Experimental Measurements of Effective Mass in Near Surface InAs Quantum Wells

Joseph Yuan\textsuperscript{1}, Mehdi Hatefipour\textsuperscript{1}, Brenden A. Magill\textsuperscript{2}, William Mayer\textsuperscript{1}, Matthieu C. Dartailh\textsuperscript{1}, Kasra Sardashi\textsuperscript{1}, Kaushini S. Wickramasinghe\textsuperscript{1}, Giti A. Khodaparast\textsuperscript{2}, Yasuhiro H. Matsuda\textsuperscript{3}, Yoshimitsu Kohama\textsuperscript{2}, Zhuo Yang\textsuperscript{3}, Sunil Thapa\textsuperscript{4}, Christopher J. Stanton\textsuperscript{1}, and Javad Shabani\textsuperscript{1}
\textsuperscript{1}Center for Quantum Phenomena, Department of Physics, New York University, NY 10003, USA
\textsuperscript{2}Department of Physics, Virginia Tech, Blacksburg, VA 24061, USA
\textsuperscript{3}Institute for Solid State Physics, University of Tokyo, 5-1-5 Kashiwanoha, Kashiwa, Chiba 277-8581, Japan
\textsuperscript{4}Department of Physics, University of Florida, Gainesville, Florida 32611, USA
(Dated: November 11, 2019)

Near surface indium arsenide quantum wells have recently attracted a great deal of interest since they can be interfaced epitaxially with superconducting films and have proven to be a robust platform for exploring mesoscopic and topological superconductivity. In this work, we present magnetotransport properties of two-dimensional electron gases confined to an indium arsenide quantum well near the surface. The electron mass extracted from the envelope of the Shubnikov-de Haas oscillations shows an average $m^*$ = 0.04 at low magnetic field. Complementary to our magnetotransport study, we employed cyclotron resonance (CR) technique. Both $m^*$ and $g^*$ have been measured and calculated for bulk InAs [18]. It is particularly interesting that confinement of the electron wave function can strongly affect these values. Confinement becomes relevant when the 2DEG is placed near the surface, as is required for epitaxial contacts. In addition, narrow gap semiconductors can lead to strong non-parabolicity of the bands modifying the $m^*$ and $g^*$. However, to date, very few experimental studies has been performed to quantify the $m^*$ and $g^*$ in near surface InAs quantum wells. Here we report on these properties in near surface InAs quantum wells using Shubnikov-de Haas (SdH) oscillations and cyclotron resonance (CR) technique.

Wafer-scale methods for the epitaxial growth of thin films of Aluminum (Al) on Indium Arsenide (InAs) heterostructures have recently been developed which yield uniform and atomically flat interfaces [1–4]. Josephson junctions fabricated on these materials yield a gate-controllable supercurrent with highly transparent contacts between the Al top layer and an InAs quantum well (QW) directly below the surface [5–9]. Tuning of the semiconductor properties will affect supercurrent and other superconducting properties due to the wavefunction overlap at the epitaxial interface. Josephson junctions made out of Al-InAs have been used for tunable superconducting qubits, the so-called “gatemon” where the Josephson energy can be tuned in-situ with an applied electric field [10, 11]. Furthermore, since InAs has large spin-orbit coupling, they can host topological superconductivity and Majorana bound states [12–15]. The key feature in these structures is that the two-dimensional electron gases (2DEG) is confined near the surface, in close proximity to the superconductor. While the epitaxial interface creates high contact transparency, it is expected that electron mobility of the 2DEG deteriorates due to increased rates of surface scattering as compared to isolated 2DEGs buried beneath the surface [1, 16, 17]. The myriad of possible applications with this platform implores a deeper study of the characteristics and material properties for near surface InAs quantum wells.

Two important material parameters of a 2DEG are the effective mass, $m^*$, and the effective $g$ factor, $g^*$. These parameters dictate the response of a material to external electric and magnetic fields. Their effect on device performance should be accounted for in the design of mesoscopic devices and realistic theoretical modeling. Both $m^*$ and $g^*$ have been measured and calculated for bulk InAs [18]. It is particularly interesting that confinement of the electron wave function can strongly affect these values. Confinement becomes relevant when the 2DEG is placed near the surface, as is required for epitaxial contacts. In addition, narrow gap semiconductors can lead to strong non-parabolicity of the bands modifying the $m^*$ and $g^*$. However, to date, very few experimental studies has been performed to quantify the $m^*$ and $g^*$ in near surface InAs quantum wells. Here we report on these properties in near surface InAs quantum wells using Shubnikov-de Haas (SdH) oscillations and cyclotron resonance (CR) technique.

The samples were grown on a semi-insulating InP (100) substrate, using a modified Gen II molecular beam epitaxy system. The In$_x$Al$_{1-x}$As buffer is grown at low temperature to help mitigate formation of dislocations originating from the lattice mismatch between the InP substrate and higher levels of the heterostructure [19–21]. The indium content of In$_x$Al$_{1-x}$As is step-graded from $x = 0.52$ to 0.81. Next, a delta-doped Si layer of $\sim 7.5 \times 10^{11}$ cm$^{-2}$ density is placed here followed by 6 nm of In$_{0.81}$Al$_{0.19}$As. The quantum well is grown next, consisting of a 4 nm thick layer of In$_{0.81}$Ga$_{0.19}$As layer,
a 4 nm thick layer of InAs, and finally a 10 nm thick top layer of In$_{0.81}$Ga$_{0.19}$As. A thin film of Al can be epitaxially grown on the final InGaAs layer. For the transport studies of the InAs quantum wells, Al films were selectively etched by Transene type-D solution while for optical studies Al was not grown from the beginning.

The samples used for our transport measurements were patterned using photolithography. The pattern used was an L-shaped Hall bar geometry allowing simultaneous measurement of longitudinal resistances ($R_{xx}$ and $R_{yy}$) and transverse resistance ($R_{xy}$). Chemical wet etching was performed after lithographic patterning leaving a 900 nm tall mesa. A 50 nm thick aluminum oxide (Al$_2$O$_3$) gate dielectric was then deposited on top of the Hall bar via atomic layer deposition. Gate electrodes were realized by subsequent deposition of 5 nm of titanium and 70 nm of gold. All measurements were performed inside a cryogen-free refrigerator with base temperature of 1.5 K with maximum magnetic field of 12 T. Carrier densities are determined based on the slope of Hall data.

Figure 1a shows the color-scale plot of longitudinal magnetotransport, $R_{xx}$, as a function of top gate voltage, $V_G$. The Landau level fan diagram is evident from the plot with crossings observed at near $n = 1.3 \times 10^{12}$cm$^{-2}$ and 8 T and another near $n = 2.2 \times 10^{12}$cm$^{-2}$ and 12 T. At lowest densities we only observe well developed integer quantum Hall states up to $n = 1.3 \times 10^{12}$cm$^{-2}$ ($V_G < -3$ V). The first Landau level crossing appears near $V_G \sim -3$ V where it signals occupation of the second electric subband. This is most evident as $\nu = 6$ stays the same before and after the crossing in Fig. 1a. Similar Landau level crossings have been studied extensively in GaAs 2DEGs [22–25]. Three magnetotransport traces are shown in Fig. 1b-d. Longitudinal and Hall resistance as a function of magnetic field are plotted for $n = 2.2, 1.3,$ and 0.68 $\times 10^{12}$ cm$^{-2}$. The beating in SdH oscillations clearly suggest occupation of two subbands at $n = 2.2 \times 10^{12}$ where below the crossing clear quantum Hall states develop with vanishing longitudinal resistance at $n = 0.68 \times 10^{12}$.

In a non-interacting quantum Hall system, the Landau level spacing increases with magnetic field as $\hbar \omega_c$, with $\omega_c = eB/(m^*m_e)$ where B is the magnetic field, and $m_e$ is the bare electron mass. Hence, measurements of energy gaps of integer quantum Hall states should be related to electron mass. Figure 2a shows the temperature dependence of the minimum in longitudinal resistance at filling factor $\nu = 2$. The natural logarithm of the minimum in resistance in a system with parabolic bands has a linear

![Figure 1](image-url)
dependence on inverse temperature as shown in Fig. 2b [26, 27]. The energy gap is directly proportional to the magnitude of the slope. We repeated these measurements as we varied the density and hence the position of \( \nu = 2 \) in magnetic field. The results are shown in Fig. 2c where extracted energy gaps are plotted as a function of magnetic field. For comparison, we also plot the energy gap expected from \( h\omega_c \) as a black dashed line. There is a large discrepancy between the measured and expected energy gap. If we allow electron mass to be a fitting parameter we should have large enough energy gaps to be clearly observed. Their very weak presence is due to either modified \( g^* \) or Landau level broadening due to disorder. To address this and the discrepancy of energy scales for gaps in even integer quantum Hall states we next measure the temperature dependence of the low magnetic field SdH oscillations. The SdH oscillation amplitude can be isolated by subtracting the background trend of the longitudinal resistance \( R_{xx} \). Figure 3a, b displays the amplitude of SdH, \( A_{SdH} \), for densities of \( n = 3.87 \) and \( 1.2 \times 10^{12} \, \text{cm}^{-2} \). At \( n = 3.87 \times 10^{12} \, \text{cm}^{-2} \) two subbands are populated and at \( n = 1.2 \times 10^{12} \, \text{cm}^{-2} \) one subband is populated. Taking the points for a single minimum or maximum, normalized by our lowest temperature value, we can fit them to the formula \( \frac{x}{\sinh(x)} \) with \( x = 2\pi^2 T/\Delta E \), where \( T \) is the temperature and \( \Delta E \) is the gap. This allows us to calculate \( m^* = \hbar eB/(m_e \Delta E) \). Figure 3c shows the data and fit for two densities shown in 3a and 3a. We have repeated these measurements for various filling factors to extract \( m^* \) as shown in 3d. The experimental values range between 0.03 - 0.05 with an average value of \( m^* = 0.04 \). This is slightly higher than bulk values of our quantum well consisting of InAs and In_{0.81}Ga_{0.19}As with \( m^* = 0.023 \) and

![Diagram](image_url)
0.03 respectively. From the exponential envelope of the SdH oscillations we can also obtain the quantum lifetime and calculate the Landau level broadening, $\Gamma = \hbar/\tau_\nu$. Figure 3c shows $\Gamma$ for the two densities $n = 3.87 \times 10^{12}$ cm$^{-2}$ and $n = 1.22 \times 10^{12}$ cm$^{-2}$. The Landau level broadening range is close to 100 K for $n = 3.87 \times 10^{12}$ cm$^{-2}$ and 200 K for $n = 1.22 \times 10^{12}$ cm$^{-2}$. The broadening in the near surface InAs 2DEG is significantly larger than in buried InAs 2DEGs where $\Gamma$ is measured to be 5 K [21]. Here the surface scattering clearly dominates the other scattering mechanisms [1]. Thankfully, the smaller electron mass in InAs enhances the energy scales and therefore enables us to resolve quantum Hall states. Our measured Landau level broadening could qualitatively describe the large discrepancy between energy gap measurements in the quantum Hall states and $\hbar \omega_c$.

A more direct way to measure $m^*$ is through infrared CR measurements using pulsed ultrahigh magnetic fields (< 150 Tesla) generated by the single-turn coil technique [28–30]. The external pulsed magnetic field was applied along the growth direction and measured by a pick-up coil around the sample. The sample and the pick-up coil were placed inside a continuous flow helium cryostat. In this study, we employed infrared radiations from a CO$_2$ laser with wavelengths ranging from 9.2-10.6 $\mu$m. The changes in transmission through the sample were collected using a fast liquid-nitrogen-cooled HgCdTe detector. A multi-channel digitizer placed in a shielded room recorded the signals from the detector and pick-up coil.

The spin resolved CR at 10.6 $\mu$m indicated by the two arrows in Fig. 4a, separated by ~ 4 Tesla, was observed at $T = 300$ K and can be expected, as the Landau levels above the Fermi level can be occupied at $T = 300$ K, allowing the transitions between $n = 0$ and $n = 1$ for two different spins. In addition, in Fig. 4a the broad resonance at ~ 55 T represents a transition between $n = 1$ and $n = 2$ which is possi-
FIG. 4: (a) The normalized transmission of 10.6 μm excitation showing cyclotron resonance (CR) taken at T = 300 K (electron-active). The transitions indicated by arrows are attributed to the spin resolved CR transitions. (b) The CR measurement displays a sharper transitions at T = 20.5 K. Unlike the measurements at 300 K, the spin resolved CR can not be resolved but the broader resonance at 55 T (the Landau level transition n = 1 to n = 2), observed at 300 K shifts to lower fields and narrows down. (c) The effective mass m* as a function of magnetic field at T = 300 K and T = 20 K, demonstrate the non-parabolicity. (d) The absolute value of effective g-factor g* as a function of magnetic field at 20.5 K.

ble when the carrier lifetime allows time for a finite population of Landau level n = 1. This transition is not predicted from the fixed Fermi energy, but can be attributed to the non-equilibrium electron distribution [31, 32].

In Fig. 4b, we present the CR measurements at 20 K with an excitation of 10.6 μm. The spin resolved CR was not observed indicating the states above the Fermi energy are no longer occupied. On the other hand, the broad resonance observed at T = 20.5 K, and T = 300 K, is due to the transition from n = 1 to n = 2, remained and narrowed. Figure 4c summarizes our measurements for m* as a function of magnetic field at T = 300 K (crosses) and T = 20 K (filled circles). We note that although the single-coil is destroyed in each shot, the sample and pick-up coil remain intact, making it possible to carry out temperature and wavelength dependence measurements on the same sample. Figure 4c shows that the m* varied and increased monotonically with magnetic field. We measured m* = 0.04 near B = 40 T and m* = 0.061 near 70 T. Correspondingly we can estimate g* as a function of magnetic field using appropriate Landau level index using Eq. 7 in the Appendix. In Fig. 4d we show measured absolute effective g-factor at 20 K as a function of magnetic field.

Next we provide a simple theoretical model to understand m* and the Landau level fan diagram in InAs which has a non-parabolic conduction band. Unlike the wide gap semiconductors such as GaAs, CR m* and g* may vary with subband index, Landau Level index, and external magnetic field. Beginning with expectations from the bulk and introducing confinement we can arrive at expressions for m* and g* (the details are presented in the Appendix):

\[ g_{j,n} = \frac{(\varepsilon_{j,n}^+ - \varepsilon_{j,n}^-)}{\mu_B B} \]  \hspace{1cm} (1)

where \( \varepsilon_{j,n} \) is the energy of the n\textsuperscript{th} Landau level, for the j\textsuperscript{th} subband index, and at magnetic field B. Plus and minus superscript represent higher and lower Zeeman split energy bands respectively. As shown in Fig. 5a, g* depends on the subband index j, the Landau level n as well as the magnetic field B. At zero magnetic field, the absolute value of g* = 12 is reduced from bulk value of g* = 14 due to confinement and monotonically decreases as magnetic field is increased. The rate depends on the Landau level index.
Similarly one can define $m^*$ obtained by CR as:

$$m^*_{j,n} = \frac{\hbar eB}{(\xi_{j,n+1}^± - \xi_{j,n}^±)}$$

(2)

We find that $m^*$, as shown in Fig. 5b also depends on the $n^{th}$ Landau level, the $j^{th}$ subband index, and the magnetic field $B$ (we plot only the ($-$) solution for clarity). At zero magnetic field we see $m^* = 0.027$ is larger than the bulk value of $m^* = 0.023$ and increases monotonically as magnetic field is increased. These values are in close agreement with values derived from magnetotransport (over a small region 3 T to 5 T) and CR (40 T $< B < 70$ T).

**SUMMARY AND CONCLUSION**

We have done magnetotransport and ultra high field cyclotron resonance characterization of surface InAs Quantum wells. The density of these structures can be tuned and our magnetotransport measurement provides insight into the Landau level broadening and the quantum Hall energy gaps. By combining magnetotransport and cyclotron resonance measurements we can obtain conduction band effective mass $m^*$ at both low and high magnetic fields respectively. A band structure model which includes the effects of strong non-parabolicity and quantum confinement can describe the extracted $m^*$ from magnetotransport and cyclotron resonance measurements. We used our experimental CR $m^*$ values to determine the effective $g^*$ factor as a function of magnetic fields and Landau level index and these values are in a good agreement with the model presented here.

**Acknowledgment:** NYU team acknowledge partial support from US Army research office and DARPA TEE Program. Giti A. Khodaparast and Christopher J. Stanton Acknowledge support from the Air Force Office of Scientific Research under award number: FA9550-17-1-0341. Giti A. Khodaparast and Brenden A. Magill acknowledge the support from the Japanese visiting program of The Institute for Solid State Physics, The University of Tokyo. Joseph Yuan acknowledges funding from the ARO/LPS QuaCGR fellowship.

**APPENDIX: SIMPLE MODEL FOR ELECTRON MASS AND $g$-FACTOR IN A NON-PARABOLIC SEMICONDUCTOR.**

The derivation of the theoretical model accounting for non-parabolicity is described in this section. In the absence of external magnetic field (and quantum confinement) a narrow gap semiconductor such as InAs has a conduction band energy, $\epsilon$ vs. wavevector $k$ given by the dispersion relationship is given by:

$$\epsilon(1 + \alpha \epsilon) = \frac{\hbar^2 k^2}{2m^*_g} = \frac{\hbar^2 (k_x^2 + k_y^2 + k_z^2)}{2m^*_g}$$

(3)

Here, $\alpha$ is the non-parabolicity factor given by:

$$\alpha = \frac{1}{\epsilon_g}$$

(4)

with $\epsilon_g$ being the band-gap, and $m^*_g$ is the cyclotron resonance (CR) $m^*$ at the band edge ($k = 0$). For
small $\alpha \varepsilon$, the energy depends quadratically on $k$ while for large $\alpha \varepsilon$, the energy depends linearly on $k$.

In the presence of a magnetic field in the $z$ direction, it can be shown [33–35] that one can write:

$$
\varepsilon(1 + \alpha \varepsilon) = \frac{\hbar^2 k_z^2}{2 m_o^*} + \left( n + \frac{1}{2} \right) \hbar \omega_c \pm \frac{1}{2} \mu_B g_o^* B
$$

(5)

Here, $n$ is the Landau level index which can take on values $[0,1,2, \ldots]$. $\omega_c$ is the band-edge CR frequency, given by:

$$
\omega_c = \frac{eB}{m_o^*}
$$

(6)

and

$$
g_o^* = 2 \left[ 1 + \left( 1 - \frac{m}{m_c^*} \right) \frac{\Delta}{3 \varepsilon_g + 2 \Delta} \right]
$$

(7)

is the band-edge $g^*$. $\Delta$ is the valence band spin-orbit splitting, and $\mu_B$ is the Bohr-magneton given by:

$$
\mu_B = \frac{e \hbar}{2 m_e^*}
$$

(8)

Note that in the Bohr magneton, as opposed to the band-edge CR frequency, it is the bare electron mass that enters the expression.

To simplify, we set the RHS of Eq. 5 to $K$ with $j$ a positive integer and $L$ the effective width of the quantum well. Substituting into equation 5 yields:

$$
\varepsilon_{j,n}^\pm(1 + \alpha \varepsilon_{j,n}^\pm) = \frac{\hbar^2 k_z^2}{2 m_o^* L^2} + \left( n + \frac{1}{2} \right) \hbar \omega_c \pm \frac{1}{2} \mu_B g_o^* B = K
$$

(13)

We assume an effective well width of $20 \text{ nm}$. The gap at low temperatures is given by $\varepsilon_g = 0.4180$ while the spin orbit splitting is $\Delta = 0.38 \text{ eV}$ and the low temperature, band-edge $m^*$ is: $m_o^* = 0.023 m$. From Eq. 7, we see this yields a band-edge $g_o^* = -14$.

The Landau fan energies in Eq. 13 can lead us to calculate and define $g^*$ for different Landau levels by:

$$
g_{j,n} = \frac{(\varepsilon_{j,n}^+ - \varepsilon_{j,n}^-)}{\mu_B B}
$$

(14)

We can see that $g^*$ depends on the subband index $j$, the Landau level $n$ as well as the magnetic field $B$.

Similarly one can define $m^*$ by:

$$
m_{j,n}^{\pm} \equiv \frac{\hbar e B}{(\varepsilon_{j,n+1}^\pm - \varepsilon_{j,n}^\pm)}
$$

(15)

Figure 5(a,b) plots the $m^*$ and $g^*$ as a function of magnetic field and the Landau level index. We plot $m^*$ only for the lowest (-) solution. Since $g^*$ will differ between Landau levels for a non-parabolic system. The + and - effective masses will differ slightly and will lead to spin-split cyclotron resonance peaks under certain conditions. The calculation shows that in presence of non-parabolicity both of these parameters depend on the subband index $j$, the Landau level $n$, and the magnetic field $B$. We note that assuming a smaller effective quantum well width (e.g. 12 nm) will shift $m^*$ to larger values (e.g. $\sim 0.035$ at $B = 0$ T) and $g^*$ will shift smaller values ($\sim -9.5$ at $B = 0$ T).

References

[1] K. S. Wickramasinghe, W. Mayer, J. Yuan, T. Nguyen, L. Jiao, V. Manucharyan, and J. Shabani, Applied Physics Letters 113, 262104 (2018).
[35] R. Bowers and Y. Yafet, Physical Review 115, 1165 (1959).