Magnetic field induced Coulomb blockade in small disordered delta-doped heterostructures

V. Tripathi and M. P. Kennett

1 Department of Theoretical Physics, Tata Institute of Fundamental Research, Homi Bhabha Road, Mumbai 400005, India and
2 Physics Department, Simon Fraser University, 8888 University Drive, Burnaby, BC, V5A 1S6, Canada

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At low densities, electrons confined to two dimensions in a delta-doped heterostructure can arrange themselves into self-consistent droplets due to disorder and screening effects. We use this observation to show that at low temperatures, there should be resistance oscillations in low density two dimensional electron gases as a function of the gate voltage, that are greatly enhanced in a magnetic field. These oscillations are intrinsic to small samples and give way to variable range hopping resistivity at low temperatures in larger samples. We discuss recent experiments where similar physical effects have been interpreted within a Wigner crystal or charge density wave picture.

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I. INTRODUCTION

The interplay between disorder and interactions in spatially inhomogeneous electronic states is an important ingredient in the physics of many strongly correlated electronic materials. Two dimensional electron gases (2DEGs) provide an ideal laboratory to gain insight into the effects of disorder and interactions as device properties such as disorder or carrier concentration can be tuned in the growth process, or by external means, such as a gate. Recent experiments on small, disordered, delta-doped devices suggest a new set of unusual resistance oscillations in a perpendicular magnetic field. The resistance as a function of gate voltage is featureless at zero magnetic field, but at non-zero fields develops peaks that are evenly spaced in gate voltage and whose magnitudes grow with increasing magnetic field whilst the positions of the peaks are relatively unaffected. Similar effects have been previously observed in low density 2DEGs, and arrays of quantum dots.

In earlier work on the same devices there was the remarkable observation that the electron tunneling distance (extracted from the magnetoresistance) is directly proportional to the average electron separation, in the 2DEG. This was used to argue in favor of a charge density wave (CDW) or Wigner crystal (WC) picture. The value of in these devices is around 5, which is much less than the prediction of for Wigner crystalization in two dimensions in a clean system, but close to the prediction of for disordered systems.

A number of theoretical studies have shown that the charge distribution in disordered delta-doped heterostructures at is likely to be neither Fermi liquid nor WC, but a droplet or “emulsion” phase. The simplest picture for the droplet phase is one where nonlinear screening by electrons is unable to dominate the disorder-induced potential barriers between regions of localized electrons. A droplet-like phase may also arise at comparatively higher electron densities in disordered quantum Hall insulators, due to the interplay of the localization effects of disorder and screening, where the screening is now sensitive to whether or not the electrons in different parts of the system belong to a partially or completely filled Landau level. A third, and somewhat different, picture proposed for low disorder is one of an “emulsion” or stripe-like phase with crystalline regions in an electron liquid background.

Nanoscale electronic inhomogeneity has also been observed in diverse strongly correlated electron systems. Some of the above mechanisms might be responsible for phase separation in many of these systems. In particular, electronic inhomogeneity in the superconductor BSCCO has been attributed to localization effects of the disorder in oxygen doping.

In this paper, we use the first of the above mentioned droplet pictures to argue that the experiments in manifest Coulomb blockade effects greatly enhanced by a magnetic field. We ignore quantum Hall physics in our treatment of the droplet phase taking note of the fact that compared to quantum Hall insulators believed to have a droplet phase, these experiments were performed on devices with strong disorder and low electron density, and no quantum Hall effect was seen at high fields. We make specific predictions about the conditions under which this novel magnetic field induced Coulomb blockade will occur, and find that our model is very successful in explaining resistance versus temperature data in Ref. In particular, such Coulomb blockade should be a signature of an electron droplet phase.

In previous work, we derived expressions for the physical parameters of electron droplets in the non-linear screening regime that is relevant to the experiments of interest here, and applied this picture to explain the experimentally observed density dependence of the tunneling distance without invoking a CDW picture.

We propose here that the resistance oscillations arise from the decrease of inter-droplet tunneling due to shrinking of the localization length in strong magnetic fields. The decrease in inter-droplet conductance is sufficient at larger magnetic fields to lead to a visible
The lengthscale on which potential fluctuations are face. This is to be contrasted with earlier predictions that the wavefunction in GaAs perpendicular to the surface is comparable with the droplet diameter. This scenario was studied in Ref. 16 where it was shown that the resistance between two droplets behaves as

\[
\frac{R(B)}{R(0)} = e^{(B/B_0)^2} \frac{1}{\cosh^2(B/B_1)},
\]

with \(B_0 \sim \phi_0/(\pi y_0 l_p)\), and \(B_1 \sim 2\phi_0/(\pi l_p^2)\). We estimate that the spread of the wavefunction under the barrier in the direction perpendicular to the tunneling is \(\delta \sim \sqrt{\pi r}\). \(B_1\) is the field below which interference effects are significant. The magneto-resistance data in Ref. 1 can be explained with Eq. (1) without invoking a CDW or WC scenario. In particular, \(B_0^{-2} \simeq A r^3\) which was the primary motivation for suggesting a CDW. The good agreement with experiment is strong evidence for the existence of electron droplets. Equation (1) expresses the inverse of the barrier transparency between droplets, and at large fields, the transparency decreases exponentially with increasing \(B\). This implies that droplets become isolated from each other with increasing field, and it is then natural to expect Coulomb blockade effects will strengthen with magnetic field, similar to those recently observed in a lattice of quantum dots.

### III. MAGNETIC FIELD AND DENSITY DEPENDENCE OF INTER-DROPLET TUNNELING

Having summarized the physical properties of the electron droplets, we briefly review their implications for magnetotransport. Unlike a dirty semiconductor where the electrons are localized at point-like impurity sites, the separation \(l_{ip}\) is comparable with the droplet diameter. This scenario was studied in Ref. 16 where it was shown that the resistance between two droplets behaves as

\[
\frac{R(B)}{R(0)} = e^{(B/B_0)^2} \frac{1}{\cosh^2(B/B_1)},
\]

with \(B_0 \sim \phi_0/(\pi y_0 l_p)\), and \(B_1 \sim 2\phi_0/(\pi l_p^2)\). We estimate that the spread of the wavefunction under the barrier in the direction perpendicular to the tunneling is \(\delta \sim \sqrt{\pi r}\). \(B_1\) is the field below which interference effects are significant. The magneto-resistance data in Ref. 1 can be explained with Eq. (1) without invoking a CDW or WC scenario. In particular, \(B_0^{-2} \simeq A r^3\) which was the primary motivation for suggesting a CDW. The good agreement with experiment is strong evidence for the existence of electron droplets. Equation (1) expresses the inverse of the barrier transparency between droplets, and at large fields, the transparency decreases exponentially with increasing \(B\). This implies that droplets become isolated from each other with increasing field, and it is then natural to expect Coulomb blockade effects will strengthen with magnetic field, similar to those recently observed in a lattice of quantum dots.

### IV. MAGNETIC FIELD INDUCED COULOMB BLOCKADE

We now consider the properties of the magnetic field induced Coulomb blockade. The charging energy of a single droplet, \(E_c\), differs from the bare value \(E_c^0 = e^2/(8\pi e_0\kappa R_p)\) (that of a single metallic sphere) if there is a gate voltage \(V_g\) that couples to the droplet through the capacitance \(C_g\) to give a gate charge \(q_g = C_g V_g\), which takes values in the interval \([0,1/2]\). For non-zero \(q_g\), \(E_c(q_g) = E_c^0(1-2q_g)\), and hence the charging energy takes values between 0 and \(E_c^0 \simeq 20\) K (for \(R_p \approx 30\) nm).

In the devices of interest there are many neighboring droplets which have a depolarizing effect, renormalizing the bare charging energy. There is charge screening for distances greater than \(R_c\), and if \(R_c \lesssim 7R_p\) (easily realized in experiment, since \(R_p \approx 0.6R_c\)), one can assume the droplet array only has nearest neighbor interactions. For a hexagonal array, we estimate the effective charging energy as \(E_c^{\text{eff}} \approx 0.22E_c^0 \approx 4.4\) K, and so expect strong renormalization of the droplet charging energy.

To calculate the excitation energy of an \(N_c\)-electron droplet, we need to consider the level separation \(\delta\) as well as the effective charging energy \((1-2q_g)E_c^{\text{eff}}\). Due
to the small size of the droplet, the Fermi energy $E_F$ can lie between levels $\varepsilon_{N_e+1}$ and $\varepsilon_{N_e}$, corresponding to $N_e+1$ and $N_e$ electrons in the droplet respectively. The energy to add an electron to the droplet is thus:

$$E_{\text{exc}}^{N_e+1,N_e} = (1-2q_B)E_c^{\text{eff}} + \min[\varepsilon_{N_e+1} - E_F, \varepsilon_{N_e} - \varepsilon_{N_e}] .$$

The gate voltage tunes both $q_B$ and $E_F$, so that the first term oscillates between 0 and $E_c^{\text{eff}}$, and the second between 0 and $\delta$. Equation (2) is for an isolated droplet, and tunneling into the droplet will reduce $E_{\text{exc}}$.

In a magnetic field, the energies $\varepsilon_{N_e}$ are modified due to both Zeeman splitting and orbital effects. For fields of the order of a tesla, as is the case in Ref. 1, the Zeeman splitting $g\mu_B B \ll \delta$ can be ignored. Orbital excitations will generically be affected by a magnetic field, and when the cyclotron energy $\hbar \omega_c = \hbar eB/m \sim \delta$, one can use Darwin-Fock theory to determine the single-particle energy levels. In our case, for $B = 1$ T, $\hbar \omega_c/2 \sim 10$ K is of the order of $\delta \approx 6$ K. However, at our level of analysis, such an improvement in accuracy does not strongly affect our results. This is also borne out by experiment, where the position of the Coulomb blockade peaks/troughs in a resistance versus gate voltage plot do not shift with field in the magnetic field range considered.

Treating electron droplets as effectively quantum dots, we now turn to consider the resistance that arises due to tunneling between droplets as a function of temperature. The tunneling rate $\Gamma$ between two droplets or between a droplet and leads is proportional to the corresponding transmission probability (and hence both droplet spacing and magnetic field) and the density of states on either side of the barrier. For sequential tunneling at temperatures $T$ such that $\hbar \Gamma \ll k_B T \ll \delta, E_{\text{exc}}$, Coulomb blockade gives rise to the droplet conductance:

$$G_{\text{CB}} = \frac{G_0}{\cosh^2 \left( \frac{E_{\text{exc}}}{2k_B T} \right)} , \quad G_0 = \frac{g_s e^2}{4\hbar k_B T} \frac{\Gamma \Gamma_r}{\Gamma_l + \Gamma_r} .$$

where $G_0$ is the peak value of the conductance of a droplet connected through (forward) scattering rates $\Gamma_{l,r}$ with its left and right neighbors, and $g_s$ is the spin degeneracy. At temperatures lower than $\hbar \Gamma/k_B$, $G_0$ saturates to $g_s e^2/h$. In Eq. (3), magnetic field controls $\Gamma$, whereas density (or $E_F$) controls $E_{\text{exc}}$. There are also parallel conductance channels, such as resonant co-tunneling, which is one order higher in the (small) tunneling probability and significant only at very low temperatures. For a single droplet (say the $i^{th}$), the transmission probability $T_{\text{cotunn}}$ associated with cotunneling is (when $k_B T \ll \hbar \Gamma$)

$$T_{\text{cotunn}}^{(i)} = \frac{\Gamma^{(i)} \Gamma_r^{(i)}}{2 \left( \Gamma_l^{(i)} + \Gamma_r^{(i)} \right) \left( E_{\text{exc}}/\hbar \right)^2 + \left( \Gamma^{(i)}/2 \right)^2} ,$$

where the total decay width $\Gamma^{(i)}$ includes inelastic as well as the elastic contributions, $\Gamma_{l,r}$. As inelastic processes are usually present, $\Gamma^{(i)} > \Gamma_l^{(i)} + \Gamma_r^{(i)}$.

For a string of $\mathcal{N}$ droplets, the total cotunneling transmission $T$ is the product of the cotunneling transmissions, $T = \prod_{i=1}^{\mathcal{N}} T_{\text{cotunn}}^{(i)}$ of individual droplets in the string, and hence the cotunneling conductance is $G_{\text{cotunn}} = g_s(c^2/\hbar)T$. This contribution is thus small unless there are uniformly spaced identical droplets, in which case there can be a contribution at resonance.

The relative importance of cotunneling and activated conduction changes as the droplet separation, $l_{ip} - 2r$, decreases towards $\xi$, implying that $\hbar^2$ approaches $E_{\text{exc}}$. This allows cotunneling to contribute to conductance at low temperatures. Assuming both contributions acting in parallel we estimate the droplet resistance as

$$R \simeq \frac{\cosh^2 \left( \frac{E_{\text{exc}}}{2k_B T} \right)}{G_0 + G_{\text{cotunn}} \cosh^2 \left( \frac{E_{\text{exc}}}{2k_B T} \right)} .$$

Equation (5) is not valid in the high or very low temperature limits since it implies $R \to 0$ as $1/T \to 0$, and does not include the fact that $G_0$ saturates as $T \to 0$. We use Eq. (5) to fit resistance versus $T$ data from Ref. 1, and in so-doing additionally assume a series resistance $R_0$, so that the total resistance $R = R_0 + R$. We also assume a phenomenological form for $G_{\text{cotunn}} = G_{\text{cotunn}}(T = 0)(1 + eT + aT^2)$, with $a$ and $c$ non-negative fitting parameters. Physically, the quadratic-$T$ dependence comes from inelastic cotunneling and the linear-$T$ dependence arises from the linear suppression of the tunneling density of states due to the Anderson orthogonality catastrophe when the droplet is coupled asymmetrically to the leads, which is the generic situation. We find excellent agreement with experiment at almost all densities, as is evident in Fig. 1 and we extract $E_{\text{exc}} \simeq 1$ K at most values of $n_e$. We note that extrapolation of the resistance in the Coulomb blockade regime to $B = 0$ leads to
a resistance that is always greater than \( h/2e^2 \), implying that there is, effectively, never more than one conductance channel open for transport, in agreement with the success of Eq. (5) in fitting the data.

As \( V_g \) is varied, the droplet energy levels cross \( E_F \), and since \( E_{\text{exc}} \) also depends on \( V_g \) [see Eq. (2)], the excitation energy for a droplet will vary between \( 0 \) and \( E_c^{\text{eff}} + \delta \). If there are many droplets, which we believe is the case here, we assume that the excitation energies are uniformly distributed in the interval \( [0, E_c^{\text{eff}} + \delta] \). In Fig. 2 we plot resistance as a function of \( n_c \) for illustrative purposes. We used Eqs. (3), (4) and (5), along with the schematic form \( E_{\text{exc}} = \frac{\hbar^2}{2m} \left[ 1 + f \cos \left( \frac{2\pi n_c}{n_0} \right) \right] \), where \( 0 \leq f \leq 1 \) (we choose \( f = 0.6 \) here), \( E_0 = 1 \text{K} \), \( T = 1 \text{K} \), \( \Delta n = 0.6 \times 10^{10} \text{cm}^{-2} \), and \( n_0 = 0.5 \times 10^{10} \text{cm}^{-2} \). We assume that \( \Gamma_I = \Gamma_r = \Gamma \) and the \( n_c \) and \( B \) dependence of \( \Gamma \) is determined by Eq. (1). The behaviour seen in Fig. 2 can be deduced very easily by looking at limits of Eq. (5). In the limit that \( G_0 \ll G_{\text{coulumn}} \), \( R \sim \frac{1}{G_{\text{coulumn}}} \), whereas when \( G_0 \gg G_{\text{coulumn}} \), \( R \sim \frac{1}{G_0} \cosh^2 \left( \frac{E_{\text{exc}}}{2k_BT} \right) \). We expect that both \( G_{\text{coulumn}} \) and \( G_0 \) will decrease with \( n_c \), but \( G_0 \) will decrease faster since \( f^{(0)} \) includes inelastic as well as elastic contributions, which are likely to be less density dependent. This would imply a crossover from oscillations to increasing resistance with decreasing density, exactly as shown in Fig. 2.

The temperature dependence of thermally activated transport depends on the size of the sample. For a sample large enough to contain many droplets there will be some with arbitrarily low activation energy, in which case transport will proceed via variable range hopping (VRH), and we do not expect to see Coulomb blockade oscillations. The crossover length for Arrhenius to Mott VRH behavior may be estimated as follows. For Mott VRH in 2D, the characteristic hopping distance \( D_{\text{Mott}} \) is given by \( D_{\text{Mott}} = (\xi/\pi n_c k_B T)^{1/3} \), where \( \nu \) is the density of available localized states per unit area:

\[
\nu \approx \frac{1}{\pi (\epsilon_p/2)^2 (\epsilon + E_c^{\text{eff}}/\hbar^2)}.
\]

This gives \( D_{\text{Mott}} \approx 60 T^{-1/3} \text{nm K}^{-1/3} \), and at \( T = 50 \text{mK} \), \( D_{\text{Mott}} \approx 160 \text{nm} \). In Ref. 1, an Arrhenius law in temperature is observed, implying that the active device area should be less than \( D_{\text{Mott}}^2 \). The total number of droplets \( N_p \) in an area \( \pi D_{\text{Mott}}^2 \) is

\[
N_p = \frac{4 \xi^2}{\epsilon_p} = \left( \frac{2\xi (\epsilon + E_c^{\text{eff}})}{\epsilon_p k_B T} \right)^{2/3},
\]

which is about 15 for the typical parameters we assume. Some low density samples do show VRH behavior, and hence we assume that the effective size of the sample is the order of \( D_{\text{Mott}} \), so the mean level separation in the sample is approximately \( \delta_{\text{sample}} \approx (\epsilon_c^{\text{eff}} + \delta)/N_p \), which is about 0.75K. The separation of the magnetoresistance peak from a trough corresponds to an energy scale of \( \approx 1 \text{K} \), in good agreement with both the activation energy deduced from experiment, and with \( \delta_{\text{sample}} \).

As \( V_g \) is varied, excited states in different droplets successively come into resonance; these resonances are associated with the minima of resistance. In between the minima, if the sample is small, no droplet in the system is in resonance with \( E_F \) and the resistance will show a maximum. The observability of resistance oscillations therefore crucially depends on the samples being small. This seems to be borne out in experiment.2

At large electron densities, the oscillations will be less visible for two reasons; firstly, as \( n_c \) increases, \( \epsilon_p \) decreases, as does \( E_c^{\text{eff}} \), which reduces \( \delta_{\text{sample}} \). This reduces the contrast between resonant and Coulomb blocked states. Secondly, the inter-droplet tunneling distance, \( D \propto 1/\sqrt{\epsilon_n} \), decreases and hence \( \Gamma \) increases. Ultimately, when the localization condition, \( D/\xi \geq 1 \), cannot be satisfied, the system becomes well-conducting and no Coulomb blockade oscillations are possible.

Increasing magnetic field reduces \( \Gamma \), with relatively little effect on excitation energies which improves the visibility of the Coulomb blockade. In experiment, the resistance oscillations tend to appear above a threshold field, which we estimate to be where the magnetoresistance switches from negative to positive, in the vicinity of \( B_1 \approx n_c \), where \( B_1 \) is defined below Eq. (1).

\section{V. Discussion}

In some respects, the issues discussed here are similar to observations in one dimensional wires of variations in the conductance periodic in \( n_c \). These were initially described in terms of a CDW, whereas later investigations appear to have convincingly demonstrated that the oscillations are due to Coulomb blockade effects. It was
found that the Coulomb blockade effects strengthened as a magnetic field was applied, consistent with the picture proposed here. However, the Coulomb blockade did not rely on magnetic field for visibility.

We did not discuss how non-linear screening is affected by the presence of a magnetic field. The screening of 2D electrons in a disordered potential in a magnetic field was discussed in Refs. [31,32], focusing on the regime where disorder is not too strong. We note that experiment appears to provide some of the solution. Measurements of localized states in the quantum Hall regime appears to provide some of the solution. Measurements of donor charges in the dopant layer – including these worsened the agreement with experiment. However, donor correlations are likely to be relevant in some cases.

In summary, we use the picture of electron droplets to establish that for low density 2DEGs in disordered delta doped heterostructures there can be a magnetic field induced Coulomb blockade. We provide evidence for this picture by using a model for resistance as a function of temperature based on the idea of electron droplets acting like quantum dots to successfully fit experimental data. The ideas we present here may have wider applicability in inhomogeneous strongly correlated electron systems.

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