Phonon Scattering of Composite Fermions

D. V. Khveshchenko[1] and M. Yu. Reizer[2]
[1]*NORDITA, Blegdamsvej 17, Copenhagen DK-2100, Denmark*
[2]*Department of Physics, Ohio State University, Columbus, Ohio 43210-1106*
(Received 23 December 1996)

We study the principal aspects of the interaction between acoustic phonons and two-dimensional electrons in quantizing magnetic fields corresponding to even denominator fractions. Using the composite fermion approach we derive the vertex of the electron-phonon coupling mediated by the Chern-Simons gauge field. We estimate the acoustic phonon contribution to electronic mobility, phonon-drag thermopower, and hot electron energy loss rate, which all, depending on the temperature regime, are either proportional to lower powers of T than their zero-field counterparts or enhanced by the same numerical factor as the coefficient of surface acoustic wave attenuation. [S0031-9007(97)02968-2]

PACS numbers: 73.40.Hm, 63.20.Kr, 71.10.Ay, 71.10.Pm

The discovery of gapless compressible states at even denominator fractions (EDF) $\nu \sim 1/\Phi$, $\Phi = 2,\ 4$, etc. [1] motivated the theoretical idea [2] to describe these states as a new kind of Fermi liquid, which is formed by spinless fermionic quasiparticles named composite fermions (CFs). On the mean-field level the CFs, regarded as spin-polarized electrons bound to Φ flux quanta, experience zero net field and occupy all states with momenta $k < k_{F,\mathrm{cf}} = (4\pi n_e)^{1/2}$, where n_e is the 2DEG density, inside the effective CF Fermi surface. In the framework of the Chern-Simons theory of Ref. [2] the residual interactions of CFs appear as the gauge forces mediated by local density fluctuations. Conceivably, a 2D Fermi system governed by long-ranged retarded gauge interactions could demonstrate quite unusual properties and thereby provide an example of a genuine non-Fermi liquid (NFL).

Therefore, a qualitative success of the mean-field CF picture in explaining the experimentally observed Fermi liquidlike features at $\nu = 1/2$, $3/2$, and $3/4$ [1] caused a great deal of theoretical activity intended to reconcile these experimental observations with the implications of the CF gauge theory [3].

The present understanding of the situation is that the electrical current relaxation processes, which correspond to smooth fluctuations of the ostensible CF Fermi surface, can be safely described by means of the Boltzmann equation, where the singular self-energy and the Landau function terms largely compensate each other [4].

Yet, even in a hydrodynamic regime described by the Boltzmann equation for CFs with a finite (mean field) $m^*_{\mathrm{cf}} \sim k_{F,\mathrm{cf}}\epsilon_0/e^2$, one might expect that, when it comes to a coupling to another subsystem, the CF "marginal Fermi liquid" will behave differently from the standard Coulomb-interacting 2DEG at zero field. Such NFL deviations from the conventional behavior would then provide new tests for the CF theory.

As an example of this sort, it was recently shown [5] that a combined effect of the CF gauge interactions and impurity scattering leads to the experimentally observed nonuniversal $\ln T$ term in the resistivity, which is strongly enhanced compared to its zero-field universal counterpart [6].

In the present Letter we discuss another such example provided by the electron-phonon (e-ph) scattering at EDF.

In the well-studied zero-field case the electron-phonon scattering in GaAs heterostructures is dominated by the piezoelectric (PE) coupling at temperatures $T < 3$–4 K (see [7] and references therein).

We will concentrate on the range of temperatures below 1 K and treat phonons as bulk acoustic modes coupled to the local 2DEG density via the vertex $M_\lambda^{\mathrm{PE}}(Q) = eh_{14}(A_\lambda/2\rho u_\lambda Q)^{1/2}$, where $\vec{Q} = (\vec{q}, q_z)$ is the 3D phonon momentum, ρ is the bulk density of GaAs, u_λ is the speed of sound with polarization λ, h_{14} is the nonzero component of the piezoelectric tensor which relates a local electrostatic potential ϕ to a lattice displacement $\vec{u}$ ($\nabla_i\phi = eh_{ijk}\frac{\partial u_j}{\partial x_k}$), and the anisotropy factor A_λ is given by the formula [7] $A_l = \frac{9q_z^2q^4}{2Q^6}$, $A_{\mathrm{tr}} = \frac{8q_z^4q^2+q^6}{4Q^6}$.

It had been a long-standing issue of whether or not the 2DEG-3D phonon vertices get screened by the 2DEG [7], and it became customary to dress the bare 3D PE vertex with the static 2D dielectric function $\epsilon(q) = 1 + H(q)(2\pi e^2\nu_F/\epsilon_0 q)$, where $\epsilon_0 \approx 12.9$, $\nu_F = \frac{sm}{2\pi}$ is the density of states on the Fermi level proportional to the number of spin components s, and $H(q) = \int\int dz dz' \xi^2(z)\xi^2(z')e^{-q|z-z'|}$ is the form factor of the quantum well given in terms of the lowest occupied subband wave function $\xi(z) \sim ze^{-z/w}$.

In fact, the PE vertex undergoes a full dynamical screening, as follows from a systematic approach to the electron-phonon coupling as resulting from fluctuations of the Coulomb potential associated with lattice vibrations [8]. Summing up the RPA sequence of diagrams, one arrives at the expression $M_\lambda^{\mathrm{PE}}(Q)/\epsilon(\omega, q)$, where the dynamic dielectric function $\epsilon(\omega, q) = 1 + H(q)V_{e\text{-}e}(q)\Pi_{00}(\omega, q)$ is given in terms of the 2D Coulomb e-e interaction $V_{e\text{-}e}(q) = 2\pi e^2/\epsilon_0 q$ and the scalar 2D polarization $\Pi_{00}(\omega, q)$. In the presence of disorder characterized by the elastic transport time $\tau = l/v_F$, the expression for the polarization $\Pi_{00} = \nu_F(1 + i\omega/v_F q)$ obtained in the clean limit $ql > 1$ and $\omega\tau > 1$ changes to $\Pi_{00}(\omega, q) = \nu_F \frac{Dq^2}{i\omega + Dq^2}$ at $ql < 1$ and $\omega\tau < 1$ (here $D = v_F^2 \tau/2$ is the diffusion coefficient).

Then, provided the ratio $u/v_F < 1$ is fairly small (in what follows we will not distinguish between longitudinal u_l and transverse u_{tr} sound velocities while making qualitative estimates), the use of $\Pi_{00}(uQ, q)$ in the screened matrix element $M(Q)$ leads to the same results as the naive static screening at all $q > u/D$ which is equivalent to $T > T_1 = \tau u^2/l^2 \sim 0.1$ mK. The latter temperature is small compared to $T_2 = u/l \sim 1$ mK, below which one must use the expression for $\Pi_{00}(uQ, q)$ which contains the diffusion pole.

In turn, T_2 is much smaller than the Debye temperature $T_D = 2uk_F \sim 10$ K, below which both in-plane q and out-of-plane q_z components of the phonon momentum Q are controlled by temperature. Indeed, for typical electron densities $n_e \sim 10^{11}$ cm^{-2} one has $\kappa > k_F$, and then the width of the quantum well $w \sim (\kappa n_e)^{-1/3}$ provides a cutoff for $q_z \sim 1/w$ which is larger than k_F (in all our estimates throughout this paper we use the typical values of the parameters from [7]). Furthermore, at these densities the PE vertex is effectively screened $[|M_\lambda^{\mathrm{PE}}(Q)/\epsilon(\Omega_q, q)|^2 \sim q^2/Q]$ at all $T < T_D$.

Phonon-limited mobility.—The low-T momentum relaxation due to e-ph scattering is usually slow compared to the impurity transport rate $1/\tau_0$, which allows one to estimate the phonon contribution to the electronic mobility of the 2DEG by means of the standard formula [7],

$$\mu_{e\text{-ph}}^{-1} = \frac{2\pi m^*}{e} \sum_{\vec{Q},\lambda} \frac{|M_\lambda^{\mathrm{PE}}(Q)|^2}{|\epsilon(\Omega_Q, q)|^2} |F(q_z)|^2 \frac{\Omega_Q}{T} N\left(\frac{\Omega_Q}{T}\right) \times \left[1 + N\left(\frac{\Omega_Q}{T}\right)\right] \cos^2\theta\, \delta\left(\frac{q^2}{2m^*} - v_F q \cos\theta\right), \tag{1}$$

where $\Omega_Q = uQ$ is the dispersion of phonons distributed with $N(x) = (e^x - 1)^{-1}$ and $F(q_z) = \int dz\, \xi^2(z) e^{izq_z}$. In the Bloch-Gruneisen regime $T < T_D$, Eq. (1) yields $\mu_{e\text{-ph}}^{-1} \sim T^5$ which changes to a linear behavior above T_D [7]. The theoretical prediction of the nonlinear T dependence of the phonon-limited mobility was experimentally confirmed in [9].

The above dependence, however, can only hold at $T_2 < T < T_D$, whereas at lower T the Matthiessen's rule breaks down because of the effects of quantum interference between impurity scattering and e-ph interactions.

This low-T regime, which is hardly accessible in the zero-field case, becomes essentially more relevant in the case of CFs, since the absolute value of the resistivity at primary EDF is more than 2 orders of magnitude higher than at zero field [1], and therefore the CF transport time τ_{cf} is much shorter than the electronic one τ_0.

The low-T mobility measurements similar to those of [9] were recently performed at $\nu = 1/2$ [10]. The authors of Ref. [10] reported a stronger temperature dependence of $\mu_{\mathrm{cf\text{-}ph}}^{-1}$ consistent with T^3 at $T < T_{D,\mathrm{cf}} = \sqrt{2}\, T_D$. They also supported their findings by the results of the analytic calculation presented without derivation.

Since, to the best of our knowledge, a systematic analysis of the CF-phonon problem so far had not been made available, to facilitate our further discussion we first derive the effective CF-phonon vertex for CFs with Fermi momentum $k_{F,\mathrm{cf}} = \sqrt{2}\, k_F$, effective mass $m_{\mathrm{cf}}^* \sim 10 m_0$, where m_0 is the band electron mass in GaAs, and $\tau_{\mathrm{cf}} \sim (10^{-1}\text{–}10^{-2})\tau_0$.

The dynamical screening in compressible EDF states is described in terms of a polarization tensor $\Pi_{\mu\nu}(\omega, q)$ of CFs coupled to a 2D Chern-Simons (CS) gauge field $a_\mu = (a_0, \vec{a})$ [2].

In the Coulomb gauge ($\vec{\nabla}\vec{a} = 0$), the CF polarization is a 2×2 matrix $\hat{\Pi} = \mathrm{diag}(\Pi_{00}, \Pi_\perp)$ corresponding to the scalar a_0 and the transverse vector $a_\perp$ components of the CS gauge field. The CS gauge propagator $U_{\mu\nu}(\omega, q)$ depends on the actual form of the e-e interaction,

$$U_{\mu\nu}^{-1}(\omega, q) = \begin{pmatrix} 0 & iq/(2\pi\Phi) \\ -iq/(2\pi\Phi) & q^2 V_{e\text{-}e}(q)/(2\pi\Phi)^2 \end{pmatrix}. \tag{2}$$

The form of the scalar CF polarization $\Pi_{00}(\omega, q)$ is similar to that of ordinary electrons, which we discussed above. The transverse vector component is given by the formula $\Pi_\perp(\omega, q) = \chi_{\mathrm{cf}} q^2 + i\omega\sigma_{\mathrm{cf}}(q)$, where $\chi_{\mathrm{cf}} \sim 1/\nu_{F,\mathrm{cf}}$ and $\sigma_{\mathrm{cf}}(q)$ equals $\nu_{F,\mathrm{cf}} D_{\mathrm{cf}}$ if $ql_{\mathrm{cf}} < 1$ and $k_{F,\mathrm{cf}}/(2\pi q)$ otherwise.

Summing up the RPA sequence of polarization diagrams, we obtain that a PE scalar potential ϕ generated by lattice vibrations induces both scalar and vector components of the gauge field acting on CFs: $a_\mu = (1 - \hat{U}\hat{\Pi})_{\mu 0}^{-1} e\phi$.

Thus, the CF-phonon vertex acquires both the densitylike and currentlike parts,

$$M_\lambda^{\mathrm{cf}} = M_\lambda^{\mathrm{cf},s} + M_\lambda^{\mathrm{cf},v} = \frac{M_\lambda^{\mathrm{PE}}(Q)}{\epsilon_{\mathrm{cf}}(\omega, q)} \times \left[1 + (2i\pi\Phi) H(q) \frac{\vec{v} \times \vec{q}}{q^2} \Pi_{00}(\omega, q)\right], \tag{3}$$

where $\epsilon_{\mathrm{cf}} = 1 + H\Pi_{00}V_{e\text{-}e} + H^2(2\pi\Phi/q)^2 \Pi_{00}\Pi_\perp$. By virtue of the vertex (3) the CFs couple to both longitudinal (l) and transverse (tr) phonons, unlike the isotropic 3D case of electrons interacting via the electromagnetic gauge field, which only generates coupling to tr-phonons [8].

In the diffusive regime $ql_{\mathrm{cf}} < 1$ corresponding to $T < T_{2,\mathrm{cf}}$ the vertex $M_\lambda^{\mathrm{cf},s}$ gets dressed by the impurity ladder $\sim \tau_{\mathrm{cf}}^{-1}(D_{\mathrm{cf}} q^2 - i\omega)^{-1}$, whereas the vector part $M_\lambda^{\mathrm{cf},v}$ does not acquire such a pole. Even in the clean regime $ql_{\mathrm{cf}} > 1$ the latter remains unscreened ($|M_\lambda^{\mathrm{cf},v}|^2 \sim 1/Q$) at all $T > T_{3,\mathrm{cf}} = u^2 k_{F,\mathrm{cf}} \epsilon_0/e^2$.

In the range of temperatures $T_{3,\mathrm{cf}} < T < T_{D,\mathrm{cf}}$, where the limiting 3D momentum of thermally excited phonons is controlled by T and, at the same time, the dynamics of the

CFs remains nondiffusive, one can use Eq. (1) and obtain $\mu^{-1}_{\text{cf-ph}} \sim (h^2_{14}\epsilon^2_0/e^3\rho u^4 k_{F,\text{cf}})T^3$. The main contribution to $\mu^{-1}_{\text{cf-ph}}(T)$ comes from the vector part of (3).

In contrast to Eqs. (1) and (2) from [10] our expression for $\mu^{-1}_{\text{cf-ph}}(T)$ contains neither m^*_{cf} nor electron band mass in GaAs. Well below T_D the ratio $\mu_{\text{cf-ph}}/\mu_{e\text{-ph}}$ varies simply as $(T/T_D)^2$.

It is worthwhile mentioning that, in principal, it could exist as a range of temperatures $T_{2,\text{cf}} < T < T_{3,\text{cf}}$, where $|M^{\text{cf}}_\lambda|^2 \sim q^2/Q$ and $\mu^{-1}_{\text{cf-ph}} \sim T^5$. However, for typical parameters, $T_{3,\text{cf}}$ appears to be close to $T_{2,\text{cf}} \sim 300$ mK, which is about the lower bound of the temperature range where the reliable data were obtained [10].

At $T < T_{2,\text{cf}}$ the processes of small momenta transfers contribute to $\mu_{\text{cf-ph}}(T)$ as

$$\mu^{-1}_{\text{cf-ph}} = \frac{m^*_{\text{cf}}}{e} \operatorname{Im} \sum_{\vec{Q},\lambda} \int \frac{d\omega}{2\pi} |F(q_z)|^2 D_\lambda(\omega, Q) \times f\left(\frac{\omega}{T}\right)\left(\frac{|M^{\text{cf},v}_\lambda|^2}{(Dq^2 - i\omega)} + \frac{v^2_{F,\text{cf}} q^2 |M^{\text{cf},s}_\lambda|^2}{(Dq^2 - i\omega)^3}\right), \tag{4}$$

where $f(\frac{\omega}{T}) = \frac{\partial}{\partial\omega}[\omega \coth(\omega/2T)]$, $D_\lambda(\omega, Q) = [\omega - \Omega_\lambda(Q) + i0]^{-1} - [\omega + \Omega_\lambda(Q) + i0]^{-1}$ is the (retarded) phonon Green function. From this expression we obtain that, within the range $T_{1,\text{cf}} < T < T_{2,\text{cf}}$, the phonon-limited CF mobility varies as $\ln(T_{2,\text{cf}}/T)$ provided the ratio $u/v_{F,\text{cf}}$ is small enough. At $T < T_{1,\text{cf}}$ the correction ceases to grow logarithmically and shows only a $\sim T^2$ downward deviation from its $T = 0$ value $\mu^{-1}_{\text{cf-ph}} \sim (h^2_{14}\epsilon^3_0 k_{F,\text{cf}}/\rho u e^5 \tau_{\text{cf}}) \ln(T_{2,\text{cf}}/T_{1,\text{cf}})$.

Besides the nonuniversal prefactor given in terms of the PE coupling, the above term is down by an extra factor of $1/\sigma_{\text{cf}}$ compared to the $\ln T$ term resulting from interference between impurity scattering and the CF gauge interactions [5].

Surface acoustic wave (SAW) attenuation.—It was the SAW anomaly at $\nu = 1/2$ which provided the first evidence of the compressible CF states at EDFs [1] and inspired the formulation of the CF theory [2].

Following the procedure of Ref. [11], one can derive the coupling of electrons to SAW phonons from the bulk PE vertex: $|M^{\text{SAW}}(q)|^2 \sim \int dq_z |M^{\text{PE}}_\lambda(Q)|^2 |F(q_z)|^2$. By contrast to the case of bulk PE phonons, the vertex $M^{\text{SAW}}(q)$ remains finite at $q \to 0$.

The SAW attenuation is given by the imaginary part of the 2DEG density-response function $K_{00}(\omega, q) = \Pi_{00}(\omega, q)/\epsilon(\omega, q)$ which incorporates the effects of the dynamical screening [12]:

$$\Gamma_q = \frac{2\pi}{u} |M^{\text{SAW}}(q)|^2 \operatorname{Im} K_{00}(\Omega_q, q) \sim q \operatorname{Im} \frac{1}{1 + i\sigma(q)/\sigma_M}, \tag{5}$$

where in the last equation we used the standard definition of the complex momentum-dependent conductivity $\sigma(q) = i\sigma_M[1 - \epsilon(\Omega_q, q)]$ and $\sigma_M = \epsilon_0 u/2\pi \sim 5 \times 10^{-7}\ \Omega^{-1}$.

In the zero-field case, $\sigma_0(q) \sim 1/q$ at $ql_l > 1$, whereas the momentum-dependent conductivity of a CF state at EDF $\nu = 1/\Phi$ is inversely proportional to the CF quasiparticle conductivity [2]: $\sigma_\nu(q) \approx (e^2\nu/2\pi)^2/\sigma_{\text{cf}}(q)$, which implies that $\sigma_\nu(q) \sim q$ at $ql_{\text{cf}} > 1$. At small q both physical conductivities approach their static values, which typically satisfy the relations, $\sigma_0 \gg e^2/h \gg \sigma_\nu \sim \sigma_M$. In this regime the SAW attenuation becomes linear in momentum; $\Gamma_q = \gamma q$ and the coefficient γ_ν appears to be strongly enhanced compared to its zero-field counterpart γ_0:

$$\gamma_\nu/\gamma_0 = \frac{\sigma_\nu}{1 + \sigma^2_\nu/\sigma^2_M} \frac{1 + \sigma^2_0/\sigma^2_M}{\sigma_0} \approx \frac{\sigma_0\sigma_\nu}{\sigma^2_M + \sigma^2_\nu} \gg 1 \tag{6}$$

in agreement with the available data on SAW propagation [1].

Phonon-drag thermopower.—Thermoelectric measurements probing the low-T dynamics of the CF and their interactions with the phonons were recently reported [13,14]. Below 100 mK the thermopower (TEP) $S(T)$ measured at EDF corresponding to $\Phi = 2$ and 4 has an approximately linear behavior and is believed to be of the diffusion origin [13]. At higher T the measured TEP shows a nonlinear dependence which was assigned to the phonon-drag contribution S_g resulting from the momentum transfer from phonons, which acquire a net flux of momentum in the presence of a thermal gradient $\vec{\nabla} T$, to the CFs through their interaction.

In analogy with the standard theory of the e-ph interaction the thermoelectric effect can be treated in the framework of the Boltzmann equation [15], which yields the closed expression for the phonon-drag TEP:

$$S_g = \frac{\tau_{\text{ph}}}{e n_e T^2} \sum_{\vec{Q},\lambda} |M^{\text{PE}}_\lambda(Q)|^2 |F(q_z)|^2 \Omega^2_Q \frac{q^2}{Q^2} N\left(\frac{\Omega_Q}{T}\right) \times \left[1 + N\left(\frac{\Omega_Q}{T}\right)\right] \operatorname{Im} K_{00}(\Omega_Q, q) \tag{7}$$

derived under the assumption that phonons equilibrate by virtue of the boundary scattering, which is characterized by the relaxation time τ_{ph} proportional to the system size.

In the case of electrons Eq. (7) yields the known result $S_{g,e} \sim T^4$ [16], which holds in the clean regime $T > T_2$. Below T_2 it crosses over to $S_{g,e} \sim T^3$. According to the above discussion this can be viewed as the change from the momentum-dependent conductivity $\sigma_0(q) \sim 1/q$ to a constant one.

However, in the case of CFs we obtain $S_{g,\text{cf}} \sim T^2$ at $T > T_{2,\text{cf}}$ (provided that $T_{3,\text{cf}} < T_{2,\text{cf}}$) and $S_{g,\text{cf}} \sim T^3$ at $T < T_{2,\text{cf}}$. Remarkably, the exponent is higher in the dirty limit, as opposed to the situation at zero field. This is a direct consequence of the fact that in the clean regime the momentum-dependent EDF conductivity $\sigma_\nu(q)$ grows linearly with momentum.

Although in the regime of strong disorder both the zero field and the EDF phonon-drag TEP share the same temperature dependence $\sim T^3$, the prefactors are drastically different. A straightforward comparison between Eqs. (5) and (7) shows that the ratio $S_{g,\mathrm{cf}}/S_{g,e}$ is equal to that of the SAW attenuation coefficients (6).

The experimental data from [14] show, according to the authors, an "only marginally weaker T dependence for CFs than for electrons," which suggests that the CFs are well in the disordered regime. The data, fitted in [14] with $S_{g,e} \sim T^{4\pm0.5}$ and $S_{g,\mathrm{cf}} \sim T^{3.5\pm0.5}$, demonstrate a 2 order of magnitude enhancement of the prefactor in the CF case. As a remark, we note that the authors of Ref. [14] analyzed their data by using the expression for S_g derived for the case of e-ph coupling via deformation potential in the clean limit [15].

It also follows from our analysis that the above similarity of the T dependence of the phonon-drag TEP at zero field and at primary EDF is not, in fact, inconsistent with the drastic difference in the corresponding phonon-limited mobilities found in [10].

Hot electron energy loss rate.—Another informative experimental probe of the e-ph interaction is provided by measurements of an effective temperature of the 2DEG as a function of applied current $T_e(I) \sim I^{2/\alpha}$, where the exponent α characterizes the energy loss rate due to phonon emission, $P \sim T_e^{\alpha}$.

In the framework of the Boltzmann equation, $P(T)$ is given by the formula

$$P(T_e, T_l) = \frac{2\pi}{n_e} \sum_{\vec{Q},\lambda} |M_{\lambda}^{\mathrm{PE}}(Q)|^2 |F(q_z)|^2 \times \Omega_Q (e^{\Omega_Q/T_l} - e^{\Omega_Q/T_e}) \times N\left(\frac{\Omega_Q}{T_l}\right) N\left(\frac{\Omega_Q}{T_e}\right) \mathrm{Im}\, K_{00}(\Omega_Q, q) . \quad (8)$$

At zero field and $T_l \ll T_e$, Eq. (8) gives in the clean limit the standard result $P_e \sim T_e^5$, which is consistent with the inelastic e-ph scattering rate $\tau_{\mathrm{in},0}^{-1} \sim T^3$ [7]. In the disordered regime, it changes to a lower power $P_{el} \sim T_e^4$.

However, the situation at EDF again appears to be reversed: the power-law dependence $P_{\mathrm{cf}} \sim T_e^3$, which refers to the clean regime $T_{2,\mathrm{cf}} < T < T_{D,\mathrm{cf}}$ and implies the inelastic cf-ph scattering rate $\tau_{\mathrm{in},\nu}^{-1} \sim T$, changes to a greater power $P_{\mathrm{cf}} \sim T_e^4$ in the dirty limit $T < T_{2,\mathrm{cf}}$. Comparing the above results obtained in the regime of strong disorder we conclude that, just like the case of the phonon-drag TEP, the ratio P_{cf}/P_e is given by the factor (6).

Recently, the dependence $P \sim T_e^4$ was discussed in the context of transitions between adjacent quantum Hall effect plateaus (both integer and fractional) without any reference to CFs [17]. It is tempting to identify the nearly 2 order of magnitude enhancement (compared to the zero-field case) of the emission rate, which was observed at the transition between $\nu = 1/3$ and $\nu = 2/5$ [17], with the CF behavior governed by the EDF state at $\nu = 3/8$. A systematic experimental verification of this conjecture could lend additional support for the CF theory.

To summarize, we carry out a comparative analysis of the effects of the PE electron-phonon interaction in the 2DEG at zero magnetic field and at strong fields corresponding to EDF states viewed as the "CF marginal Fermi liquid." We show that in the latter case the acoustic phonon contribution to the electronic mobility, the phonon-drag TEP, and the energy loss rate for hot electrons, depending on temperature, either contain smaller powers of T or are enhanced by the numerical factor related to the ratio between the SAW attenuation at zero field and that at the EDF. Our results reconcile the seemingly contradicting conclusions which otherwise could have been drawn from the experimental data on phonon-limited mobilities [10] and phonon-drag TEP [14]. In addition to the already existing experimental observations, we predict a strong enhancement of the hot electron energy loss rate at EDFs, which is expected to be comparable with that of the SAW attenuation.

One of the authors (M. Yu. R.) acknowledges the support from US Office of Naval Research.

[1] R. L. Willett *et al.*, Phys. Rev. Lett. **65**, 112 (1990); H. W. Jiang *et al.*, *ibid.* **65**, 633 (1990); R. L. Willett *et al.*, *ibid.* **71**, 3846 (1993); W. Kang *et al.*, *ibid.* **71**, 3850 (1993); V. J. Goldman *et al.*, *ibid.* **72**, 2065 (1994).

[2] B. I. Halperin, P. A. Lee, and N. Read, Phys. Rev. B **47**, 7312 (1993).

[3] S. H. Simon and B. I. Halperin, Phys. Rev. B **48**, 17 368 (1993); B. L. Altshuler, L. B. Ioffe, and A. Millis, *ibid.* **50**, 14 048 (1994); D. V. Khveshchenko and P. C. E. Stamp, *ibid.* **49**, 5227 (1994); Y. B. Kim, A. Furisaki, X.-G. Wen, and P. A. Lee, *ibid.* **50**, 17 917 (1994); Y. B. Kim, P. A. Lee, X.-G. Wen, and P. C. E. Stamp, *ibid.* **52**, 17 275 (1995); A. Stern and B. I. Halperin, *ibid.* **52**, 5890 (1995); S. H. Simon, A. Stern, and B. I. Halperin, *ibid.* **54**, R11 114 (1996).

[4] Y. B. Kim, P. A. Lee, and X.-G. Wen, Phys. Rev. B **52**, 17 275 (1995).

[5] D. V. Khveshchenko, Phys. Rev. Lett. **77**, 362 (1996); Phys. Rev. B **54**, R14 317 (1996).

[6] L. P. Rokhinson *et al.*, Phys. Rev. B **52**, R11 588 (1995).

[7] W. Walukiewicz *et al.*, Phys. Rev. B **30**, 4571 (1984); **32**, 2645 (1985); S. J. Manion *et al.*, *ibid.* **35**, 9203 (1987); Y. Okuyama and N. Tokuda, *ibid.* **40**, 9744 (1989); T. Kawamura and S. Das Sarma, *ibid.* **45**, 3612 (1992).

[8] M. Yu. Reizer, Phys. Rev. B **40**, 5411 (1989).

[9] H. L. Stormer *et al.*, Phys. Rev. B **41**, 1278 (1990).

[10] W. Kang *et al.*, Phys. Rev. Lett. **75**, 4106 (1995).

[11] A. Knabchen *et al.*, Phys. Rev. B **54**, 10 696 (1996).

[12] A. L. Efros and Y. M. Galperin, Phys. Rev. Lett. **64**, 1959 (1990).

[13] X. Ying *et al.*, Phys. Rev. B **50**, 4969 (1994); V. Bayot *et al.*, *ibid.* **52**, 8621 (1995).

[14] B. Tieke *et al.*, Phys. Rev. Lett. **76**, 3630 (1996).

[15] D. G. Cantrell and P. N. Butcher, J. Phys. C **20**, 1985 (1987); *ibid.* **20**, 1993 (1987); M. J. Smith and P. N. Butcher, J. Phys. Condens. Matter **2**, 2375 (1990).

[16] S. K. Lyo, Phys. Rev. B **38**, 6345 (1988); S. S. Kubakaddi *et al.*, *ibid.* **40**, 1377 (1989).

[17] E. Chow *et al.*, Phys. Rev. Lett. **77**, 1143 (1996).