Phonon Scattering of Composite Fermions

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Abstract

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 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.
The discovery of gapless compressible states at even denominator fractions (EDF) \( \nu \sim 1/\Phi \), \( \Phi = 2, 4, etc \) became a new challenge for the theory of the interacting two-dimensional electron gas (2DEG) in quantizing magnetic fields. A number of metal-like features exhibited by strongly correlated EDF electronic states [1] motivated the theoretical idea [2] to describe these states as a new kind of Fermi liquid, which is formed by spinless fermionic quasiparticle named composite fermions (CFs). On the mean field level the CFs, regarded as spin-polarized electrons bound to \( \Phi \) flux quanta, experience zero net field and occupy all states with momenta \( k < k_{F,cf} = (4\pi n_e)^{1/2} \), where \( n_e \) is the 2DEG density, inside the effective CF Fermi surface. However, as a cost of such a simplification, the residual interactions of the CFs, as well as their interactions with charged impurities (remote ionized donors sitting some \( 10^2 \text{nm} \) apart from the 2DEG) turn out to be essentially more singular than the conventional Coulomb ones. In the framework of the Chern-Simons theory of Ref. [2] these interactions 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). Indeed, beyond the mean field approximation, the CF theory [2] predicts such NFL phenomenon as a divergence of the CF effective mass \( m_{\text{cf}}^* \) already in the lowest order of perturbation theory.

Therefore, a qualitative success of the mean field CF picture in explaining the experimentally observed Fermi liquid-like 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]. In particular, physical electromagnetic response functions in the low-frequency long-wavelength regime can be computed in the framework of the random phase approximation (RPA) while ignoring the
$m^*_{cf}$ divergence. The latter, however, is expected to manifest itself in processes which involve rough fluctuations of the CF Fermi surface. Those are responsible, for example, for the Subnikov-de-Haas-type oscillations of the magnetoresistivity $\Delta \rho_{xx}(B)$ in a residual field $\Delta B = B - 2\pi \Phi n_e$ in the vicinity of primary EDFs, such as $\nu = 1/2$. The theoretical predictions of an essentially NFL shape of the corresponding Dingle plot [5] is qualitatively consistent with experiment, although the actual divergence of $m^*_{cf}(\Delta B)$ derived from the data [6] is stronger than predicted.

Yet, even in a hydrodynamic regime described by the Boltzmann equation for CFs with a finite (mean field) $m^*_{cf} \sim k_{F,cf} \epsilon_0/e^2$, 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 [7] that a combined effect of the CF gauge interactions and impurity scattering leads to the experimentally observed non-universal $\ln T$ term in the resistivity, which is strongly enhanced compared to its zero field universal counterpart [8].

In the present Letter we discuss another such example provided by the electron-phonon 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 $T < 3 - 4K$ (see [10] and references therein). Above few Kelvin the coupling via deformation potential becomes important as well, whereas the coupling to optical phonons remains negligible up to $T \sim 40K$.

In what follows we will concentrate on the range of temperatures below $1K$ where the only important coupling is the PE one, and treat phonons as bulk acoustic modes coupled to the local 2DEG density via the vertex $M^P_F(Q) = \epsilon h_{14}(A\lambda/2\rho u_\lambda Q)^{1/2}$ where $\bar{Q} = (\bar{q}, q_z)$ is the 3D phonon momentum, $\rho$ is the bulk density of GaAs, $u_\lambda$ is the speed of sound with polarization $\lambda$, $h_{14}$ is the non-zero component of the piezoelectric tensor which relates a local
electrostatic potential $\phi$ to a lattice displacement $\vec{u}$: $\nabla_i \phi = e \hbar \epsilon_{ijk} \frac{\partial u_j}{\partial x_k}$, and the anisotropy factor $A_\lambda$ is given by the formulae \[10\]: $A_l = \frac{9 \epsilon_0^2 q^4}{2Q^6}$, $A_{tr} = \frac{8 \epsilon_0^2 q^4 + q^6}{4Q^6}$.

It had been a long-standing issue of whether or not the 2DEG-3D phonon vertices get screened by the 2DEG \[10\], 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 $\epsilon_0 \approx 12.9$, $\nu_F$ is the density of states on the Fermi level, and $H(q) = \int \int dzdz' \xi^2(z) \xi^2(z') e^{-q|z-z'|}$ is the formfactor 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 dynamic screening, as follows from a systematic approach to the electron-phonon coupling as resulting from fluctuations of the Coulomb potential associated with lattice vibrations \[9\]. Summing up the RPA sequence of diagrams, one arrives at the expression $M^{PE}_\lambda(Q)/\epsilon(\omega, q)$, where the dynamic dielectric function $\epsilon(\omega, q) = 1 + H(q) V_{e-e}(q) \Pi_{00}(\omega, q)$ is given in terms of the 2D Coulomb $e-e$ interaction $V_{e-e}(q) = 2\pi e^2/\epsilon_0 q$ 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 $\Pi_{00}(\omega, q) = \nu_F \frac{D q^2}{i\omega + D q}$ 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 $T_2 = u/l \sim 1mK$ 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 [10]). Furthermore, at these densities the PE vertex is effectively screened \(|M^P_{\lambda}(Q)/\epsilon(\Omega_q, q)|^2 \sim q^2/Q\) at all \(T < T_D\).

**Phonon-limited mobility**

The low-\(T\) the 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 [10]

\[
\mu_{e-ph}^{-1} = \frac{2\pi m^*}{e} \sum_{Q, \lambda} |M^P_{\lambda}(Q)|^2 |F(q_z)|^2 \frac{\Omega_q}{T} N\left(\frac{\Omega_q}{T}\right)(1 + N\left(\frac{\Omega_q}{T}\right)) \cos^2 \theta \delta\left(\frac{q^2}{2m^*} - v_F q \cos \theta\right)
\]

(0.1)

where \(\Omega_q = uQ\) is the dispersion of phonons distributed with \(N(x) = (e^x - 1)^{-1}\) and \(F(q_z) = \int dze^izq_z\). In the Bloch-Gruneisen regime \(T < T_D\) Eq.(1) yields \(\mu_{e-ph}^{-1} \sim T^5\) which changes to a linear behavior above \(T_D\) [10]. The theoretical prediction of the non-linear \(T\)-dependence of the phonon-limited mobility was experimentally confirmed in [11].

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 two orders of magnitude higher than at zero field [1,6], and therefore the CF transport time \(\tau_{cf}\) is much shorter than the electronic one \(\tau_0\).

The low-\(T\) mobility measurements similar to those of [11] were recently performed at \(\nu = 1/2\) [12]. The authors of Ref. [12] reported a stronger temperature dependence of \(\mu_{cf-ph}^{-1}\) consistent with \(T^3\) at \(T < T_{D,cf} = \sqrt{2}T_D\). 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 \(k_{F,cf} = \sqrt{2}k_F\), effective mass \(m_{cf}^* \sim 10m_0\), where \(m_0\) is the band electron mass in GaAs, and \(\tau_{cf} \sim 10^{-2}\tau_0\).
The dynamic 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 $a_\mu = (a_0, \vec{a})$.

In the Coulomb gauge ($\vec{\nabla} \vec{a} = 0$) the CF polarization is a $2 \times 2$ matrix $\tilde{\Pi} = \text{diag}(\Pi_{00}, \Pi_\perp)$ 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

$$U^{-1}_{\mu\nu}(\omega, q) = \begin{pmatrix} 0 & iq/(2\pi\Phi) \\ -iq/(2\pi\Phi) \end{pmatrix} q^2 V_{\epsilon-e}(q)/(2\pi\Phi)^2 \quad (0.2)$$

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 $\Pi_\perp(\omega, q) = \chi_{cf} q^2 + i\omega \sigma_{cf}(q)$, where $\chi_{cf} \sim 1/\nu_{F,cf}$ and $\sigma_{cf}(q)$ equals $\nu_{F,cf} D_{cf}$ if $ql_{cf} < 1$ and $k_{F,cf}/(2\pi q)$ otherwise.

Summing up the RPA sequence of polarization diagrams we obtain that a PE scalar potential $\phi$ generated by lattice vibrations induces both scalar and vector components of the gauge field acting on CFs:

$$a_\mu = (1 - \tilde{U}\tilde{\Pi})^{-1}_{\mu0} e\phi.$$ 

Thus, the CF-phonon vertex acquires both the density- and the current-like parts

$$M^{cf}_{\lambda} = M^{cf,s}_{\lambda} + M^{cf,v}_{\lambda} = M^{PE}_{\lambda}(Q) \epsilon_{\epsilon_{cf}}(\omega, q) (1 + (2i\pi\Phi)H(q)\vec{q} \times \vec{q} q^2 \Pi_{00}(\omega, q)) \quad (0.3)$$

where $\epsilon_{cf} = 1 + H\Pi_{00}V_{\epsilon-e} + H^2(2\pi\Phi/q)^2\Pi_{00}\Pi_\perp$. 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.

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

In the range of temperatures $T_{3,cf} < T < T_{D,cf}$ where the limiting 3D momentum of thermally excited phonons is controlled by $T$ and, at the same time, the dynamics of the CFs remains non-diffusive, one can use Eq.(1) and obtain $\mu_{\epsilon_{cf}^{-1} - ph} \sim (h_\uparrow^2 \epsilon_0^2/e^3\rho u^4 k_{F,cf}) T^3$. The main contribution to $\mu_{\epsilon_{cf}^{-1} - ph}(T)$ comes from the vector part of (3).
By contrast to Eq.(1-2) from \[12\] our expression for \(\mu_{cf-ph}^{-1}(T)\) contains neither \(m_{cf}^*\) nor electron band mass in GaAs. Well below \(T_D\) the ratio \(\mu_{cf-ph}/\mu_{el-ph}\) varies simply as \((T/T_D)^2\).

It is worthwhile mentioning that, in principal, it could exist a range of temperatures \(T_{2,cf} < T < T_{3,cf}\) where \(|M^{cf}_\lambda|^2 \sim q^2/Q\) and \(\mu_{cf-ph}^{-1} \sim T^5\). However, for typical parameters \(T_{3,cf}\) appears to be close to \(T_{2,cf} \sim 300\text{mK}\), which is about the lower bound of the temperature range where the reliable data were obtained \[12\].

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

\[
\mu_{cf-ph}^{-1} = \frac{m_{cf}^*}{e} \Im \sum_{Q,\lambda} \int \frac{d\omega}{2\pi} |M^{PE}_\lambda|^2 |F(q_z)|^2 D(\omega, Q) \left( \frac{|M^{cf,v}|^2}{(Dq^2 - i\omega)} + \frac{v_{F,cf}^2 q^2 |M^{cf,s}|^2}{(Dq^2 - i\omega)^3} \right) \frac{\partial}{\partial \omega} \left[ \omega \coth(\omega/2T) \right]
\]

(0.4)

where \(D(\omega, Q) = (\omega - \Omega_\lambda(Q) + i0)^{-1} - (\omega + \Omega_\lambda(Q) + i0)^{-1}\) is the (retarded) phonon Green function. From this expression we obtain that within the range \(T_{1,cf} < T < T_{2,cf}\) the phonon-limited CF mobility varies as \(\ln(T_{2,cf}/T)\) provided the ratio \(u/v_{F,cf}\) is small enough. At \(T < T_{1,cf}\) the correction ceases to grow logarithmically and shows only a \(\sim T^2\) downward deviation from its \(T = 0\) value \(\mu_{cf-ph}^{-1} \sim (\hbar^2 e_0^5 / puk_{F,cf}^5 \tau_{cf}^3) \ln(T_{2,cf}/T_{1,cf})\).

Besides the non-universal prefactor given in terms of the PE coupling, the above term is down by an extra factor of \(1/\sigma_{cf}\) compared to the \(\ln T\) term resulting from interference between impurity scattering and the CF gauge interactions \[7\].

**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. \[13\] one can derive the coupling of electrons to SAW phonons from the bulk PE vertex:

\[ |M^{SAW}(q)|^2 \sim \int dz |M^{PE}_\lambda(Q)|^2 |F(q_z)|^2. \]

By contrast to the case of bulk PE phonons, the vertex \(M^{SAW}(q)\) 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 \[14\]:

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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 \( \sigma_M = \epsilon_0u/2\pi \sim 5 \times 10^{-7}\Omega^{-1} \).

In the zero field case \( \sigma_0(q) \sim 1/q \) at \( q l_{el} > 1 \) whereas the momentum-dependent conductivity of a CF state at EDF \( \nu = 1/\Phi \) is inversely proportional to the CF quasiparticle conductivity \( \approx (e^2/2\pi)^2/\sigma_{cf}(q) \), which implies that \( \sigma_0(q) \sim q \) at \( q l_{cf} > 1 \). At small \( q \) both physical conductivities approach their static values, which typically satisfy the relations: \( \sigma_0 >> e^2/h >> \sigma_\nu \sim \sigma_M \). In this regime the SAW attenuation becomes linear in momentum: \( \Gamma_q = \gamma q \) and the coefficient \( \gamma_\nu \) appears to be strongly enhanced compared to its zero field counterpart \( \gamma_0 \):

\[
\gamma_\nu/\gamma_0 = \frac{\sigma_\nu}{1 + \sigma_\nu^2/\sigma_M^2} \frac{1 + \sigma_\nu^2/\sigma_M^2}{\sigma_0} \approx \frac{\sigma_0\sigma_\nu}{\sigma_M^2 + \sigma_\nu^2} >> 1
\]

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 [17,18]. 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 [17]. At higher \( T \) the measured TEP shows a non-linear 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 \( e-\text{ph} \) 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:

\[
S_g = \frac{\tau_{ph}}{e n e T^2} \sum_{\tilde{Q}} |M^P\tilde{E}(Q)|^2|F(q_z)|^2 \Omega^2 Q^2 N(\frac{\Omega Q}{T})(1 + N(\frac{\Omega Q}{T})) ImK_{00}(\Omega Q, q)
\]

(0.7)

derived under the assumption that phonons equilibrate by virtue of the boundary scattering, which is characterized by the relaxation time \( \tau_{ph} \) proportional to the system size.
In the case of electrons Eq.(7) yields the known result $S_{g,el} \sim T^4$, which holds in the clean regime $T > T_2$. Below $T_2$ it crosses over to $S_{g,el} \sim T^3$. 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.

On the contrary, in the case of CFs we obtain $S_{g,cf} \sim T^2$ at $T > T_{2,cf}$ (provided that $T_{3,cf} < T_{2,cf}$) and $S_{g,cf} \sim T^3$ at $T < T_{2,cf}$. 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 Eq.(5) and (7) shows that the ratio $S_{g,cf}/S_{g,el}$ is equal to that of the SAW attenuation coefficients (6).

The experimental data from [18] 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 [18] with $S_{g,el} \sim T^{4 \pm 0.5}$ and $S_{g,cf} \sim T^{3.5 \pm 0.5}$, demonstrate a two order of magnitude enhancement of the prefactor in the CF case. As a remark, we note that the authors of Ref. [18] 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 [12].

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 $\alpha$ 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
\[ P(T_e, T_l) = \frac{2\pi}{n_e} \sum_{Q, \lambda} |M^{PE}_\lambda(Q)|^2 |F(q_z)|^2 \Omega_Q(e^{\Omega_Q/T_l} - e^{\Omega_Q/T_e}) N(\frac{\Omega_Q}{T_l}) N(\frac{\Omega_Q}{T_e}) ImK_{00}(\Omega_Q, q) \]

At zero field and \( T_l << T_e \) Eq.(8) gives in the clean limit the standard result \( P_{el} \sim T_e^5 \), which is consistent with the inelastic \( e - ph \) scattering rate \( \tau_{in,0}^{-1} \sim T^3 \) [10]. However, in the disordered regime it changes to a lower power \( P_{el} \sim T_e^4 \).

On the contrary, the situation at EDF again appears to be reversed: the power-law dependence \( P_{cf} \sim T_e^3 \), which refers to the clean regime \( T_{2,cf} < T < T_{D,cf} \) and implies the inelastic \( cf - ph \) scattering rate \( \tau_{in,\nu}^{-1} \sim T \), changes to a greater power \( P_{cf} \sim T_e^4 \) in the dirty limit \( T < T_{2,cf} \). 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 \( P_{cf}/P_{el} \) 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 [19]. It is tempting to identify the nearly two 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 \( \nu = 2/5 \) [19], with the CF behavior governed by the EDF state at \( \nu = 3/8 \). A systematic experimental verification of this conjecture could lend an 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 EDF. Our results reconcile the seemingly contradicting conclusions which could of been drawn from the experimental data on phonon-limited mobilities [12] and phonon-drag TEP [18]. In addition to the already existing experimental observations, we...
predict a strong enhancement of the hot electron energy loss rate at EDF, which is expected to be commensurate with that of the SAW attenuation.
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