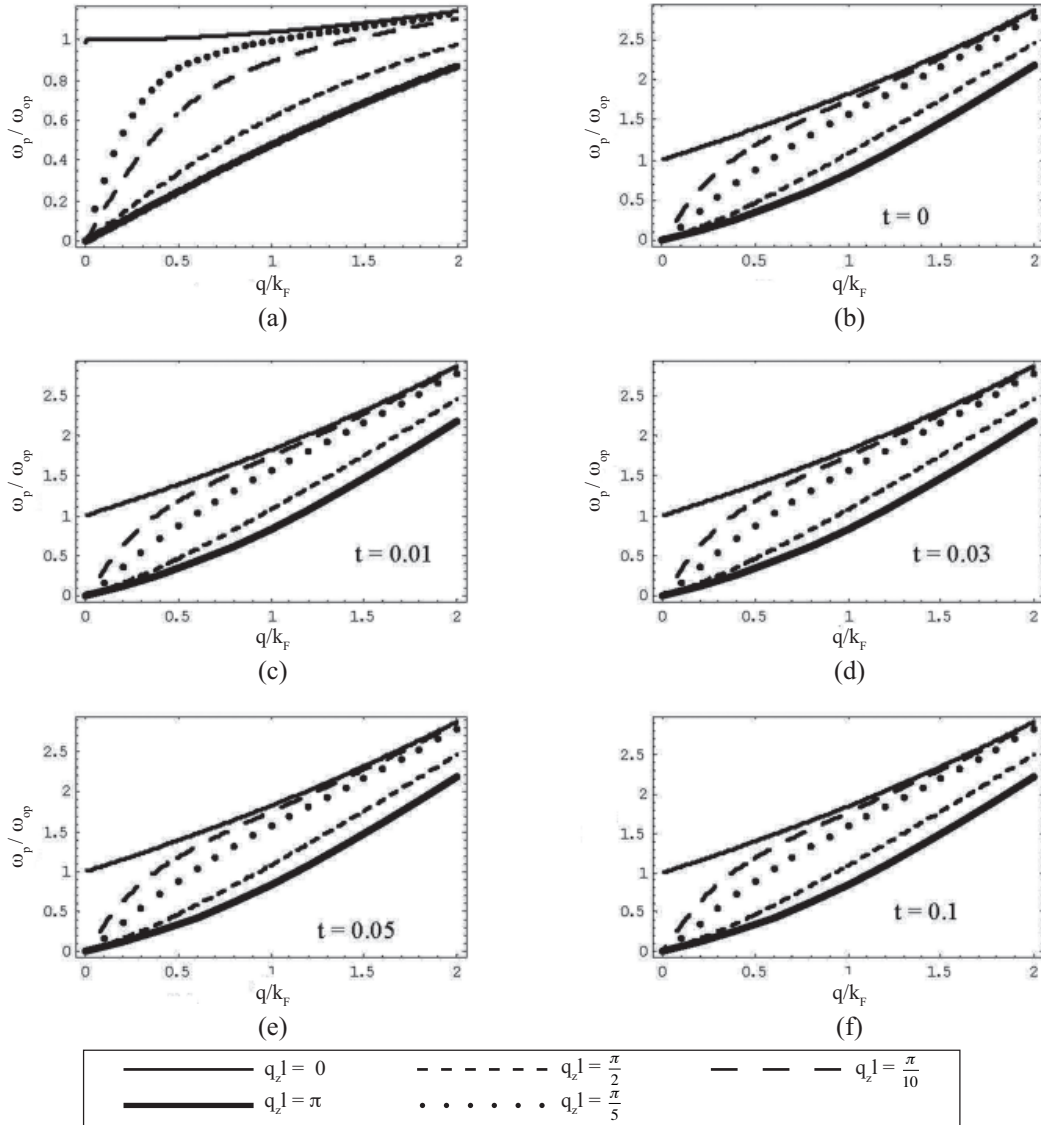


Figure 2. Layer plasmon dispersion as given by Eqn. (13) for which:

(a) containing only the first term

(b) with the second term at $t = 0$

(c) - (f) with the second term at various small t

large qL , $R(Q) \rightarrow 1$ as shown in Figure 1 and it shows different characteristics for small values of qL and $q_z L$. The layer plasmon dispersion relation at low temperatures given by Eqn. (13) contains an additional term (the second term in the bracket) which is to be compared with the previous results. Only the first term in the bracket

gives the same result as obtained by the previous works (Morawitz *et al.*, 1993; Bill *et al.*, 2003) and shown in Figure 2(a). The layer plasmon dispersion relation with and together with factor 3/4 in the second term is identical to the results by Rajagopal (1977) and Hwang and Sarma (2001), who obtained 2D plasmon

dispersion for $t = 0$, but includes the exchange correlation. The effect of the second term at $t = 0$ to the layer plasmon dispersion is shown in Figure 2(b). The dispersion relation at low temperatures is shown in Figure 2(c)-(f). The temperature effect is quite small and negligible for $t \ll 1$.

It is easily seen from Figure 2 that the inclusion of the second term in the dispersion relation makes the slope of the acoustic plasmon increase significantly. The finite-wave-vector nonlocal effect is therefore, important at low temperature limit. The temperature dependence shows a smaller effect compare to the effect by higher order in q . The dispersion relation at high temperature limit given by Eqn. (14) is proved to be valid compared with the classical limit of the 2D system such as the system of electrons on the surface of liquid helium, but valid only in a very narrow range of q near zero.

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