Interaction Corrections to Two-Dimensional Hole Transport in Large $r_s$ Limit

H. Noh,¹ M. P. Lilly,² D. C. Tsui,¹ J. A. Simmons,² E. H. Hwang,³ S. Das Sarma,³ L. N. Pfeiffer,⁴ and K. W. West⁴

¹Department of Electrical Engineering, Princeton University, Princeton, New Jersey 08544
²Sandia National Laboratories, Albuquerque, New Mexico 87185
³Condensed Matter Theory Center, Department of Physics, University of Maryland, College Park, Maryland 20742
⁴Bell Labs, Lucent Technologies, Murray Hill, New Jersey 07974

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The metallic conductivity of dilute two-dimensional holes in a GaAs HIGFET (Heterojunction Insulated-Gate Field-Effect Transistor) with extremely high mobility and large $r_s$ is found to have a linear dependence on temperature, consistent with the theory of interaction corrections in the ballistic regime. Phonon scattering contributions are negligible in the temperature range of our interest, allowing comparison between our measured data and theory without any phonon subtraction. The magnitude of the Fermi liquid interaction parameter $F_0$ determined from the experiment, however, decreases with increasing $r_s$ for $r_s \gtrsim 22$, a behavior unexpected from existing theoretical calculations valid for small $r_s$.

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In two-dimensional (2D) charge carrier systems, it is well known that any amount of disorder in the absence of interactions between the carriers will localize the carriers, leading to an insulator with zero conductivity ($\sigma$) as the temperature ($T$) is decreased to zero. Recent experiments on high mobility dilute 2D systems, on the other hand, have shown a “metallic” behavior at low $T$, characterized by an increasing $\sigma$ with decreasing $T$, and an apparent metal-insulator transition (MIT) as the carrier density is lowered. There are three important energy scales in these systems. The first two are the Fermi energy and the interaction energy. Their ratio, which is $r_s$, is around 10 or higher for the systems where the MIT is observed, implying that interaction must be playing a role. The other energy scale is related to the disorder in the system given by $\hbar/\tau$, where $\tau$ is the elastic scattering time. It has been found from more recent experiments that disorder is also playing a significant role. In particular, the critical density ($n_c$ for electrons and $p_c$ for holes), above which a system shows the metallic behavior, is found to decrease when disorder in the 2D system is decreased.

An important question is whether this apparent metallic state is truly a new ground state of the 2D charge carriers or simply a novel finite temperature behavior of the 2D gas, since all experiments are done at finite $T$. What is measured in such experiment is the temperature coefficient, $d\sigma/dT$. The metallic behavior evinced by the observation of negative $d\sigma/dT$ at finite $T$ does not necessarily mean, however, a true metal with nonzero conductivity at $T = 0$. Recently, Zala et al. calculated the Fermi liquid interaction corrections to the conductivity in the asymptotic low temperature regime ($T/T_F \ll 1$ where $T_F$ is the Fermi temperature), and pointed out that the metallic behavior seen in the high mobility samples could be understood by taking into account of interaction corrections in the “high temperature” ballistic regime ($k_B T \gg \hbar/\tau$). They found that the conductivity of interacting 2D carriers changes linearly with $T$ in the ballistic regime, $T_F \gg T \gg \hbar/k_B \tau$, with the sign as well as the magnitude of $d\sigma/dT$ depending on the strength of the interaction, while in the low temperature diffusive regime ($k_B T \ll \hbar/\tau$) the conventional logarithmically changing conductivity is recovered. A linear dependence of $\sigma$ on $T$ has also been predicted in earlier theories based on temperature dependent screening, and this screening contribution is included in the theory by Zala et al.

Experimentally, however, it is not straightforward to identify the interaction corrections unequivocally in the ballistic regime for two main reasons. First, scattering by phonons can give significant contributions at high temperatures. In order to have the ballistic regime at sufficiently low $T$ to minimize the phonon contributions, the 2D charge carrier system must have a very high mobility so that $\hbar/k_B \tau$ becomes very low. Second, the temperature constraint, $T_F \gg T \gg \hbar/k_B \tau$, satisfying the dual conditions of being in the ballistic regime (i.e. $T \ll \hbar/k_B \tau$, which is a high-temperature constraint) and of also being in the asymptotic low temperature regime of $T \ll T_F$ (so that the thermal expansion in $T/T_F$, essential for obtaining the linear-T term in the conductivity, applies) is not easily satisfied experimentally, and indeed most experimentally measured $\rho(T)$ data in 2D systems do not manifest any clear cut linear-T behavior at low temperatures. An additional issue we are addressing in this work is whether the theory of interaction corrections to the conductivity can describe 2D transport in the large $r_s$ limit as well. This is particularly germane in view of the fact that the interaction theory is a systematic many-body diagrammatic expansion in the interaction parameter $r_s$ (albeit an infinite order formal expansion), and the question of the radius of convergence of such an $r_s$-expansion becomes quite important for large $r_s$ values obtained in our samples.

In this paper, we report our experiments on the low
temperature conductivity and the in-plane magnetoresistance (MR) of two-dimensional (2D) holes with extremely high mobility and very low density ($r_s \approx 17$ to 80) to study the interaction corrections. From the temperature dependence of conductivity, we clearly observed a temperature region where the the conductivity shows a linear dependence on $T$ even in this large $r_s$ limit for a range of densities in the metallic side of the transition. However, the Fermi liquid interaction parameter $F_0$, determined from a comparison of the data with the theory by Zala et al., shows a surprising non-monotonic dependence on the carrier density with its value lying between $-0.5$ and $-0.7$. $F_0$ increases in magnitude with decreasing density for $p \gtrsim 2 \times 10^{10}$ cm$^{-2}$ ($r_s \lesssim 22$) and then decreases with decreasing density for $p \lesssim 2 \times 10^{10}$ cm$^{-2}$ ($r_s \gtrsim 22$), a behavior unexpected from a simple extrapolation of the predicted dependence of $F_0$ on small $r_s$. A separate measurement of effective g-factor ($g^*$) from the in-plane MR provides a further confirmation of the unexpected behavior of $F_0$. $g^*$ decreases with decreasing density, consistent with the behavior expected from $F_0$ for $p \lesssim 2 \times 10^{10}$ cm$^{-2}$ through $g^* = g_0 / (1 + F_0^2)$. The sample used in this study is a heterojunction insulated-gate field-effect transistor (HIGFET) made on a (100) surface of GaAs. A metallic gate, separated by an insulator (AlGaAs) from the semiconducting GaAs, is used to induce the 2D holes at the interface between the GaAs and AlGaAs. Ohmic contacts to the 2D holes are made by using a self-aligned contact technique which allows the diffusion of the contact material under the gate region. We would like to emphasize that there is no intentional doping in the sample and the 2D holes are induced by phonon contributions are negligible. The extremely high mobility also allows us to measure the temperature dependence of conductivity down to very low densities reaching $p = 1.5 \times 10^9$ cm$^{-2}$, with $r_s$ near 80 (assuming a hole mass of $0.38m_e$).

In Fig. 1, we show the temperature dependence of the resistivity ($\rho$) at various densities. For $p \geq 1.7 \times 10^{10}$ cm$^{-2}$, $\rho$ decreases monotonically with decreasing $T$, showing a metallic behavior. For $p$ between 1.2 and $0.4 \times 10^{10}$ cm$^{-2}$ $\rho$ shows a nonmonotonic dependence on $T$. It initially increases with decreasing $T$ at high $T$, which was interpreted as the classical to quantum crossover, and then decreases with decreasing $T$ at low $T$ showing a metallic behavior as the system goes into the degenerate regime. This crossover shifts to lower temperature with decreasing density and the range where the metallic behavior is seen becomes very narrow, especially for $p = 0.5$ and $0.4 \times 10^{10}$ cm$^{-2}$. For $p \leq 3 \times 10^9$ cm$^{-2}$, $\rho$ increases monotonically with decreasing $T$, exhibiting an insulating behavior. To identify the critical density, we measured $\rho$ as a function of $p$ at different temperatures and the data is shown in the inset. The critical density determined from the crossing point, which shows a temperature independent resistivity, is $p_c = 3 \times 10^9$ cm$^{-2}$. This low critical density, the lowest ever observed in 2D systems which exhibit the MIT, is consistent with the previous observation by Yoon et al. that the critical density becomes lower with decreasing disorder in the system. If we use the hole effective mass $m^* = 0.38m_e$, this critical density corresponds to $r_s = 57$, which is much larger than the $r_s = 37$ predicted for the Wigner crystallization in 2D.

In Fig. 2(a), we replotted the data for $p = 3.2 \times 10^{10}$ cm$^{-2}$ to $0.7 \times 10^{10}$ cm$^{-2}$ as $\sigma$ vs $T$. The data are scaled by $\sigma_0$, the value of $\sigma$ extrapolated to $T = 0$, and offset by 0.1 for clarity. The metallic behavior is identified by increasing $\sigma$ with decreasing $T$ at low $T$ for all these densities. Clearly, there is a region where $\sigma$ shows a linear dependence on $T$ as shown by the best fits in the figure with solid lines. We note that $r_s$ for these densities ranges from 17 to 37 and $\sigma$ shows a linear dependence on $T$ for such large $r_s$. To compare our results with the theory by Zala et al., we need to consider several points. First, as discussed below phonon contributions are negligible in the temperature range where the linear dependence is
observed, which is below 200 mK for \( p = 3.2 \times 10^{10} \text{ cm}^{-2} \) and becomes lower for lower densities. Second, this linear region is in the ballistic regime \( (k_B T > \hbar / \tau) \) since \( \hbar / k_B \tau \) calculated from \( \sigma_0 \) ranges from 16 mK for \( p = 3.2 \times 10^{10} \text{ cm}^{-2} \) to 80 mK for \( p = 0.7 \times 10^{10} \text{ cm}^{-2} \). Finally, this region is also much lower than the Fermi temperature \( (T_F) \) of the system, which is 500 mK for the lowest density \( p = 0.7 \times 10^{10} \text{ cm}^{-2} \). Zala et al. have pointed out that the regime where \( F_0^* \) can be treated as a momentum independent constant is \( T \ll (1 + F_0^*)^2 T_F \). A self-consistency check after we have determined \( F_0^* \) approximately (but somewhat weakly) satisfies this condition. All these allow a direct comparison of our data with their theory. From their theory, the slope of this linear dependence is directly related to the Fermi liquid interaction parameter \( F_0^* \) by the relation

\[
\text{Slope} = \frac{m^* k_B}{\pi \hbar^2 p} \left[ 1 + \frac{3F_0^*}{(1 + F_0^*)} \right].
\]

Using \( m^* = 0.38 m_e \) (which was obtained from cyclotron resonance of high density \( 5 \times 10^{14} \text{ cm}^{-2} \) 2D holes on (100) surface of GaAs\ROLE{25}{[25]}, we have obtained from our data \( F_0^* \) as a function of \( p \). The result is shown in Fig. 2(b). The value of \( F_0^* \) lies between 0.5 and 0.7 in the density range we measured from \( 3.2 \times 10^{10} \text{ cm}^{-2} \) to \( 0.7 \times 10^{10} \text{ cm}^{-2} \). For \( p \gtrsim 2 \times 10^{10} \text{ cm}^{-2} \) \( (r_s \lesssim 22) \), the magnitude increases with decreasing density. Proskuryakov et al.\ROLE{26}{[26]} have also found that \( F_0^* \) for 2D holes in (311)A GaAs/AlGaAs heterostructure increases in magnitude with decreasing \( p \) for \( p = 2 \) to \( 8 \times 10^{10} \text{ cm}^{-2} \) with values between \(-0.3 \) and \(-0.45 \). The change of \( F_0^* \) per density is similar in both experiments, while the magnitude of \( F_0^* \) in our measurement is much larger. We note that \( F_0^* \) found from experiments on 2D electrons in Si-MOSFET\ROLE{15}{[15]} also has much smaller magnitude, ranging from \(-0.14 \) to \(-0.5 \) for electron densities 1 to \( 4 \times 10^{11} \text{ cm}^{-2} \) while the value found from p-SiGe by Coleridge et al.\ROLE{14}{[14]} is somewhat comparable to ours, between \(-0.55 \) and \(-0.65 \). What is surprising in our experiment, which explores the much lower density \( (r_s \sim 22) \) regime, is that the magnitude of \( F_0^* \) does not increase monotonically with decreasing density. When the density is decreased below \( 2 \times 10^{10} \text{ cm}^{-2} \) \( (r_s \lesssim 22) \) the magnitude of \( F_0^* \) decreases again. This is opposite to the predicted dependence of \( F_0^* \) on \( r_s \) valid for small \( r_s \). To our best knowledge, the dependence of \( F_0^* \) on \( r_s \) when \( r_s \) is large has not been calculated theoretically, and it is not possible to compare our result with any theoretical predictions at this time.

Experimentally, however, an additional test for this unexpected behavior of \( F_0^* \) can be made from the MR measurements under an in-plane magnetic field \( (B_y) \). The in-plane MR provides a way to measure the effective g-factor, which is directly related to \( F_0^* \) by the relation \( g^* / g_0 = 1/(1 + F_0^*) \), where \( g_0 \) is the bare g-factor. Figure 2(a) shows the in-plane MR measured in our sample for various densities at 65 mK. The MR increases as \( \exp(B_y^*) \) at low \( B_y \) and \( \exp(B_y) \) at high \( B_y \), consistent with an earlier observation for 2D holes on (311)A GaAs.\ROLE{14}{[14]} Similarly strong MR has also been observed for 2D electrons in Si-MOSFET\ROLE{15}{[15]} and the 2D electrons in GaAs.\ROLE{15}{[15]} It has been established that this crossover from low field to high field dependences corresponds to full spin polarization of the carriers and its position allows the determination of \( g^* \). The crossover field \( B^* \), determined from the position where the second derivative of the \( p \) vs \( B_y \) curve becomes maximum, is marked by an arrow for \( p = 3.2 \times 10^{10} \text{ cm}^{-2} \) in Fig. 3(a), and the dependence of \( B^* \) on \( p \) is shown in Fig. 3(b). For \( p > p_c \), \( B^* \) decreases linearly, with decreasing \( p \) (best fit given by the solid line), and saturates in the insulating side of the MIT for \( p < p_c \). This behavior is also consistent with earlier observation by Yoon et al.\ROLE{14}{[14]} for the 2D holes on (311)A. A different way of determining \( B^* \), using the inflection point between high and low field dependences, yields a result within the error bar of this plot, and produces an error of less than 15 % in \( g^* \).

For the metallic side, \( g^* \) determined from the relation \( 2E_F = g^* \mu_B B^* \) (where \( E_F \) is the Fermi energy and \( \mu_B \) is the Bohr magneton) is shown in Fig. 3(c) as a function of \( p \). \( g^* \) decreases monotonically with decreasing \( p \) for the density range measured. The solid line in Fig. 3(c) is the result when the best linear fit in Fig. 3(b) is used. Although a quantitative comparison between \( g^* \) in Fig. 3(c) and the value of \( g^* \) expected from \( F_0^* \) cannot be made since the bare g-factor \( g_0 \) is not well known for holes in GaAs, we can qualitatively compare their density de-
pends. From the density dependence of \( F_0^g \) shown in Fig. 3(b) and the relation \( g^*/g_b = 1/(1 + F_0^g) \), we expect that \( g^* \) increases with decreasing \( p \) for \( p \geq 2 \times 10^{10} \text{ cm}^{-2} \) and then decreases with increasing \( p \) for \( p < 2 \times 10^{10} \text{ cm}^{-2} \). The behavior of \( g^* \) shown in Fig. 3(c), therefore, agrees well with that expected from \( F_0^g \) for \( p < 2 \times 10^{10} \text{ cm}^{-2} \), where \( F_0^g \) decreases in magnitude with decreasing \( p \). While this region of \( p \) is where we have observed the unexpected density dependence of \( F_0^g \), in this region there is good agreement in the behavior of \( F_0^g \) and \( g^* \). Thus, our in-plane MR measurements confirm the unexpected behavior of \( F_0^g \) found from the \( T \) dependence of \( \sigma \).

The decrease in magnitude of \( F_0^g \) with decreasing \( p \) (increasing \( r_s \)) is of great interest, and needs further examination. We, therefore, have also analyzed our data without the assumption of a density independent mass \( m^* = 0.38m_e \). In this analysis, we used three independent relations (equation (1), \( 2E_K = g^*\mu_B B^* \), and \( g^*/g_b = 1/(1 + F_0^g) \)) between \( F_0^g \), \( m^* \), and \( g^* \) to calculate each quantity. Since \( g_b \) is not well known, we used it as a parameter, and found that the density dependence of each quantity does not depend on a specific value of \( g_b \) while \( g_b = 0.5 \) gives the best quantitative agreement with \( F_0^g \) and \( g^* \) determined earlier. From this analysis, \( F_0^g \) also exhibits a nonmonotonic dependence on \( p \). The magnitude of \( F_0^g \) increases with decreasing \( p \) for \( p \gtrsim 1.5 \times 10^{10} \text{ cm}^{-2} \) and decreases with decreasing \( p \) for \( p \lesssim 1.5 \times 10^{10} \text{ cm}^{-2} \). The decrease in magnitude of \( p \) for large \( r_s \) is still observed and confirmed once again.

Any explanation for this surprising result should take into account the large \( r_s \) values in our system. For \( r_s \gtrsim 37 \), the 2D system is expected to be a pinned Wigner crystal. The critical density for MIT in our system corresponds to \( r_s = 57 \), considerably larger than the critical \( r_s \) predicted for this crystallization. The \( r_s \) values for which we observed the unexpected behavior of \( F_0^g \) range between 22 and 37, close to that predicted for Wigner crystallization. We note that there has been a Monte Carlo calculation of Fermi-liquid parameters for \( r_s \) up to 5, where \( F_0^g \) still increases monotonically in magnitude with increasing \( r_s \). To our knowledge, there is no theoretical calculation of the dependence of \( F_0^g \) on \( r_s \) when \( r_s \) is large, relevant to our experiment. The question whether the crystallization is preceded by a ferromagnetic instability with \( F_0^g = -1 \) has to be addressed as well. From Monte Carlo calculations, Tanatar and Ceperley have showed both possibilities of a diverging and a finite-valued spin susceptibility as \( r_s \) increases toward the critical \( r_s \) for the Wigner crystallization. Our result appears to imply that the ferromagnetic instability does not occur in the large \( r_s \) regime of our 2D hole system.

We now address the important issue of phonon scattering contribution to our measured hole resistivity, which we have ignored in our analysis. The question of phonon contribution to the resistivity is crucial since, if it is significant, it would then be impossible to compare our measured resistivity to the interaction theory. We have therefore theoretically directly calculated the phonon scattering contribution by including both deformation potential and piezoelectric coupling of the 2D holes to GaAs acoustic phonons. Following ref. 13 we have carried out a detailed calculation of the phonon scattering contribution to the hole resistivity in the parameter range of our experiment with the results being shown in Fig. 4. Our theoretical phonon-only resistivity, as shown in Fig. 4, demonstrates that for \( T \leq 200 \text{ mK} \), the temperature range we concentrate on in comparing our experimental resistivity with the interaction theory (see Fig. 3), the phonon contribution to the resistivity is miniscule (less than 1% of the measured resistivity). We are therefore justified in neglecting phonon scattering effects in the discussion of our experimental results as long as we restrict ourselves to \( T < 200 \text{ mK} \) as we have done in analyzing our data. As is obvious from our theoretical results pre-

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**FIG. 3:** (a) \( \rho \) vs \( B_0 \) at \( T = 65 \text{ mK} \) and \( p = 3.2, 2.2, 1.7, 1.2, 0.9, 0.7, 0.5, 0.4, 0.27, 0.22, \) and \( 0.18 \times 10^{10} \text{ cm}^{-2} \) from the bottom. \( B^* \) is marked by an arrow for \( p = 3.2 \times 10^{10} \text{ cm}^{-2} \). (b) \( B^* \) vs \( p \). Solid line is a linear fit to the data for metallic side, \( p > p_c \). (c) \( g^* \) determined from \( B^* \). Solid line is the result obtained by the linear fit in (b).
sent in Fig. 4, hole-phonon scattering contribution to hole resistivity becomes non-negligible for $T > 200 mK$ and is, in fact, significant for $T \geq 500 mK$ with its quantitative importance increasing with decreasing carrier density.

We should point out in this context that we disagree with the methodology employed recently by Proskuryakov et al. in subtracting out a calculated phonon scattering contribution to their measured resistivity in analyzing their hole transport data in the context of a quantitative comparison with the interaction theory of Zala et al. First, it is well-known that Matthiessen’s rule does not apply to 2D systems at finite temperatures, and therefore, subtraction of phonon contribution (even if this contribution were accurately known, which is questionable) in order to obtain the nonphonon part is unjustified and may be subject to large errors. (This problem is worse in the presence of screening of hole-phonon interaction, which must be included in the theory.) Second, the calculation of phonon scattering contribution to hole resistivity, following ref. 21, carried out by Proskuryakov et al. is rather crude and approximate (compared, for example, with our theoretical calculations shown in Fig. 3 of this paper). We note that in ref. 11 the measured hole resistivity (before any phonon subtraction) hardly manifests any clear cut linear temperature regime, and the subtracted phonon contribution is a large fraction of the measured resistivity, thereby casting substantial doubt on the accuracy of the subtracted resistivity eventually compared with the interaction theory. Our analysis in this work avoids these serious pitfalls of ref. 11 by directly considering the measured resistivity in the context of interaction theory, which we justify by explicitly calculating the phonon contribution to the hole resistivity in the temperature range of our interest and showing it to be negligible so that no arbitrary and unjustifiable phonon subtraction is required (in contrast to ref. 11).

We calculate the temperature dependence of the hole resistivity by considering screened acoustic-phonon scattering. We include both deformation potential and piezoelectric coupling of the 2D holes to 3D acoustic phonons of GaAs. Details of the acoustic-phonon scattering theory are given in Ref. 19. In this calculation we use the parameters corresponding to GaAs: $c_{1} = 5.14 \times 10^{5}$ cm/s, $c_{4} = 3.04 \times 10^{5}$ cm/s, $\rho = 5.3 g/cm^{3}$, $e_{h14} = 1.2 \times 10^{7}$ eV/cm, and $D = -8.0$ eV. In the low temperature range ($T < 0.2 K$) we find $\rho(T) \propto T^{3} - T^{7}$ because the deformation potential scattering dominates over piezoelectric coupling. In the intermediate temperature range ($0.2 K < T < T_{BG}$) we have $\rho(T) \propto T^{3}$ mostly due to the piezoelectric scattering. Above the Bloch-Grüneisen temperature, $T_{BG} = 2k_{F}c_{1}/k_{B} \approx 1 - 2K$, both scattering processes give rise to linear temperature dependence of the resistivity, $\rho(T) \propto T$. The phonon contribution to the resistivity shows very weak temperature dependence and is negligible when the temperature is substantially below $T_{BG}$. We emphasize that the phonon contribution to the resistivity cannot be linear for $T < T_{BG}$, and for $T \leq 200 mK$ the phonon contribution is negligible.

In summary, we have measured the temperature dependence of the metallic conductivity of extremely high mobility dilute 2D holes in GaAs in large $r_s$ limit. We find that the conductivity shows a linear dependence on temperature in the ballistic regime. The Fermi liquid interaction parameter $F_0$ obtained from our data is found to exhibit a nonmonotonic dependence on density and decrease in magnitude with increasing $r_s$ for $r_s \gtrsim 22$.

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1. E. Abrahams, P. W. Anderson, D. C. Licciardello, and T. V. Ramakrishnan, Phys. Rev. Lett. 42, 673 (1979).
2. For a review, see B. L. Altshuler, D. L. Maslov, and V. M. Pudalov, Physica E (Amsterdam) 9(2), 209 (2001); E. Abrahams, S. V. Kravchenko, M. P. Sarachik, Rev. Mod. Phys. 73, 251 (2001).
3. J. Yoon, C. C. Li, D. Shahar, D. C. Tsui, and M. Shayegan, Phys. Rev. Lett. 82, 1744 (1999).
4. G. Zala, B. N. Narozhny, and I. L. Aleiner, Phys. Rev. B 64, 214204 (2001).
5. B. L. Altshuler, A. G. Aronov, and P. A. Lee, Phys. Rev. Lett. 44, 1288 (1980); H. Fukuyama, J. Phys. Soc. Jpn. 48, 2169 (1980); A. M. Finkelstein, Sov. Phys. JETP 57, 97 (1983) [Zh. Eksp. Teor. Fiz. 84, 168 (1983)].
A. Gold and V. T. Dolgopolov, Phys. Rev. B 33, 1076 (1986); S. Das Sarma, Phys. Rev. B 33, 5401 (1986).

S. Das Sarma and E. H. Hwang, Phys. Rev. Lett. 83, 164 (1999); Phys. Rev. B 61, R7838 (2000).

B. E. Kane et al., Appl. Phys. Lett. 63, 2132 (1993).

B. Tanatar and D. M. Ceperley, Phys. Rev. B 39, 5005 (1989).

S. Das Sarma and E. H. Hwang, Phys. Rev. Lett. 83, 164 (1999); Phys. Rev. B 61, R7838 (2000).

B. E. Kane et al., Appl. Phys. Lett. 63, 2132 (1993).

B. Tanatar and D. M. Ceperley, Phys. Rev. B 39, 5005 (1989).

H. L. Stormer et al., Phys. Rev. Lett. 51, 126 (1983).

Y. Y. Proskuryakov, A. K. Savchenko, S. S. Safonov, M. Pepper, M. Y. Simmons, and D. A. Ritchie, Phys. Rev. Lett. 89, 076406 (2002).

A. A. Shashkin, S. V. Kravchenko, V. T. Dolgopolov, and T. M. Klapwijk, cond-mat/0111478; S. A. Vitkalov, K. James, B. N. Narozhnny, M. P. Sarachik, and T. M. Klapwijk, cond-mat/0204566; V. M. Pudalov, M. E. Gershenson, H. Kojima, G. Brunthaler, A. Prinz, and G. Bauer, cond-mat/0205449.

P. T. Coleridge, A. S. Sachrajda, and P. Zawadzki, Phys. Rev. B 65, 125328 (2002).

J. Yoon, C. C. Li, D. Shahar, D. C. Tsui, and M. Shayegan, Phys. Rev. Lett. 84, 4421 (2000).

D. Simonian, S. V. Kravchenko, M. P. Sarachik, and V. M. Pudalov, Phys. Rev. Lett. 79, 2304 (1997).

E. Tutuc, S. Melinte, and M. Shayegan, Phys. Rev. Lett. 88, 036805 (2002).

T. Okamoto et al., Phys. Rev. Lett. 82, 3875 (1999); S. A. Vitkalov et al., Phys. Rev. Lett. 85, 2164 (2000); V. T. Dolgopolov and A. Gold, JETP Lett. 71, 27 (2000); I. F. Herbut, Phys. Rev. B 63, 113102 (2001); E. Tutuc et al., Phys. Rev. Lett. 86, 2858 (2001).

Y. Kwon, D. M. Ceperley, and R. M. Martin, Phys. Rev. B 50, 1684 (1994).

T. Kawamura and S. Das Sarma, Phys. Rev. B 42, 3725 (1990); Phys. Rev. B 45, 3612 (1992).

F. Stern, Phys. Rev. Lett. 44, 1469 (1980).

V. Karpus, Semicond. Sci. Technol. 5, 691 (1990).