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Electron transport in a quasi-one-dimensional channel on suspended helium films

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Quasi-one-dimensional electron systems have been created using a suspended helium film on a structured substrate. The electron mobility along the channel is calculated by taking into account the essential scattering processes of electrons by helium atoms in the vapor phase, ripplons, and surface defects of the film substrate. It is shown that the last scattering mechanism may dominate the electron mobility in the low-temperature limit changing drastically the temperature dependence of the mobility in comparison with that controlled by the electron-ripplon scattering.

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I. INTRODUCTION

There has been great interest in the study of the quasi-one-dimensional electron system (Q1DES) produced on the liquid-helium surface[1–9] in which the motion of usual quasi-two-dimensional (Q2D) surface electrons (SE) (Ref. 10) is restricted in one more spatial direction. In most methods to realize the Q1D electron system, the charged channels are formed when parallel strips of a dielectric substrate are filled with liquid helium due to the capillary forces. The curvature radius R of the liquid into the strip can be varied in a wide range. The profile of the liquid surface across the channel (y direction) can be assumed to have a semicircular form $z_0 = R[1-\sqrt{1-(y/R)^2}] \simeq y^2/2R$ for $y \ll R$, where $R = \alpha/\rho g H$ with α and ρ are the surface tension and the helium density, respectively, g is the acceleration of the gravity, and H is the bulk helium level below the stripped structure. The electron confinement across the channel is achieved by the action of a holding electric field $E_\perp$ along the normal direction to the liquid surface (z axis). As a result, the electron near the channel bottom is subjected to the parabolic potential $U(y) = eE_\perp z_0 \simeq m\omega_0^2 y^2/2$, where $\omega_0 = (eE_\perp/mR)^{1/2}$.[11] In most previous experiments,[2] substrate effects can be discarded because the electron distance from the substrate are much larger than typical values of R. In this case, the electron is trapped in the direction perpendicular to the plane, by the combination of the image potential $-\Lambda_0/z$, where $\Lambda_0 = e^2(\varepsilon_{\mathrm{He}}-1)[4(\varepsilon_{\mathrm{He}}+1)]$ with $\varepsilon_{\mathrm{He}}$ the dielectric constant of the liquid helium, and the field potential $eE_\perp z$. The energy spectrum and the wave function for electron in the plane are $E_{n,k_x} = \hbar^2 k_x^2/2m + \hbar\omega_0(n+1/2)$ and $\chi_n(x,y) = [\exp(ik_x x)\exp(-y^2/2l^2)H_n(y/l)]/(\pi^{1/2} l L_x 2^n n!)^{1/2}$ respectively, where $l = (\hbar/m\omega_0)^{1/2}$, L_x is the system size along the channel axis, and $H_n(x)$ are the Hermite polynomials. The parameter l gives the scale of electron localization in the y direction (l yields 3.4×10^{-6} cm for $\omega_0 = 10^{11}$ Hz). Typical values of ω_0 are in the range 10^{10}–10^{11} Hz (for $E_\perp = 1$–3 kV/cm, and $R = 5\times 10^{-4}$ cm). Then the approximation $z_0 \simeq y^2/R$ is rather good and the condition $y \ll R$ is well satisfied and allows us to discard curvature effects in the hydrodynamic properties of the helium surface.[12]

Recently it was shown that very stable suspended helium films over a structured substrate can be created with arbitrary thickness d, as depicted in Fig. 1.[7,13] Under certain conditions (the distance between the elevations must be smaller than the capillary length), the helium film does not follow the substrate form but fills the depressions through capillary condensation. The value of d is controlled by the channel width W and the bulk level height H, and can vary in a wide range, from the bulk scale $\sim 10^{-4}$ cm down to small thicknesses $\lesssim 10^{-6}$ cm for van der Waals films. The method is useful to laterally confine the electrons in a submicrometer case and generate Q1DES above helium films.

The Q1DES over the helium film can be modelled in a similar form as that over bulk liquid, but an effective holding field $E_\perp^*$ replaces now $E_\perp$ in the confinement frequency ω_0 due to the contribution of the polarization interaction of the electron with the solid substrate: $E_\perp^* = E_\perp + (\Lambda_1/e)\langle 1|(z+d)^{-2}|1\rangle = E_\perp + (2\Lambda_1\gamma^2/e)F(\gamma d)$ where the angular brackets denote averaging over the electron wave function $\langle z|1\rangle = 2\gamma^{3/2} z\exp(-\gamma z)$ for the motion in the z direction (γ

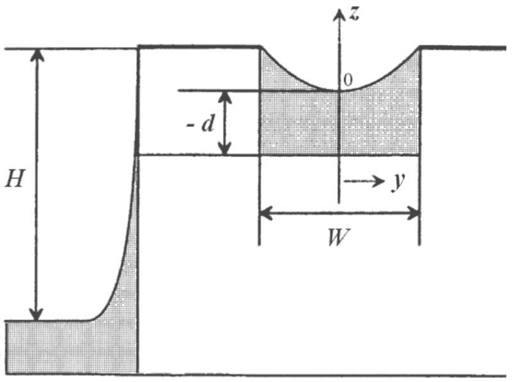

FIG. 1. Schematic diagram of a suspended film on a structured substrate.

is the localization parameter[15]), $\Lambda_1=e^2\varepsilon_{\mathrm{He}}(\varepsilon_s-\varepsilon_{\mathrm{He}})/[(\varepsilon_{\mathrm{He}}+1)^2(\varepsilon_s+\varepsilon_{\mathrm{He}})]$ with ε_s the dielectric constant of the substrate, respectively; d is film thickness near $y=0$, and $F(x)=[1+2x+4x(1+x)\exp(2x)Ei(-2x)]$, with $Ei(x)$ being the integral exponential function. According to our estimates, polarization effects become dominant for $d\lesssim 10^{-6}$ cm, especially for substrates with large ε_s. In such a condition the effective holding field $E_\perp^*$ reaches values $\lesssim 10$ kV/cm and the energy gap between the lowest and first excited levels in the z direction corresponds to a few tens K,[14] which allows to restrict to the ground-level contribution.

For wide enough rectangular strips, the curvature effects can be discarded especially for small d and high electron densities which lead to a flat profile in the center of the channel. However, charged electrodes may be arranged in such a way that the applied gate voltage produces an effective lateral electrostatic confinement.[6,9] The electrostatic potential should be calculated using the Poisson equation with appropriate boundary conditions. However, for $|y|\ll W$, the same model can be used, but $U(y)$ now depends on $\omega_{\mathrm{conf}}^2=\omega_0^2+\omega_{\mathrm{es}}^2$, where ω_{es} is the characteristic electrostatic frequency. Very recently, the experimental study of transport properties of the Wigner SE crystal confined in the conducting channel with effective width $W\simeq 10$ μm was reported.[16]

In this paper, we study the transport properties of the Q1DES over a helium film.[17] Despite different possible ways to create the system, the electron states inside the channel can be modelled for a flat helium surface by the potential

$$V(y,z)=eE_\perp z-\frac{\Lambda_0}{z}-\frac{\Lambda_1}{z+d}+\frac{m\omega_{\mathrm{conf}}^2y^2}{2}, \tag{1}$$

where the characteristics of the channel geometry are given by ω_{conf} within the harmonic approximation. As we will see, the electron multisubband spectrum leads to rather interesting transport properties along the channel. In Sec. II, we describe the electron scattering processes, and calculate the scattering potentials. In Sec. III, we obtain explicit formulas for the collision frequencies and discuss the results for the electron mobility.

II. SCATTERING POTENTIALS

The main scattering mechanisms are the electron interaction with atoms in the vapor phase predominating at $T>1$ K, the electron-ripplon interaction at lower temperatures, and the electron scattering by surface defects at the helium-substrate interface ($z=-d$). Quite recently it was shown that the latter scattering can dominate the 2DSE mobility over a helium film in the low-temperature regime.[18]

The electron-atom interaction can be calculated by a contact-type pseudopotential $U_g\delta(\mathbf{r}-\mathbf{R}_a)$ with $U_g^2=\pi\hbar^4A/m^2$, where $A\simeq 4.676\times 10^{-16}$ cm^{-2} is the cross section for the electron-atom scattering, $\mathbf{r}$ and $\mathbf{R}_a$ are two-dimensional (2D)-electron and 3D-atom positions, respectively.[19]

The electron-ripplon interaction is related to the variation of the potential energy for the electron at $z>0$ due to oscillations of the vapor-liquid interface. The adiabatic theory of electron-ripplon interaction was developed by Shikin and Monarkha.[20] Here we generalize Ref. 20 to take into account the roughness of the substrate surface. The electron potential can be written as $V(r,z)=V(z)+\delta V(r,z)$ where $V(z)$ is given by the first terms of Eq. (1) and

$$\delta V(\mathbf{r},z)=-\frac{1}{\pi}\int d\mathbf{r}'\left[\Lambda_0\int_0^{\xi_1(\mathbf{r}')}\frac{dz'}{(|\mathbf{r}-\mathbf{r}'|^2+|z-z'|^2)^2}+\Lambda_1\int_{-d}^{-d+\xi_2(\mathbf{r}')}\frac{dz'}{(|\mathbf{r}-\mathbf{r}'|^2+|z-z'|^2)^2}\right]. \tag{2}$$

Here $\xi_1(\mathbf{r})$ and $\xi_2(\mathbf{r})$ are the displacements of the free helium surface and the profile of the substrate, respectively. By Fourier transforming $\xi_j(r)=S^{-1/2}\Sigma_{\mathbf{q}}\xi_{j\mathbf{q}}\exp(i\mathbf{q}\cdot\mathbf{r})$ and expanding Eq. (2) up to first-order terms, one can rewrite Eq. (2) as

$$\delta V(\mathbf{r},z)=\delta V_1(\mathbf{r},z)+\delta V_2(\mathbf{r},z), \tag{3a}$$

where

$$\delta V_1(\mathbf{r},z)=-\frac{\Lambda_0}{\sqrt{S}}\sum_{\mathbf{q}} q\frac{K_1(qz)}{z}\xi_{1\mathbf{q}}\exp(i\mathbf{q}\cdot\mathbf{r}) \tag{3b}$$

and

$$\delta V_2(\mathbf{r},z)=-\frac{\Lambda_1}{\sqrt{S}}\sum_{\mathbf{q}} q\frac{K_1[q(z+d)]}{z+d}\xi_{2\mathbf{q}}\exp(i\mathbf{q}\cdot\mathbf{r}). \tag{3c}$$

Here S is the area of the liquid surface and $K_1(z)$ is the modified Bessel function. $\delta V_2(r,z)$ comes from the deviation of the helium-solid boundary from the equilibrium position at $z=-d$ and should be omitted when we consider only the electron-ripplon scattering. The potential $\delta V(r,z)$ is a small perturbation of $V(z)$ if the conditions $|\xi_j(r)|<<\langle z\rangle$ and $|\xi_j(r)|<<d$ are fulfilled, where $\langle z\rangle=(3/2)\gamma^{-1}\sim 10^{-6}$ cm is the mean electron distance from the helium surface. However $\delta V_1(r,z)$ cannot be simply considered as the electron-ripplon scattering potential. Indeed, only long-wavelength ripplons with wave number $q\leqslant 2k$ do contribute to the electron-ripplon scattering and the characteristic scale of electron wave number k is its thermal value $k_T=\sqrt{mT}/\hbar\simeq 3.4\times 10^5$ cm^{-1} for $T=1$ K. In such a condition the electron follows adiabatically the shape of the liquid surface and its wave function is close to $f_1[z-\xi_1(\mathbf{r})]$.[20] To calculate the scattering matrix elements one should replace z by the variable $\zeta=z-\xi_1(r)$. After eliminating the ξ dependence in the wave equation and redefining z as the original transverse coordinate we arrive at

$$\tilde{V}(\mathbf{r},z)=V(z)+\delta V_{\mathrm{er}}(\mathbf{r},z)+\delta V_{\mathrm{ed}}(\mathbf{r},z), \tag{4a}$$

where

$$\delta V_{\mathrm{er}}(\mathbf{r},z)=\frac{1}{\sqrt{S}}\sum_{\mathbf{q}}\xi_{1\mathbf{q}}V_{r\mathbf{q}}(z)\exp(i\mathbf{q}.\mathbf{r}) \tag{4b}$$

and $\delta V_{\rm ed}(r,z)=\delta V_2(r,z)$ are now the electron-ripplon and electron-defect scattering potentials, respectively. Here $\xi_{1\mathbf{q}}=[\hbar q\tanh(qd)/2\rho\omega_q]^{1/2}(c_{\mathbf{q}}+c^{+}_{-\mathbf{q}})$ where $c_{\mathbf{q}}(c^{+}_{-\mathbf{q}})$ are ripplon annihilation (creation) operators. In Eq. (4b), we take $V_{r\mathbf{q}}(z)=eE^*_\perp$ neglecting the contribution of the polarization interaction of SE with helium which is proportional to Λ_0 and is quite small for thin films ($\lesssim 10^{-5}$ cm) and high $E_\perp$. Indeed, taking the solid neon ($\varepsilon_s=1.19$) as the substrate, one estimates that the contribution to $E^*_\perp$ due to the polarization of the substrate is approximately 600 V/cm for $d\simeq 10^{-6}$ cm. This contribution, which is proportional to Λ_1, is essentially larger for substrates with higher dielectric constant. Correspondingly, the contribution coming from the helium polarization to the electron-ripplon scattering is equivalent to that for a holding field of 200 V/cm,[21] and hence can be discarded for $E_\perp\gtrsim 1$ kV/cm.

The Fourier-transformed potential $V_{d\mathbf{q}}(z)$ is written as

$$V_{d\mathbf{q}}(z)=-\Lambda_1 e^2\frac{qK_1[q(z+d)]}{z+d}\tag{4c}$$

and does not depend on the form of the electron wave function since we obtain the same result for both $f_1[z-\xi_1(\mathbf{r})]$ and $f_1(z)$. Indeed, we can show that $\delta V_2(r,z)$ is not influenced by the variable change $\zeta=z-\xi_1(r)$ if we consider first-order terms in the integral expansion of Eq. (2). Because the electron-defect scattering potential is independent of the specific choice to the electron wave function and the de Broglie electron wavelength is of the same order or larger than the characteristic defect width, we use $f_1(z)$ to calculate the matrix elements of the electron-defect interaction. We emphasize that the electron does not follow adiabatically the form of the interface helium-substrate resulting in different scattering potentials $V_{d\mathbf{q}}(z)$ and $V_{r\mathbf{q}}(z)$. The situation changes if we consider the electron motion over a solid surface without the helium blanket. In such a condition the electron follows the form of the surface with defects in the long-wavelength scale and the structure of $V_{d\mathbf{q}}(z)$ is similar to that of $V_{r\mathbf{q}}(z)$.[22]

It must be stressed that the last term of Eq. (2) was obtained summing over the polarization potential for all atoms of the substrate for nonzero $\xi_2(\mathbf{r})$. The result is valid quantitatively when the dielectric constant of the substrate is small such as in solid hydrogen or solid neon, because we can discard, as a first approximation, the screening of the electric field inside the solid substrate. For substrates with larger ε_s, screening effects must be considered and only atoms near the surface are polarized. In this situation, a more detailed analysis of surface atoms contributions to the interaction potential is necessary.

III. COLLISION FREQUENCIES AND THE ELECTRON MOBILITY

The electron mobility in the Q1D channel over bulk helium was calculated in Ref. 12 taking into account the population of the excited subbands in the y direction when $\hbar\omega_{\rm conf}\lesssim T$ (note that $\hbar\omega_{\rm conf}\simeq 0.8$ K for $\omega_{\rm conf}\simeq 10^{11}$ Hz). However, the contribution of $n>1$ subbands can be discarded for $T\ll\hbar\omega_{\rm conf}$. Here we limit ourselves to this regime and neglect quantum-statistical effects (the Fermi energy is much smaller than T).

The electron mobility along the channel in the limit $\hbar\omega_{\rm conf}/T\gg 1$ is given as[12]

$$\mu=\frac{2}{\sqrt{\pi}}\frac{e}{m}\left(\frac{\hbar\omega_{\rm conf}}{T}\right)^{3/2}\int_0^\infty\frac{\sqrt{x}\exp(-\hbar\omega_{\rm conf}x/T)}{[\nu_g(x)+\nu_r(x)+\nu_d(x)]},\tag{5}$$

where $x=\hbar k_x^2/(2m\omega_{\rm conf})$. The collision frequencies $\nu_g(x),\nu_r(x)$, and $\nu_d(x)$ denote the electron collisions with vapor atoms, ripplons, and defects, respectively.

The collision frequency with helium atoms is given by $\nu_g(x)=3\hbar n_g\gamma A/8m\sqrt{x}$, where n_g is the volume concentration of helium atoms.[19] Note that the contribution of $\nu_g(x)$ becomes negligible for $T<1$ K because n_g decays exponentially with T. The electron mobility is dominated by contributions of $\nu_r(x)$ and $\nu_d(x)$ in that temperature range.

In order to calculate $\nu_r(x)$, we use the expression[12]

$$\nu_r(k_x)=\frac{2\pi}{\hbar S}\sum_{\mathbf{q}}|\langle 0|\exp(iq_yy)|0\rangle|^2|\times\langle 1|U_{r\mathbf{q}}(z)|1\rangle|^2(2N_q+1)\frac{q_x}{k_x}\delta(E_{k_x-q_x}-E_{k_x}),\tag{6}$$

where $E_{k_x}=\hbar^2k_x^2/2m$ and $U_{r\mathbf{q}}(z)=V_{r\mathbf{q}}(z)[\hbar q\tanh(qd)/2\rho\omega_q]^{1/2}$. Here $\omega_q^2=[(\alpha/\rho)q^3+g'q]\tanh(qd)$ is the ripplon dispersion law with $g'=g+3n_{He}\beta/(\rho d^4)$ where β is the van der Waals constant of the substrate, n_{He} is the liquid-helium volume concentration. Long-wavelength ripplons do mainly contribute to the scattering with $N_q\simeq 2T/\hbar\omega_q$. Straightforward calculation of Eq. (6) leads to

$$\nu_r(x)=\frac{e^2(E^*_\perp)^2T}{4\alpha\hbar^2\omega_{\rm conf}}\cdot\frac{\exp[4(x+x_c)]}{\sqrt{x^2+x_cx}}[1-{\rm erf}(2\sqrt{x+x_c})],\tag{7}$$

where $x_c=\hbar\rho g'/8\alpha m\omega_{\rm conf}$. In the limit $d\to\infty$ and taking $g=0$ one obtains

$$\nu_r(x)=\frac{e^2(E^*_\perp)^2T}{4\alpha\hbar^2\omega_0}\cdot\frac{\exp(4x)}{x}[1-{\rm erf}(2\sqrt{x})].\tag{8}$$

which reproduces the result of Ref. 11.

The general expression for $\nu_d(k_x)$ can be written as

$$\nu_d(k_x)=\frac{2\pi}{\hbar S}\sum_{\mathbf{q}}|\langle 0|\exp(iq_yy)|0\rangle|^2|\times\langle 1|V_{d\mathbf{q}}(z)|1\rangle|^2|\xi_{2\mathbf{q}}|^2\frac{q_x}{k_x}\delta(E_{k_x-q_x}-E_{k_x}).\tag{9}$$

To calculate the electron-defect contribution, we use the well-known Gaussian two-parameter model for surface defects in which the correlation function $\langle\xi_2(\mathbf{r})\xi_2(\mathbf{r}')\rangle=\xi_0^2\exp[-|\mathbf{r}-\mathbf{r}'|^2/a^2]$ depends on ξ_0 and a playing the roles of characteristic defect height and width, respectively.[23] This model leads to $\langle |\xi_{2\mathbf{q}}|^2 \rangle = \pi \xi_0^2 a^2 \exp(-q^2 a^2/4)$ and was used to explain the SE transport properties over a thin helium film[18] and over solid hydrogen.[22] The final expression for $\nu_d(x)$ is

$$\nu_d(x) = \frac{32 m \Lambda_1^2 \xi_0^2 a^2 \exp(-2a^2 x/l^2)}{\hbar^3 l^4 x^{1/2}} \Phi_d(x), \qquad (10)$$

where

$$\Phi_d(x) = \int_0^\infty \frac{dy}{\sqrt{y}} (x+y)^2 \exp\left[-\left(4 + \frac{2a^2}{l^2}\right) y\right] \varphi_d^2\left(\frac{\sqrt{2(x+y)}}{\gamma l}\right)$$

and

$$\varphi_d(x) = \frac{\exp(2\gamma d)}{x} \int_{2\gamma d}^\infty \frac{ds}{s} (s - 2\gamma d)^2 K_1(xs) \exp(-s).$$

For a ideal substrate interface $[\nu_d(x) = 0]$ the electron mobility is determined by $\nu_r(x)$ and is given by

$$\mu_r \simeq 6 \mu_\perp \left[1 + \frac{32}{3\pi} \left(\frac{T}{\hbar \omega_{\text{conf}}}\right)^{1/2}\right] \qquad (11)$$

for large enough d and x_c so small that the condition $T \gg \hbar \omega_{\text{conf}} x_c$ is fulfilled, even though $T \ll \hbar \omega_{\text{conf}}$. Here $\mu_\perp = \alpha \hbar / [em(E_\perp^*)^2]$. In the case of thin films with $T \ll \hbar \omega_{\text{conf}} x_c$, the asymptotic expression for the ripplon-limited mobility is

$$\mu_r \simeq \frac{8 \mu_\perp}{\sqrt{\pi}} \left(\frac{\hbar \omega_{\text{conf}} x_c}{T}\right)^{1/2} \frac{\exp(-4 x_c)}{1 - \mathrm{erf}(2\sqrt{x_c})}. \qquad (12)$$

By comparing Eqs. (11) and (12), one can see the drastic change in the temperature dependence of the mobility from a thick to a thin film. One estimates this transition at $d \sim 10^{-6}$ cm for actual substrate materials.

The situation becomes rather interesting when one includes the defect contributions of solid substrates. Indeed, for realistic values of $a = 10^{-6}$ cm, d in the same range, and $\xi_0 = 10^{-7}$ cm, $\nu_d(x)$ is near two orders of magnitude larger than $\nu_r(x)$ for $x = x_T = T/\hbar \omega_{\text{conf}}$ and gives the major contribution to the integral of Eq. (5). In such a condition, the defect-limited mobility is given as

$$\mu_d \simeq \frac{e}{m \nu_d^{(0)}} \left(\frac{T}{\hbar \omega_{\text{conf}}}\right)^{1/2}, \qquad (13)$$

which is valid for the above-mentioned values of ξ_0 and a and $T \gtrsim 0.1$ K. Here $\nu_d^{(0)} = \pi m \Lambda_1^2 \xi_0^2 a^2 \gamma^4 F^2(\gamma d)/\hbar^3$. One should emphasize that the approach used to describe the electron scattering by surface defects is valid for smooth roughness with height significantly smaller than the helium film thickness. Furthermore trapped charges in the substrate may affect the electron transport and a comparison of the result given by Eq. (13) with experimental ones may be unreliable in such a condition.

The results for mobility, given by Eqs. (11)–(13), are obtained in the single-electron approximation (SEA). However, by increasing the electron density, the effects of electron-electron collisions may be pronounced and change transport properties. Usually, these effects are included through many-body approaches beyond the Hartree-Fock approximation for degenerate electron systems. For nondegenerate systems, like semiconductor plasmas, electron-electron collisions effects can be taken into account in the so-called complete control approach (CCA) or the Boltzmann shifted-distribution approximation, which consists of solving the equation $\hat{S}_{e-e}\{f_0(k)\} = 0$, where $\hat{S}_{e-e}$ is the collision integral for electron-electron scattering in the kinetic Boltzmann equation. In this case, scattering process is dominated by the electron momentum distribution because it is assumed that the electron-electron scattering frequency is much larger than the frequency for other scattering processes ($\nu_e >> \nu_r$, ν_g, and ν_d). As a consequence, we obtain $f_0(k) \sim \exp[-\hbar^2 k_x^2/2mT + \hbar k_x u/T]$ and the electrons have equal drift velocity u calculated from balance equation.[24] This approach has been successfully used in calculations of the electron mobility in the Q2DES over bulk helium which should be the unique low-dimensionality charge system where the complete control regime was realized at experimental conditions.[25] In particular, for the Q2DES over bulk helium, the results of CCA mobilities are in excellent agreement with experimental data and becomes twice smaller than the values predicted by SEA as the SE density increases. For larger electron densities, the highly correlated Q2DES must be described by transport theories in which the conductivity is obtained in terms of many-body quantities like the dynamic structure factor of the electron liquid[26] or the fluctuating electric field driven to the electron due to thermal fluctuations of the density.[27]

The method of calculation of the Q1DSE mobility in the CCA is described in detail in Ref. 12. The final expression for the ripplon-limited mobility in the CCA is

$$\mu_r^{(\text{cca})} \simeq 2 \mu_\perp \left[1 + \frac{4}{\pi} \left(\frac{T}{\hbar \omega_{\text{conf}}}\right)^{1/2}\right] \qquad (14)$$

for thick enough films where $T \gg \hbar \omega_{\text{conf}} x_c$. In the opposite limit $\mu_r^{(\text{cca})} = (\pi/4) \mu_r$, for $T \ll \hbar \omega_{\text{conf}} x_c$ where μ_r is given by Eq. (12). Comparing with Eqs. (11) and (12), one concludes that CCA gives the same qualitatively dependences of the ripplon-limited mobility on temperature and effective holding field as SEA. However, the absolute values of mobility are smaller in the CCA demonstrating the influence of electron-electron collisions on transport in Q1DES. The defect-limited mobility in CCA is $\mu_d^{(\text{cca})} = \mu_d/4$ where μ_d is given by Eq. (13).

In conclusion, we have investigated theoretically the properties of Q1DES over suspended helium films. We have derived the interaction potential for the electron scattering by interface defects. Film effects modify the confinement potential across the channel and the electron mobility at low temperatures is limited by ripplon scattering and mainly by surface defects at the helium film substrate interface. The latter scattering is dominant for thin films with $d \sim 10^{-6}$ cm and leads to the increase of electron mobility with temperature whereas the ripplon-limited mobility should decrease in this limit. Such a prediction can be tested in experimental attempts to observe the influence of different scattering mechanisms in the electron transport of Q1DES over helium film for temperatures below 1 K.

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