Thermodynamic magnetization of a strongly interacting two-dimensional system

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We report thermodynamic magnetization measurements of a 2-dimensional electron gas for several high mobility Si-MOSFETs. The low-temperature magnetization is shown to be strongly sub-linear function of the magnetic field. The susceptibility determined from the zero-field slope diverges as $1/T^\alpha$, with $\alpha = 2.4 \pm 0.2$ even at high electron densities, in apparent contradiction with the Fermi-liquid picture.

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Magnetic properties of the strongly-interacting electron gas have long been a subject of intensive theoretical and experimental investigations. The Coulomb interaction, as reflected by the exchange term, favors parallel spins, and therefore leads to ferromagnetism, whenever it is strong enough. The strength of the interaction is determined by the ratio between the typical Coulomb and kinetic energies, and is customarily characterized by a dimensionless parameter $r_s = \rho/a_B$, $a_B$ being the Bohr radius in the material, and $\rho \sim n^{-1/3}$ is the effective distance between electrons in $d$ dimensions; in two dimensions $\rho = 1/\sqrt{n}$, $n$ being the electron density. When $n$ decreases, the relative effect of the interactions becomes stronger, contrary to naive expectations. At zero temperature and sufficiently low density the clean single-valley system of itinerant electrons was predicted to become ferromagnetic, a phenomenon called Stoner instability\textsuperscript{3}. In two dimensions, according to the Mermin-Wagner theorem, finite temperature destroys the ferromagnetic order.

In a real system there is always some degree of disorder, which favors an antiferromagnetic interaction. Indeed, the ground state of two localized spins is a singlet, much like a hydrogen molecule. A system of many localized spins preserves the tendency to order neighboring spins in opposite directions, that is, the coupling is antiferromagnetic; see\textsuperscript{2} for a review. Intensive investigations of magnetic properties of doped semiconductors, particularly phosphorus doped Si, in the 1980’s led to a substantial understanding of the interplay between interactions and disorder. The experimentally observed divergence of the susceptibility as $1/T^\alpha$, with $\alpha \approx 0.6$, on the low-density side of the metal-insulator transition has been well-understood\textsuperscript{2}; however the persistence of the divergence on the high-density side remains a puzzle.

Two-dimensional gated structures like MOSFETs allow for a gradual change in electron density, and hence interaction strength, without strongly affecting sample disorder. Interest in the magnetic properties of two-dimensional electron gas (2DEG) was stimulated by the observation of strong suppression of conductivity in Si MOSFETs by an in-plane magnetic field\textsuperscript{2}, which polarizes electrons and drives the system into an insulating phase. Scaling analysis of the magnetoresistance led the authors of\textsuperscript{2, 3} to suggest a quantum phase transition into a ferromagnetic state at the density of the metal-insulator transition, $n_c$\textsuperscript{3}. This claim was contested in\textsuperscript{2, 4} and\textsuperscript{3} on the basis of thermodynamic and Shubnikov-de Haas measurements, respectively.

In this Letter we use the technique developed and described in\textsuperscript{2} to study the magnetization at elevated temperatures. This technique measures the recharging of the gate-to-2DEG capacitor due to a change in the 2DEG chemical potential $\mu$. By modulating an external magnetic field with the amplitude $\delta B(\omega)$, while keeping the gate voltage constant we are able to measure a gate charge modulation $\delta Q(\omega)$ given by

$$
\delta Q(\omega) = C(\omega) \frac{\partial \mu}{\partial B} \delta B(\omega),
$$

(1)

where the capacitance $C(\omega)$ is measured independently by modulating the gate voltage at the same frequency. We note that all the quantities in Eq. (1) also depend on the electron density $n$; this dependence is not explicitly indicated. By virtue of the Maxwell relation $\partial \mu/\partial B = -\partial m/\partial n$, $\partial \mu/\partial B$ can be expressed as the derivative of the magnetization per unit area $m$ with respect to the density. In principle, integrating $\partial m/\partial n$ over the density from $n = 0$ would give $m(n)$.

In this Letter we present measurements from deep into the insulating region to far into the metallic phase, allowing us to integrate $\partial m/\partial n$ from almost $n = 0$. The magnetization $m(n)$ thus obtained exhibits a surprisingly strong temperature dependence at low magnetic field $\mu_B B \ll T$ even at high densities. This result contradicts expectations from the Fermi liquid theory. Indeed, at high densities interactions are relatively weak, therefore magnetization should be temperature independent at temperatures small compared to the Fermi energy $E_F$.

Extension of the measurements deep into the insulating regime became possible due the lower sample resis-
tance at elevated temperatures, together with good quality of the contacts. It is important to note that the expression (1) holds even if the capacitance drops and acquires imaginary part due to the contact and channel resistance, which happens at the resistance about 500 MOhms. This facilitates measurements deep into insulating region down to less than half the density of the metal-insulator transition in the sample.

Our measurements were performed on several high-mobility Hall bar shaped Si-MOS structures with 2DEG located on the (100) interface between Si surface and SiO$_2$. Such a 2DEG possesses a two-fold valley degeneracy in addition to its spin degeneracy. We used two groups of samples: those fabricated in Russia (R), similar to those used in Ref. [8, 9, 10] and those made in Holland (H), used in Ref. [11, 12]. The in-plane magnetic field was modulated at frequency $\omega/2\pi = 6.1$ Hz with amplitude 40 mT. Standard He4 pumping and heating were used to set the temperature in the range 1.7-13K. We present the results obtained from the R sample with peak mobility 3.4m$^2$/Vs and H sample with 3.3m$^2$/Vs at 1.7K, measured most extensively; the data for several other R and H samples were very similar. An example of the data collected at different temperatures and electron densities is shown in Fig. 1. At low field, $\partial m/\partial n(B)$ grows linearly below some temperature-dependent field $B^\ast$ indicated by the gray area in Fig. 1 above which the slope changes. The maximal $\partial m/\partial n(B)$ at lowest achievable density and temperature reaches 0.9$\mu_B$ at $B \approx 0.3$ T, indicating almost full spin polarization.

Let us first consider the low-field slope of the data, $\partial^2 m/\partial n \partial B|_{B=0} = \partial \chi/\partial n|_{B=0}$, plotted in Fig. 2 as a function of density at different temperatures. The slope is large and positive at low temperature and density. As the density increases, the slope decreases, changes sign, and finally almost tends to zero at highest densities. As a function of temperature the slope decreases rapidly almost to zero at $T = 13K$. There is some small negative slope left at highest temperatures in the R samples, which we attribute to the diamagnetic contribution due to the finite thickness of the 2DEG [2]. This negative slope is absent for H samples [11].

In order to obtain the zero field susceptibility we integrate $\partial \chi/\partial n$:

$$\chi(n, T) = \int_0^n \partial \chi/\partial n(n', T) dn'$$

We were able to measure $\partial \chi/\partial n$ down to densities $0 < n_L(T) < n_c$, so we must extrapolate the integrand in (2) down to zero density. For simplicity we took $\partial \chi/\partial n$ to be constant and equal to its value at $n_L(T)$. In order to get the spin part of susceptibility we need to subtract the temperature-independent diamagnetic susceptibility. This we did by equating the high temperature susceptibility to the non-renormalized Pauli one. The results of the integration under these assumptions are depicted in Fig. 2b, together with the subtracted diamagnetic contribution.

Before discussing the results, let us first examine possible sources of inaccuracy. The diamagnetic contribution is small, and therefore can affect only the high-temperature susceptibility. The largest error may come from the extrapolation of $\partial \chi/\partial n$ to zero density, as well from the uncertainty $\approx 10^{10}$cm$^{-2}$ in the density itself [11]. We note that even if we make the extreme assumption that the magnetization is zero below $n_L(T)$, thus significantly underestimating the susceptibility, we would get qualitatively similar results, as also shown in Fig. 2b.

The most striking features of the susceptibility $\chi$ are (i) its low-temperature value, at maxima exceeding the Pauli one by the factor of 40, and (ii) strong temperature and relatively weak density dependence at high densities. The temperature dependence of the susceptibility per electron for several densities is shown in Fig. 3. It can be fit reasonably well with a power law $\chi \propto 1/T^\alpha$, with $\alpha \approx 2.4 \pm 0.2$. Not only does the susceptibility diverge faster than the independent-spin Curie susceptibility, $\mu_B^2/T$, but its low-temperature low-density value exceeds the Curie value, indicated by the solid line in Fig. 3.

We emphasize that neither the susceptibility value, nor its temperature dependence change qualitatively when the system passes through the metal-insulator transition at $n_c \approx 8.5 \cdot 10^{10}$cm$^{-2}$. Remarkably, the results for the R and H samples, shown in Figs. 2a and 3 by the bold and hollow...
ceptibility solely in the low field domain. As seen in
the temperature, whereas we observe the divergent sus-
ceptibility to be determined by excitations in the vicin-
ity of the Fermi level. Interaction corrections to sus-
ceptibility in a clean system were calculated in [14, 15].

Fig. 1a, \( m_1 \)

\( \partial \chi / \partial n \)

\( \propto \)

\( \propto \)

\( B / 2T \)

\( 7K \)

\( -1.7K \)

\( -2.4K \)

\( -4K \)

\( 10^{11} \text{cm}^{-2} \)

\( 10^{11} \text{cm}^{-2} \)

\( T = 1.7K, m, n \) in Fig. 4

\( 1/T^n \)

\( \partial \chi / \partial n \) also diverges as \( 1/T^n \), even deep in the metallic regime, as indicated in Fig. 3.

The magnetic moment can be obtained similarly to the susceptibility, by integration of \( \partial m(n, T, B) / \partial n \) over \( n \) for given values of \( B \) and \( T \). Since at low-temperature \( \partial m / \partial n \) saturates, and even drops for \( B > B^* \), as seen in Fig. 4, \( m(B) \) behaves similarly; see the inset in Fig. 4.

This saturation resolves the apparent contradiction between the current thermodynamic and earlier transport susceptibility measurements. Indeed, all the transport measurements to date were performed in the high-field domain, when the Zeeman energy is larger than the temperature, whereas we observe the divergent susceptibility solely in the low field domain. As seen in

\( B^* \) depends linearly on temperature. Additionally, we note that \( B^* \) appears to be density-independent for \( n \ll 3 \cdot 10^{13} \text{cm}^{-2} \). The low-field domain, at which the zero-field susceptibility can be determined, lies in the bottom-right corner. In contrast, both Shubnikov-de Haas, and magnetoresistance data belong to the top left corner. If one takes the magnetic moment at e.g. 2T, which is the typical total field in [4, 5, 9], he would get the value of the susceptibility, defined as \( \chi(B) = m(B)/B \), similar to obtained in [4, 9], as shown in Fig. 2.

In our previous work [7] we could only measure \( \partial m / \partial n \) down to \( n \ll n_c \), and therefore we chose to integrate it from a high density at which we assumed \( m \) to be known and temperature independent. We used the susceptibility obtained from the Shubnikov-de Haas measurements in the high-field domain, (similar assumption was later used in [13]) as the initial value for the integration in both high and low field domains. As the present work shows the temperature dependence of the susceptibility in the low field domain cannot be ignored.

A susceptibility diverging faster than \( 1/T \) indicates a ferromagnetic interaction between spins. If, however, the low temperature extrapolation of \( B^* = k_B T / 5.7 \mu_B + 0.03T \) in Fig. 4 is taken seriously, it means that the divergence should be cut off at about 100mK. A divergent susceptibility contradicts numerical calculations [12] for a two-valley system, which predicts no tendency to ferromagnetism. At high enough density one would expect the susceptibility to be determined by excitations in the vicinity of the Fermi level. Interaction corrections to susceptibility in a clean system were calculated in [14, 15].

Although the results differ in prefactors, both predict a correction to the susceptibility of the order of \( T / E_F \). It is important to stress that not only does \( \chi \) diverge

FIG. 2: (a) Derivative of the low-field susceptibility with re-
spect to the density \( \partial \chi / \partial n \) for different temperatures. Hollow symbols-R sample, bold symbols - H sample. (b) \( \chi \) obtained by integration; Dotted line (also in c) - non-renormalized Pauli susceptibility, dashed line—subtracted diamagnetic con-
tribution. × - the same as □ integrated from the density \( n_L \). (c) Susceptibility defined as \( \chi = m/B \) at 2T. Solid line - the Shubnikov-de Haas susceptibility [2]. The temperatures are indicated by the symbols: □ - 1.7K; ○ - 2.4K; △ - 4K; ◆ - 7K; × - 13K. Inset - maximum of the magnetization in the field \( 0 < B < 2T \) at \( T = 1.7K \). The symbols respectively, are very similar. The largest contribution to the integral [2], which determines \( \chi \), even at high densities comes from \( \partial \chi / \partial n \) at low densities. Therefore it is important to stress that \( \partial \chi / \partial n \) also diverges as \( 1/T^n \), even deep in the metallic regime, as indicated in Fig. 3.

FIG. 3: Low-field susceptibility \( \chi \) per electron vs. temper-
ature for different densities, indicated in units of \( 10^{11} \text{cm}^{-2} \). Hollow symbols-R sample, bold symbols - H sample. Also \( |\partial \chi / \partial n| \) is shown for comparison with × and +. Dashed line line-\( \chi \propto 1/T^2 \), solid line-Curie susceptibility.

Although the results differ in prefactors, both predict a correction to the susceptibility of the order of \( T / E_F \). It is important to stress that not only does \( \chi \) diverge...
as $1/T^2$, but so does $\partial \chi / \partial n$, even deep in the metallic regime at densities $n = 2 - 4 \times 10^{11} \text{cm}^{-2}$, in clear contradiction with [13] [13]. We therefore think the divergent susceptibility is disorder-related and signals the presence of local moments in the 2D system up to high densities; numerical calculations indeed predict susceptibility enhancement by disorder [10]. This suggestion does not contradict numerous existing transport data, since localized states are very rarely seen in transport, which probe the electrons on the time scale of ps. In contrast, thermodynamic measurements probe all the states which can be recharged on the time scale set by the magnetic field modulation. On the other hand, it is important to emphasize that the maximal magnetization of the system, shown in the inset in Fig. 2, is very close to $\mu_0 B$ at low densities. This rules out the contribution of distant magnetic centers, existing in addition to the 2DEG (such as, e.g., localized states in the bulk Si, or deep levels at the interface).

A clue about the origin of the anomalous susceptibility may come from the magnetization drop at low temperatures in the intermediate magnetic field, above $B^*$ and below some 2.5T, visible in the inset in Fig. 4; the magnetization grows again at even higher fields. For a magnetic field coupled only to the spins, the magnetization should grow with the field. Therefore this drop must be related either to the orbital effects, or to the field induced variation in the number of electrons contributing to the signal.

In conclusion, we observed a strongly enhanced divergent susceptibility at low temperatures in high-mobility MOSFET’s. We find no qualitative change of the susceptibility behavior in the vicinity of the metal-insulator transition density $n_c$, and the divergence persists deep into the high density metallic phase, similar to the observations on phosphorus doped Si.

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FIG. 4: Comparison with transport measurements. ■ marks almost density-independent $B^*$. Arrows indicate the domains of the previous measurements of $\chi$. Inset: field dependence of magnetization for R sample at 1.7K.

References:

[1] E. C. Stoner, Ferromagnetism, Rep. Prog. Phys. 11 43 (1947)
[2] R. N. Bhatt, Physica Scripta. T14 7-16, (1986).
[3] D. Simonian, S. V. Kravchenko, M. P. Sarachik, and V. M. Pudalov, Phys. Rev. B 55, R13 421 (1997); 57, R9420 (1998); Phys. Rev. Lett. 79, 2304 (1997).
[4] A. A. Shashkin, S. V. Kravchenko, V. T. Dolgopolov, T. M. Klapwijk, Phys. Rev. Lett., 87, 86801 (2001)
[5] S. A. Vitkalov, H. Zheng, K. M. Mertes, M. P. Sarachik, T. M. Klapwijk, Phys. Rev. Lett. 87, 086401 (2001).
[6] $n_c$ is the density below which the metallic temperature dependence of resistivity switches to the insulating one.
[7] O. Prus, Y. Yaish, M. Reznikov, U. Sivan, and V.M. Pudalov, Phys. Rev. B 67, 205407 (2003).
[8] V. M. Pudalov, M. E. Gershenson, and H. Kojima, cond-mat/0110160
[9] V. M. Pudalov, M. E. Gershenson, H. Kojima, N. Butch, E. M. Dizhur, G. Brunthaler, A. Prinz, and G. Bauer, Phys. Rev. Lett. 88, 196404 (2002).
[10] p-type substrate doping estimated from C-V measurements is $2 \times 10^{15} \text{cm}^{-3}$ for R samples, and $3 \times 10^{16} \text{cm}^{-3}$ for the H ones. Higher doping means thinner 2-D gas, and therefore smaller diamagnetic contribution. Indeed, in [13] diamagnetic contribution was not observed.
[11] Since the magnetic field was in the sample plain, we could not measure the Hall simultaneously, and zero density was determined from comparison between the Hall and capacitance measurements in a different run.
[12] M. Marchi et. al. PRB 80, 035103 (2009); G. Fleur, X. Waintal, arXiv: 0902.3171v1 (2009).
[13] A. A. Shashkin, S. Anissimova, M. R. Sakr, S. V. Kravchenko, V. T. Dolgopolov, and T. M. Klapwijk, Phys. Rev. Lett. 96, 036403 (2006).
[14] A. Shekhter and A. M. Finkelstein, Phys. Rev. B 74, 205122 (2006)
[15] A. V. Chubukov, D. L.Maslov Phys. Rev. B 68, 155113 (2003); A. V. Chubukov, D. L. Maslov, S. Gangadharaiah, and L. I. Glazman, Phys. Rev. Lett. 95, 026402 (2005); D. L.Maslov A. V. Chubukov, Phys. Rev. B 79, 075112 (2009).
[16] S. De Palo, S. Moroni, G. Senatore J. Phys. A: Math. Theor. 42, 214043, (2009).